Deexcitons in an inverted 2D semiconductor magnetoplasma

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An inverted two-dimensional electron–hole magnetoplasma with an integer carrier filling factor is analyzed. It has "deexciton" elementary excitations, i.e., exciton-like quasiparticles characterized by a momentum and a negative energy. The recombination radiation accompanying the formation of deexcitons lies in a narrow band.

Maan *et al.*[1,2] have recently observed and studied magnetooscillation luminescence spectra of a new type. The measurements were carried out on high-quality GaAs/$Ga_{1-x}Al_xAs$ quantum wells at a pump level so high that up to 13 Landau levels were filled in a strong magnetic field $H = 18$ T. These levels were filled by electrons in the conduction band (the c band) and by holes in the valence band (the v band). Electron–hole correlations in a 2D gas are manifested in the emission spectrum of this highly inverted plasma, which exhibits a clearly defined structure. Similar spectra are observed[3] in InGaAs/InP quantum wells. Spectroscopy of electron correlations at $H = 0$ does not hold much promise, since the spectra are smooth. Magnetospectroscopy of a highly inverted gas opens up some new opportunities here.

A theory is derived below for the case of strong magnetic fields,

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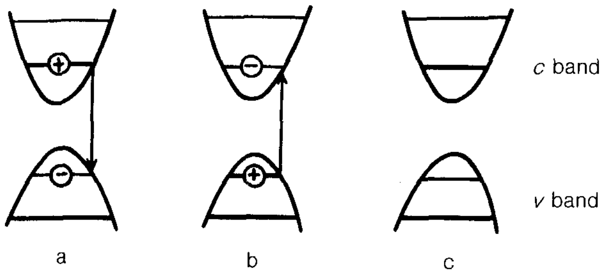

FIG. 1. Exciton and diexciton. The lower Landau level is inverted ($N=0$). The heavy lines represent the filled Landau levels and the thin lines denote the empty levels: (1) diexciton corresponding to an allowed transition $c0\to v0$ in the emission spectrum of the inverted system. The finite state with one diexciton is shown; (b) magnetic exciton corresponding to the $v0\to c0$ transition in the absorption spectrum of an unexcited crystal. This magnetic exciton is the reciprocal of the diexciton in Fig. 1a; (c) the initial state of the inverted system which precedes the formation of a diexciton.

$\gamma\equiv\omega_c/(e^2/\epsilon l)\gg 1$, where $\omega_c(H)$ is the electron cyclotron frequency, $l(H)$ is the magnetic length, and $\hbar=1$; in this case, the concept of 2D magnetoexcitons is effective.[4–7] In contrast with the case of an equilibrium system, the exciton-like formations which are peculiar to an inverted system are "deexcitations": When they form, they move the system toward an equilibrium state. In this sense one can speak of "deexcitons." The basic purpose of this letter is to derive a theory of elementary excitations in an inverted magnetized plasma in these terms. We assume a temperature $T=0$ and an integer filling of the Landau levels in the c and v bands. The electrons are assumed to have no spin.

The underlying physics can be explained most easily under the assumption that the number of inverted levels, $N+1$ ($N=0,1,\ldots$), is small. Figure 1a illustrates the case in which only the lowest-lying level, $N=0$, is inverted. A deexciton consists of a Fermi hole in the c band and an electron in the v band. This neutral formation is stable and has a momentum $\mathbf{k}$, by analogy with the corresponding exciton (Fig. 1b). This deexciton forms in an inverted system (Fig. 1c) through the radiative recombination of an electron–hole pair. The selection rule $\mathbf{k}=0$ means that the emission must lie in a narrow band; the Coulomb interaction does not broaden the emission band, simply shifting it. Since the formation of a deexciton is associated with a decrease in the energy of the system, the energy of the deexciton is negative:

$$E^d(\mathbf{k})=-E_G-(\omega_c+\omega_v)/2+\epsilon(\vec{k})\equiv E^0+\epsilon(\vec{k})<0, \quad (1)$$

where ω_v is the hole cyclotron frequency, E_G is the gap, and $\epsilon(\mathbf{k})$ is the Coulomb energy. Figure 2a shows a stable deexciton in the case of two inverted levels ($N=1$). A theory for such deexcitons can be derived rigorously in terms of the parameter $\gamma^{-1}\ll 1$, as has been done for the theory of magnetoexcitons.[5–7] The situation involving the deexciton in Fig. 2b is more complicated at $N=1$. If the Coulomb interaction is ignored, the configuration in Fig. 2b is energetically degenerate with respect to the two other configurations shown in Figs. 2c and 2d. These other configurations have two excitonic excitations. In Fig. 2c it is a deexciton with quantum numbers $(v0,c1)$

