Photoluminescence from fractional quantum Hall systems: Role of separation between electron and hole layers

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The photoluminescence (PL) spectrum of a two-dimensional electron gas (2DEG) in the fractional quantum Hall regime is studied as a function of the separation d between the electron and valence hole layers. The abrupt change in the response of the 2DEG to the optically injected hole at d of the order of the magnetic length λ results in a complete reconstruction of the PL spectrum. At $d<\lambda$, the hole binds one or two electrons to form neutral (X) or charged (X^-) excitons, and the PL spectrum probes the lifetimes and binding energies of these states rather than the original correlations of the 2DEG. At $d>2\lambda$, depending on the filling factor ν, the hole either decouples from the 2DEG to form an ''uncorrelated'' state h or binds one or two Laughlin quasielectrons (QE's) to form fractionally charged excitons hQE or $h\mathrm{QE}_2$. The strict optical selection rules for bound states are formulated, and the only optically active ones turn out to be h, hQE* (an excited state of the dark hQE), and $h\mathrm{QE}_2$. The ''anyon exciton'' $h\mathrm{QE}_3$ suggested in earlier studies is neither stable nor radiative at any value of d. The critical dependence of the stability of different states on the presence of QE's in the 2DEG explains the observed anomalies in the PL spectrum at $\nu=\frac{1}{3}$ and $\frac{2}{3}$.

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I. INTRODUCTION

The optical properties of quasi-two-dimensional (2D) electron systems in high magnetic fields have been extensively studied in the recent years both experimentally[1–20] and theoretically.[21–38] In symmetrically doped quantum wells (QW), where both conduction electrons and valence holes are confined in the same 2D layer, the photoluminescence (PL) spectrum of a two-dimensional electron gas (2DEG) probes the binding energy and optical properties of neutral and charged excitons (bound states of one or two electrons and a hole, $X=e$-h and $X^-=2e$-h), rather than the original correlations of the 2DEG itself. The experiments[10–20] and theory[30–36] agree that the X^- can exist in the form of a number of different bound states, whose binding energies depend strongly on the well width w and composition, magnetic field B, etc., but (at least in dilute systems) much less on the electron filling factor ν. In particular, the only bound X^- state that occurs at zero or low B is the optically active singlet[30,35,36] X_s^-, while more bound states form at higher B. Of these states, one is observed in PL,[12–14,20] and it has only recently been identified[35] as an excited ''bright'' triplet X_{tb}^-. The lowest energy ''dark'' triplet X_{td}^- has been predicted earlier,[31] but it is expected to have very long optical lifetime[32] and its recombination has not yet been detected experimentally.[14,35,36]

The PL spectra containing more information about the original electron correlations of the 2DEG are obtained in asymmetrically doped wide QW's or heterojunctions, where the spatial separation d of electron and hole layers weakens the e-h interaction.[24] Unless d is smaller than the magnetic length λ, the PL spectra of such bilayer systems show no recombination from X^- states. Instead, they show anomalies[1–7] at the filling factors $\nu=\frac{1}{3}$ and $\frac{2}{3}$ at which Laughlin incompressible fluid states[39] are formed in the 2DEG and the fractional quantum Hall (FQH) effect[40] is observed in transport experiments.

The present paper is a continuation of our earlier work[29] where we studied the energy spectra of 2D fractional quantum Hall systems in the presence of an optically injected valence hole. There, we have identified the possible bound states in which a valence hole can occur. They included the ''uncorrelated'' state h in which the free hole moves in the rigid electron Laughlin fluid at a local filling factor $\nu=\frac{1}{3}$, and the fractionally charged exciton (FCX) states hQE and $h\mathrm{QE}_2$ in which the hole binds one or two Laughlin quasielectrons (QE's). The charge neutral ''anyon exciton'' state[27] $h\mathrm{QE}_3$ was found unstable at any value of d. Here, we give a detailed analysis of the optical properties of these states and explain the features observed in the PL spectra of bilayer systems. Based on the analysis of the involved dynamical symmetries[41,42] (those of charged particles moving in a translationally invariant space and in a perpendicular magnetic field), we formulate the optical selection rules for the FCX complexes. These rules are verified in exact numerical calculations for finite systems in Haldane's spherical geometry[43,44] (using Lanczos-based algorithms[45] we are able to calculate the exact spectra of up to nine electrons and a hole at $\nu\approx\frac{1}{3}$). It turns out that the only radiative bound states involving the hole are h, hQE* (an excited state of hQE), and $h\mathrm{QE}_2$, and that emission from both hQE and $h\mathrm{QE}_3$ is forbidden. The fact that the previously suggested[24] recombination from a h-QE pair state can only occur through the excited state hQE* diminishes the importance of this