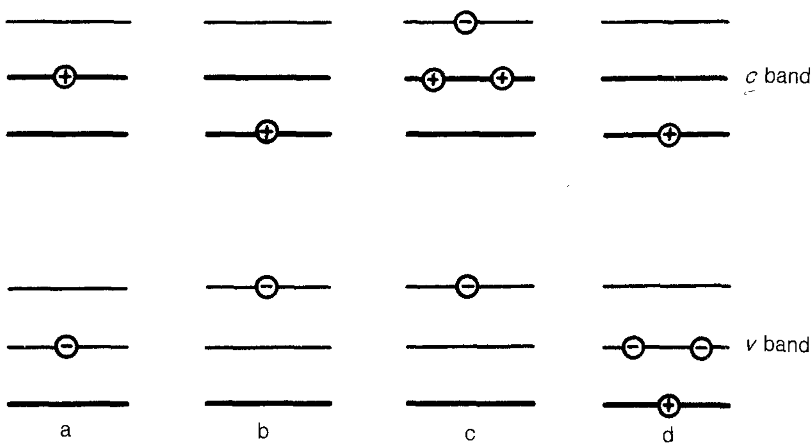

FIG. 2. Mutually degenerate electronic configurations. Two lower Landau levels are inverted ($N=1$). The light and heavy lines are the same in Fig. 1: (a) diexciton ($v1$, $c1$); (d) diexciton ($v0$, $v0$); (c) diexciton ($c0$, $c1$) and magnetic plasmon ($c2$, $c1$); (d) ($v1$, $c0$) and magnetic plasmon ($v1$, $v2$). The last three configurations are mutually degenerate and become strongly mixed as a result of Coulomb interaction.

and an intraband exciton ($c\,2$,$c1$) (magnetoplasmon). Depending on whether there exists only a two-exciton continuum or also a bound state of the type in Fig. 2b, the emission spectrum does not or does contain narrow deexciton bands in addition to a two-frequency continuum. At $N>1$, the situation quickly becomes more complicated, but a deexciton (vN,cN) is stable at arbitrary N.

The following operator creates an exciton which has a momentum **k** and which is formed by an electron in level cm and a hole in level vn:

$$A^{+}_{cm,vn}(\mathbf{k}) = (2\pi/S)^{1/2}\sum_{p} a^{+}_{cm}(p)a_{vn}(p+k_x)\exp\{ik_x(p+k_y/2)\}. \tag{2}$$

Here a_{cm} and a_{vn} are electron operators, and S is the area of the sample in units of l^2; below we set $l=1$. If a hole forms in the c band, a_{vn} is replaced by a_{cn}.

The matrix element of the electron–electron interaction, including the Coulomb and exchange contributions, is

$$V_{mn}(q) = V(q)w_m(q^2)w_n(q^2) - \int V(q_1)w_{mn}(q_1^2)\exp(i\mathbf{q}\mathbf{q}_1)d\mathbf{q}_1/2\pi\,. \tag{3}$$

Here $V(q)=2\pi/q$ is a Fourier component of the Coulomb potential $V(r)$;

$$w_m(q^2) = \exp(-q^2/4)L_m(q^2/2);$$

$$w_{mn}(q^2) = (n!/m!)(q^2/2)^{m-n}\exp(-q^2/2)[L_n^{m-n}(q^2/2)]^2,\quad m\geq n; \tag{4}$$

L_n^m are the Laguerre polynomials; and $w_{mn}=w_{nm}$. If states m and n belong to different bands, the second term in (3) should be omitted. We introduce the energy of an elementary excitation created by the operator $A^{+}_{cm,vn}$ (an exciton), $A^{+}_{vn,cm}$ (a deexciton), or $A^{+}_{cm,cn}$ (a plasmon). This energy is defined as an expectation value of the type $\langle AA^{+}\rangle$ in terms of the corresponding vacuum. It contains E^0, the distance between the levels in the case $V(q)=0$ (for a deexciton, $E^0_{vn,cm} = -E^0_{cm,vn} < 0$), and a Coulomb part $\epsilon(\mathbf{k})$, which consists of a k-dependent term

$$\epsilon_{mn}(\mathbf{k}) = -\int V_{nm}(q)\exp(i\mathbf{qk})/d\mathbf{q}/(2\pi)^2 \tag{5}$$

and a constant component

$$\epsilon^{(0)}_{mn} = \int V(q)w_n^2(\vec{q}^{\,2})d\vec{q}/(2\pi)^2 + \sum_s V_{ms}(q=0)/2\pi - \sum_s V_{ns}(q=0)/2\pi. \tag{6}$$

The first term in (6) is the change in the exchange energy in level n (cn or vn) when an electron is removed from this level. The sums over s are extended to all filled levels in the c and v bands, except the level from which the electron is removed. For an exciton and a deexciton, these sums differ in magnitude because of a difference in the original filling of the levels.

After several manipulations, we find the following expression for the deexciton energy:

$$E^d_{vn,cm}(\mathbf{k}) = E^0_{vn,cm} - \int V(\mathbf{r}+\mathbf{k})w_{mn}(\mathbf{r}^2)d\mathbf{r}/2\pi$$
$$+\int V(r)[w_n(r^2)+w_m(r^2)]w_N^{(1)}(r^2)\,d\mathbf{r}/2\pi, \tag{7}$$

where $w_N^{(1)}$ is found from W_N by replacing L_N by L^1_N. In the important case $m=n=N$ we have

$$E^d_{vN,cN}(\mathbf{k}=0) = E^0_{vN,cN} + \int V(r)w_N(r^2)[2w_N^{(1)}(r^2) - w_N(r^2)]d\vec{r}/2\pi. \tag{8}$$

Here and below, E^0 incorporates the renormalization of the gap E_G caused by the interaction of the hole with all electrons in the valence band. The energy of the exciton in the noninverted system is

$$E^{ex}_{cm,vn}(\mathbf{k}) = E^0_{cm,vn} - \int V(\mathbf{r}+\mathbf{k})w_{mn}(r^2)d\mathbf{r}/2\pi. \tag{9}$$

One can show that

$$|E^d_{vN,cN}(\mathbf{k}=0)| \le E^{ex}_{cN,vN}(\mathbf{k}=0).$$

The difference is zero at $N=0$ and increases as $(4/\pi)(2N)^{1/2}(e^2/\epsilon l)$ as $N\to\infty$. This difference determines the difference between the frequencies of a stable deexciton and of the corresponding exciton. These conclusions agree qualitatively with the results of Ref. 1.

In the energy region corresponding to a deexciton $(v\,0,c\,0)$, it is not possible to find the spectrum in its general form, because of the contribution from the configurations in Figs. 2c and 2d. This spectrum is found from an integral equation for a biexciton which generalizes Eq. (34) of Ref. 6. The nature of the problems that arise here can be illustrated in the $N=1$ case. Using (7) for the energy of a deexciton and (5) and (6) for the energy of a magnetoplasmon, one can show that in the case $N=1$ we have

$$E^{d}_{v0,c1}(\mathbf{k})+E^{ex}_{c2,c1}(-\mathbf{k})=E^{0}_{v0,c0}+(e^2/\epsilon l)\begin{cases}(9/4)\sqrt{\pi/2}+k, & k\to 0,\\ (51/16)\sqrt{\pi/2}-2/k, & k\to\infty,\end{cases} \tag{10}$$

and that the entire two-particle spectrum, with a zero total momentum, lies within the limits which follow from (10). According to (7), the approximate value of the energy of a deexciton $(v\,0,c\,0)$ (Fig. 2b) is

$$E^{d}_{v0,c0}(\mathbf{k}=0)=E^{0}_{v0,c0}+2\sqrt{\pi/2}(e^2/\epsilon l). \tag{11}$$