process at low temperatures. The result that the $h\mathrm{QE}_3$ complex (or any of its excitations) is neither stable nor radiative questions applicability of the theory of ''anyon excitons'' put forward by Rashba and Portnoi[27] to account for the anomalies observed in the PL spectra at $\nu=\frac{1}{3}$ and $\frac{2}{3}$. Instead, these anomalies are explained in terms of emission from the competing *radiative* bound states, $h\mathrm{QE}^*$ and $h\mathrm{QE}_2$, and from an uncorrelated hole state h.

II. MODEL

The model considered here is identical to that of Ref. 29. A 2DEG in a strong magnetic field B fills a fraction $\nu<1$ of the lowest Landau level (LL) of a narrow QW, whose width w we set to zero. A small number ($\nu_h \ll \nu$) of valence holes are optically injected into a parallel 2D layer of width $w_h=0$, separated from the electron layer by a distance d. The single-particle states $|m\rangle$ in the lowest LL are the eigenstates of the orbital angular momentum, $m=0, -1, -2, \ldots$ for the electrons and $m_h=-m=0, 1, 2, \ldots$ for the holes. Since $\nu_h \ll \nu$ and the strongly bound complexes at large B involve only one hole, it is enough to study the many-electron–one-hole Hamiltonian which can be written as

$$H=\sum_{ijkl}\,(c_i^\dagger c_j^\dagger c_k c_l V_{ijkl}^{ee}+c_i^\dagger h_j^\dagger h_k c_l V_{ijkl}^{eh}), \qquad (1)$$

where $c_m^\dagger$ ($h_m^\dagger$) and c_m (h_m) create and annihilate an electron (hole) in state $|m\rangle$. The constant energy of the lowest LL is removed from H, which hence includes only the e-e and e-h interactions whose two-body matrix elements V^{ee} and V^{eh} are defined by the intra- and interlayer Coulomb potentials, $V_{ee}(r)=e^2/r$ and $V_{eh}(r)=-e^2/\sqrt{r^2+d^2}$. At $d=0$, the e-h matrix elements are equal to the e-e exchange ones, $V_{ijkl}^{eh}=-V_{ikjl}^{ee}$, and at $d>0$ the e-h attraction at short range is reduced. The convenient units of length and energy are the magnetic length λ and the energy e^2/λ.

The 2D translational invariance of H results in the conservation of two orbital quantum numbers: the projection of total angular momentum $\mathcal{M}=\Sigma_m(c_m^\dagger c_m-h_m^\dagger h_m)m$ and an additional angular momentum quantum number $\mathcal{K}$ associated with partial decoupling of the center-of-mass motion of an e-h system in a homogeneous magnetic field.[41,42] For a system with a finite total charge, $\mathcal{Q}=\Sigma_m(h_m^\dagger h_m-c_m^\dagger c_m)e\neq 0$, the partial decoupling of the center-of-mass motion means that the energy spectrum consists of degenerate LL's.[41] The states within each LL are labeled by $\mathcal{K}=0, 1, 2, \ldots$ and all have the same value of $\mathcal{L}=\mathcal{M}+\mathcal{K}$. Since both $\mathcal{M}$ and $\mathcal{K}$ (and hence also $\mathcal{L}$) commute with the PL operator $\mathcal{P}$, which annihilates an optically active (zero-momentum, $k=0$) e-h pair (exciton), $\mathcal{M}$, $\mathcal{K}$, and $\mathcal{L}$ are all simultaneously conserved in the PL process.

The 2D symmetry of a planar system is preserved in the finite size, N-electron–one-hole (Ne-h) calculation in Haldane's geometry,[43] where all particles are confined to a spherical surface of radius R and the radial magnetic field is produced by a Dirac monopole. The conversion of the numerical results between the spherical and planar geometries follows from the exact mapping[35,46] between the planar quantum numbers $\mathcal{M}$ and $\mathcal{K}$, and the 2D algebra of the total angular momentum $\mathbf{L}$ on a sphere. The detailed description of the Haldane sphere model can be found elsewhere[43,44,47] (see also Refs. 33–35 for the application to e-h systems). The strength $2S$ of the magnetic monopole is defined in the units of flux quantum $\phi_0=hc/e$, so that $4\pi R^2 B=2S\phi_0$ and the magnetic length is $\lambda=R/\sqrt{S}$. The single-particle states are the eigenstates of angular momentum $l\geqslant S$ and its projection m, and are called monopole harmonics. The single-particle energies fall into degenerate angular momentum shells (LL's). The lowest shell has $l=S$, and thus $2S$ is a measure of the system size through the LL degeneracy. The charged many-electron–one-hole states form degenerate total angular momentum (L) multiplets (LL's) of their own. The total angular momentum projection L_z labels different states of the same multiplet just as $\mathcal{K}$ or $\mathcal{M}$ did for different states of the same LL on a plane. Different multiplets are labeled by L just as different LL's on a plane were labeled by $\mathcal{L}$. The pair of optical selection rules on the sphere, $\Delta L_z=\Delta L=0$ (equivalent to $\Delta\mathcal{M}=\Delta\mathcal{K}=0$ on a plane) results from the fact that an optically active exciton has zero angular momentum.

III. BOUND STATES

A. Small layer separation

Depending on the separation d between the electron and hole layers, different bound states can occur in an e-h system.[29] In the ''strong-coupling'' regime, at d less than about 1.5λ, the interaction between the hole and the electrons is stronger than the characteristic correlation energy of the 2DEG. The response of the 2DEG to the optically injected valence hole occurs through spontaneous creation of charge excitations (QE-QH pairs) which completely screen its charge. As a result, the original e-e correlations of the 2DEG are locally (in the vicinity of the hole) replaced by (stronger) e-h correlations. These new correlations are most conveniently described in terms of two types of new quasiparticles formed in the system, neutral (X) or charged (X^-) exciton states, in which the hole binds one or two electrons, respectively.

In an ideal system with no LL mixing and $w=w_h=d=0$, the ''dark'' (nonradiative; see Sec. IV) triplet charged exciton X_{td}^- is the only bound state (other than the neutral exciton X) that is stable in the presence of the surrounding 2DEG (Refs. 31, 32, and 35) (e.g., $X_2^-+e\rightarrow 2X^-$ for the charged biexciton). The X_{td}^- unbinds[29,32] at $d\approx\lambda$, and a different X^- state, a dark singlet X_{sd}^-, forms[29,48] at $0.4\lambda\leqslant d\leqslant 1.5\lambda$. In more realistic systems, when the effects due to LL mixing, finite widths of electron and hole layers, and their finite separation are taken into account, a few other bound X^- states occur.[35] Most important of these states are the bright singlet X_s^- and the bright triplet X_{tb}^-. The four X^- states are distinguished by the total electron spin J, and the total angular momentum L: the X_s^-, X_{tb}^-, X_{td}^-, and X_{sd}^- states have $J=0, 1, 1, 0$, and $L=S, S, S-1, S-2$, respectively (on a plane, $L=S$, $S-1$, and $S-2$ correspond to $\mathcal{L}=0$, -1, and -2, respectively). The binding energies of X^-