This energy lies outside the two-particle spectrum. The interaction of quasiparticles leads to a repulsion of the one-particle levels from the boundaries of the two-particle spectrum, which remain unshifted. At $N=1$, the one-particle level of the type $(v\,0,c\,0)$ is therefore stable with respect to two-particle decays. However, level (11) falls in a three-particle continuum: a deexciton $(v1,c1)$ plus two magnetoplasmons, one in the c band and the other in the v band. Whether the shift of the level in (11) caused by the interactions is sufficient for its stabilization can be determined only by solving the corresponding equation. It appears that deexcitons are usually metastable. Expression (7) is rigorous in terms of the parameter $\gamma^{-1}\ll 1$ only in the case $m=n=N$. In other cases, the contribution from configurations of the types in Figs. 2c and 2d, complicates the situation dramatically. The problem can be simplified by assuming a pronounced deviation from a parabolic shape: $\gamma'\equiv\Delta/(e^2/\epsilon l)\gg 1$, were Δ is the difference between the energies of the configurations in Figs. 2b, 2c, and 2d. We assume $\Delta\ll\omega_c$. The parameter $\gamma'\gg 1$ leads to the existence of one-particle deexciton states which generate a narrow-band omission spectrum. This parameter makes it possible to carry out a perturbation-theory calculation of the intensity of the continuum corresponding to transitions to two-particle states. The admixture of these states in the deexciton state (vn,cn), with momentum $\mathbf{k}=0$ and approximately the same energy, is

$$2(2\pi S)^{-1/2}\sum_{q}\sum_{ms}(V(q)/\Delta)J_{m,m-s}(\vec{q})J_{n,n+s}(-\mathbf{q})A^{+}_{vn,cn+s}(\mathbf{q})\,A^{+}_{cm,cm-s}(-\mathbf{q})\,|0> \tag{12}$$

plus the corresponding expression for the v band. Here $n<N$, $m>N$, $n+s$, $m-s\leqslant N$, and

$$J_{m,m-s}(\mathbf{q})=[(m-s)!/m!]^{1/2}\exp(-q^2/4)[(iq_x+q_y)/\sqrt{2}]^{s}L^{s}_{m-s}(q^2/2), \tag{13}$$

where $J_{n,m}(\mathbf{q})=J_{m,n}(q_x,-q_y)\ldots$. It follows from (12) that the integral intensity of the emission due to transitions to a two-exciton state $(vn,\,cn+s,\,cm,\,cm-s)$ is, in units of the intensity of the allowed transition to the (vn,cn) state,

$$(2\pi\Delta^2 2^{2s-2})^{-1}[n!(m-s)!/(n+s)!m!]$$

$$\times \int V^2(q) q^{4s} e^{-q^2} [L^s_{m-s}(q^2/2) L^s_n(q^2/2)]^2 \, d\mathbf{q}/(2\pi)^2. \tag{14}$$

With $V = 2\pi/q$, the integral can be evaluated exactly, but the triple sum that arises is very complicated. In the $n = 0$ case, it reduces to the single sum

$$\frac{2(e^2/\epsilon l)^2}{s\Delta^2} \sum_l (-)^{m-s-l} 2^{-2(s+l)} (2(s+l)-1)!/$$

$$[l!(s+l)!(m-s-l)!(2s+l-m-1)!]^{-1}.$$

In the case $m \gg 1$, its s dependence is determined by $s^{-3/2}(m-s)^{-1/2}$, where $s, m-s \gg 1$. In other words, transitions involving small changes in quantum numbers are predominant among the Auger processes.

We have discussed only processes which originate within one energy shell. In other words, we have ruled out an interaction of configurations whose energies differ by an amount on the order of ω_c. Among the Auger processes associated with this interaction is the creation of a stable deexciton of the type in Fig. 1a, in a process accompanied by the excitation of a plasmon. Their intensity is suppressed in accordance with the parameter $\gamma^{-1} \ll 1$, but these are the processes which are responsible for the long-wave tail in the emission.

In summary, we have shown that Coulomb correlations in the final state, which lead to the formation of deexcitons, give rise to the existence of narrow bands in the spectrum of recombination radiation from an inverted 2D magnetoplasma.

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