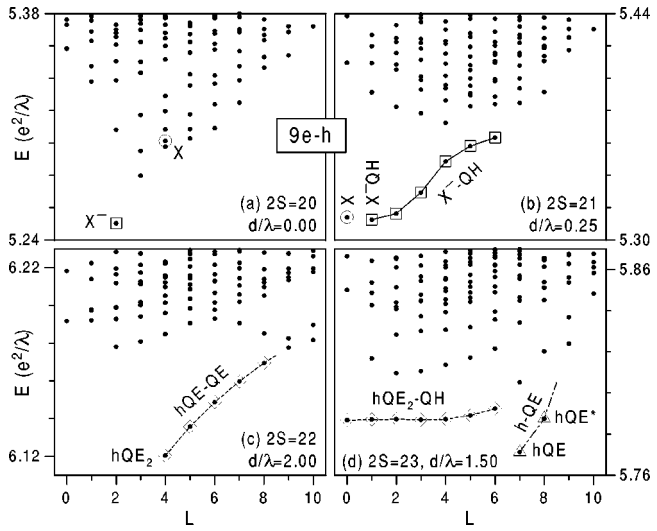

FIG. 1. The energy spectra (energy E vs angular momentum L) of the $9e$-h system calculated on a Haldane sphere with different monopole strengths $2S$ and at four different layer separations d: (a) $2S=20$ and $d=0$, (b) $2S=21$ and $d=0.25\lambda$, (c) $2S=22$ and $d=2\lambda$, (d) $2S=23$ and $d=1.5\lambda$. Different symbols and lines mark states and bands containing different quasiparticles: circles indicate X, squares indicate X^-, diamonds indicate $h\mathrm{QE}_2$, triangles indicate $h\mathrm{QE}$ and $h\mathrm{QE}^*$. λ is the magnetic length.

states depend strongly on B and d. In a narrow ($w\approx w_h \approx 10$ nm) and symmetric ($d=0$) GaAs QW, the X_s^- is the most strongly bound X^- state at B smaller than about 30 T, and at larger B, the X^- ground state changes to the X_{td}^- (the X_{tb}^- has always smaller binding energy than both X_s^- and X_{td}^-). At $d>0$, the binding energy of the X_s^- is reduced more than that of the two triplets, and the critical value of B at which the singlet-triplet crossing occurs is significantly decreased.[49]

It has not been clearly spelled out until recently[29,50] that only the ''decoupled''[21–23] $k=0$ state of the charge neutral X exists in the 2DEG. The X at $k>0$ has a finite electric dipole moment (proportional to k) whose strong interaction with the underlying 2DEG leads to the binding of the second electron and the formation of an X^-. The numerical calculations show[29,50] that the low-lying band of e-h states at $L>0$, previously interpreted[23,25,26] as the dispersion of a charge neutral ''dressed exciton'' (an X with $k>0$ coupled to the QE-QH pair excitations of the 2DEG), in fact describe an X^-.

As an example, in Figs. 1(a) and 1(b) we present the $9e$-h energy spectra at $2S=20$ and 21 and at small layer separation ($d=0$ and 0.25λ, respectively) in which the lowest energy states have been identified as containing an X^- state (this is the ''dark'' triplet state X_{td}^- bound in the lowest LL) Laughlin correlated with the remaining seven electrons and an X decoupled from the remaining eight electrons. While a detailed discussion of these states has been given elsewhere,[29] let us only note that the low-energy band of states connected with a line in Fig. 1(b) describes an X_{td}^- interacting with a QH of the two-component e-X_{td}^- incompressible fluid with Laughlin correlations[39,51] and not the ''dressed exciton'' dispersion (the angular momenta $L=1$, 2, . . . , 6 of this band can be predicted from a generalized, two-component e-X^- composite fermion picture[34]).

It is quite remarkable that the wave function of the 2DEG in the strong-coupling regime can be well represented in terms of wave functions of competing bound X and X^- states, neglecting the distortion of these states due to the coupling to the surrounding electrons. For the X^- states, this is a consequence of the short range[52] of the e-X^- repulsion that results in Laughlin correlations and the effective isolation[35] of the X^- states from the 2DEG. For the X state at $k=0$ (whose charge and electric dipole moment vanish), this is a result of weak (zero at $d=0$) coupling to the 2DEG.

B. Large layer separation

At d larger than about 1.5λ, in the ''weak-coupling'' regime, the e-h attraction becomes too weak compared to the characteristic e-e correlation energy, its range becomes too large compared to the characteristic e-e separation, and the X and X^- states unbind. In this regime, the perturbation associated with the potential of the optically injected hole does not cause the reconstruction of the e-e (Laughlin) correlations of the 2DEG, whose response involves only the existing Laughlin QE's. Since no additional QE-QH pairs are spontaneously created to screen the hole charge, a discontinuity occurs at the Laughlin fillings such as $\nu=\frac{1}{3}$. At $\nu \leqslant \frac{1}{3}$, no QE's that could bind to the hole occur in the 2DEG, the existing QH's are repelled from it, and the electrons in the vicinity of the hole form a Laughlin state with the (local) filling factor $\nu=\frac{1}{3}$. In this ''uncorrelated'' state, the hole causes no (local) response of the 2DEG. At $\nu>\frac{1}{3}$, the hole binds one or two QE's to form fractionally charged excitonic states $h\mathrm{QE}$ or $h\mathrm{QE}_2$ (it has been shown[29] that the charge neutral ''anyonic excitons'' $h\mathrm{QE}_3$ are unstable at any value of d). Just as in the case of the X or X^- at small d, the $h\mathrm{QE}_n$ states are well-defined quasiparticles of the e-h system at larger d, and they can be attributed such single-particle properties as the binding energy Δ, angular momentum L, PL energy ω, and oscillator strength τ^{-1}, etc. Because of their low density, the $h\mathrm{QE}_n$ quasiparticles can be to a good approximation regarded as noninteracting, free particles moving in a ''rigid'' Laughlin $\nu=\frac{1}{3}$ reference state.

While the binding of $h\mathrm{QE}_n$ states as a function of d and ν has been discussed in great detail in Ref. 29, in Figs. 1(c) and 1(d) we present the $9e$-h energy spectra at $2S=22$ and $d=2\lambda$ (c), and at $2S=23$ and $d=1.5\lambda$ (d), in which the lowest energy states contain the $h\mathrm{QE}$, $h\mathrm{QE}^*$ (the first excited state of the h-QE pair), and $h\mathrm{QE}_2$ complexes. As for the X and X^- states in the strong-coupling regime, it is quite remarkable that the complicated correlations of a many-body e-h system at larger d can be well represented in terms of rather simple and well-defined free $h\mathrm{QE}_n$ quasiparticles.

IV. OPTICAL SELECTION RULES

A number of different selection rules govern the optical recombination of bound e-h complexes. In general, any symmetry resulting in a conservation of a certain quantum num-

ber $\mathcal{W}$ results in a strict selection rule

$$\Delta\mathcal{W}=\text{const} \tag{2}$$

if the commutator between $\mathcal{W}$ and the PL operator

$$\mathcal{P}=\sum_m (-1)^m c_m h_m \tag{3}$$

that annihilates an optical ($k=0$) exciton (on a Haldane sphere) is proportional to $\mathcal{P}$.

The so-called ''hidden symmetry''[21–23] is the exact particle-hole symmetry in the lowest LL of narrow QW's in which electrons and valence holes are confined to the same layer (equal widths, $w=w_h$, and zero separation, $d=0$). As a result, the optically active ($k=0$) excitons annihilated by $\mathcal{P}$ decouple from the excess electrons. The $k=0$ exciton is the only radiative bound state of an e-h system, and the emission from the so-called ''multiplicative'' (MP) many-body states that contain a number (N_X) of $k=0$ excitons occurs at the bare exciton energy (independent of the electron density) and follows the $\Delta N_X=-1$ selection rule. Because of the exciton decoupling, all bound states other than the exciton (e.g., the triplet charged exciton X_{td}^-) have $N_X=0$ and cannot recombine. The hidden symmetry holds only to some extent in realistic systems, where the asymmetry ($w\neq w_h$) and separation ($d>0$) of electron and hole layers, as well as the asymmetric LL mixing (due to different electron and hole cyclotron energies) result, for example, in the binding of the radiative singlet (X_s^-) and triplet (X_{tb}^-) charged exciton states.

The 2D translational/rotational symmetry results in the conservation of $\mathcal{M}$ and $\mathcal{K}$ (or L_z and L on a sphere) in the emission process.[33–35,42] The $\Delta\mathcal{M}=\Delta\mathcal{K}=0$ (or $\Delta L_z=\Delta L=0$) selection rules hold strictly when applied to the entire e-h system or to an isolated bound state. These rules (independently from the $\Delta N_X=-1$ rule) forbid emission from the X_{td}^- state that has $\mathcal{L}=\mathcal{M}+\mathcal{K}=-1$ (or $L=S-1$), while the electron left in the final state has $\mathcal{L}=0$ (or $L=S$). For bound states coupled to the surrounding 2DEG (or to any QW imperfections that break the translational symmetry) these selection rules are only approximate, and the strength of the optical transitions from otherwise nonradiative states is a measure of the distortion of these states due to their coupling to the 2DEG. We have shown[35] that the e-X^- Laughlin correlations limit high-energy e-X^- collisions in dilute ($\nu\leqslant\frac{1}{3}$) systems, and thus that the approximate selection rules remain valid for the X^- states formed in the 2DEG.

Yet another set of selection rules are associated with the electron and hole spin degrees of freedom. If the absence of the mixing of valence subbands, the total electron and heavy hole spins, J and J_h, and projections, J_z and J_{zh}, are all conserved by H. The recombination events must obey $\Delta J_z=\mp\frac{1}{2}$ and $\Delta J_{zh}=\pm\frac{3}{2}$, and the two types of transitions with $\Delta(J_z+J_{zh})=\pm 1$ correspond to two different polarizations of emitted light. In the presence of the valence subband mixing, the spin of the hole is coupled to the hole orbital angular momentum and, through the Coulomb interaction, to the orbital angular momentum of the electron. We assume here that the subband mixing can be neglected and that all electron and hole spins are polarized by a large Zeeman energy, so that the spin selection rules are always obeyed.

V. PHOTOLUMINESCENCE OF NEUTRAL AND CHARGED EXCITONS AT SMALL LAYER SEPARATION

A. Laughlin correlated e-X^- liquid

In narrow QW's ($w\leqslant 20$ nm), the X's decouple and the X^- with the remaining electrons form a two-component incompressible fluid with Laughlin e-X^- correlations.[33,34] Laughlin correlations mean that a number of e-X^- pair eigenstates that correspond to the smallest average e-X^- separation (on a sphere, these are the states with maximum L; on a plane, these are the ones with minimum relative angular momentum) are completely avoided.[52] The avoiding of a number (m_{eX^-}) of highly repulsive e-X^- pair states is described by a Jastrow prefactor $\Pi_{ij}(z_e^{(i)}-z_{X^-}^{(j)})^{m_{eX^-}}$ in the wave function (which leads to a generalized, two-component composite fermion model[34]). This is equivalent to saying that the high energy e-X^- collisions do not occur, and that the X^-'s are effectively isolated from the 2DEG. The isolation of the X^- states is even enhanced at $d>0$ where the perpendicular dipole moment of bound e-h states increases their repulsion from one another and from electrons.

Because of the isolation of the X^- states, their binding energies and oscillator strengths remain almost unaffected by the presence of the surrounding 2DEG. This is a somewhat surprising result, and one might rather expect that the interaction of an X^- with Laughlin quasiparticles could affect its recombination. For example, since the X^--QE or X^--QH scattering breaks the $\Delta\mathcal{M}=\Delta\mathcal{K}=0$ selection rule of an isolated X^-, one might expect the (ν-dependent) recombination of the X_{td}^- state embedded in a 2DEG. The exact numerical calculations for finite Ne-h systems with $N\leqslant 9$ show[29] that the X_{td}^- repels QE's and attracts QH's. Although this might suggest discontinuous behavior of PL at Laughlin fillings, we find that the oscillator strength of the X_{td}^- remains negligible compared to the excitonic emission at any ν in the whole range of d in which it is bound.

For example, at $\nu>\frac{1}{3}$, all bound X^- states keep far away from QE's, and the correlations in the vicinity of each X^- are given precisely by the two-component Laughlin wave function[51] $[m_{ee}m_{X^-X^-}m_{eX^-}]$ with Jastrow exponents $m_{ee}=3$ and $m_{eX^-}=2$ (the value of $m_{X^-X^-}$ is irrelevant at small X^- density).[34,35,29] The oscillator strength of an X^- ''locked'' in such (locally) incompressible state [see the $L=2$ ground state in the $9e$-h spectrum in Fig. 1(a)] increases very slowly as a function of d and remains negligible compared to the excitonic emission until X^- unbinds at $d\approx\lambda$. At $\nu<\frac{1}{3}$, the X^- can bind one or two QH's to form a new bound complex X^-QH or X^-QH$_2$ [see, e.g., the X^-QH state at $L=1$ in the $9e$-h spectrum in Fig. 1(b)]. Both these complexes have negligible oscillator strength compared to an exciton. The reason why binding of one or two QH's to an X^- state does not strongly affect its recombination appears to be that the binding of QH's means separation of the X^- from neighboring electrons by an additional (compared to that of a

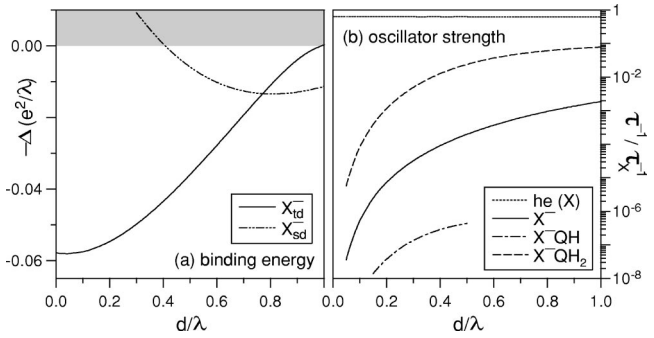

FIG. 2. (a) The binding energy Δ of the isolated dark triplet (X_{td}^-) and dark singlet (X_{sd}^-) charged excitons as a function of layer separation d and calculated on a Haldane sphere at $2S=60$. (b) The PL oscillator strength τ^{-1} of a charged exciton state X_{td}^- binding up to two Laughlin quasiholes QH of the $6e$-X^- incompressible fluid, as a function of d and calculated for an $8e$-h system on a Haldane sphere. The he state contains an exciton and originates from the multiplicative state at $d=0$.

Laughlin state) charge depletion region, without disturbing the X^- state itself. This only weakly modifies the electron wave function in the vicinity of the hole that is probed by PL (PL can be regarded as a one-electron Green function describing the removal of an electron from a state initially occupied by a valence hole). The dependence of the PL oscillator strength of the X^-, X^-QH, and X^-QH_2 states on d (calculated for the $8e$-h system) has been compared to the excitonic emission in Fig. 2(b). In Fig. 2(a) we have also plotted the binding energy of an isolated X_{td}^- and compared it to that of a dark singlet X_{sd}^-.

Summarizing, the PL of a 2DEG in the ''strong-coupling'' regime (small d) occurs from a number of competing radiative bound states: X, X_s^-, and X_{tb}^-, whose optical properties are rather insensitive to the presence (or density) of the surrounding 2DEG. Which of these bound states occur in the 2DEG (and their relative numbers) depends on their binding energies (which in turn depend on B, w, w_h, and d) and on their characteristic formation ($e+X\leftrightarrow X^-$) and recombination times. The X_{td}^- state remains dark, and no other bound states than bright X, X_s^-, and X_{tb}^-, and dark X_{td}^- occur at any ν or d. It is noteworthy that the PL spectrum at small d does not probe the interaction of X or X^- states with the 2DEG (at least at filling factors up to $\nu\sim\frac{1}{3}$). As a result, no information about the original correlations of the 2DEG (before it is perturbed by optically injected valence holes) can be obtained in a PL experiment in the strong-coupling regime. Indeed, the experimental spectra of symmetrically doped QW's are rather insensitive to ν and show no features at the filling factors such as $\nu=\frac{1}{3}$ or $\frac{2}{3}$, at which the Laughlin-Jain incompressible fluid states of the 2DEG occur and the FQH effect is observed in transport experiments.

B. Uncorrelated e-X^- system

The d dependence of the energy spectrum of a $3e$-h system (the simplest system in which to study interaction of X^- states with electrons) shows another interesting feature that might have consequence on PL. At $d=0$, the lowest $3e$-h states describe e-X^- pairs (where X^- is any of the X_s^-, X_{td}^-, or X_{tb}^- bound states),[34,35] and the dependence of energy E on angular momentum L for these states is (up to the appropriate X^- binding energy) equal to the e-X^- interaction pseudopotential $V_{eX^-}(L)$, defined[52,53] as the dependence of the pair the interaction energy on the pair angular momentum. Due to the dipole-dipole e-X repulsion within an X^- at $d>0$, the e-X^- energies anticross the energies of the $2e$-X states (at the same L), in which a $k=0$ exciton is almost decoupled from two interacting electrons (the states that evolve from the MP states at $d=0$). Because the crossings at larger L (i.e., larger pair energy and smaller average e-X^- or e-e separation) occur at smaller d, the stability of the X^- in a e-X^- collision depends critically on both d and L. As we argued in the preceding section, high-energy (i.e., high L) e-X^- collisions do not occur in a Laughlin correlated system. However, if Laughlin correlations were weakened or destroyed by finite QW width ($w>20$ nm), large electron density ($\nu>\frac{1}{3}$), or temperature, such collisions could, for example, result in the breakup of otherwise long-lived X^- states ($e+X^-\rightarrow 2e+X$) and/or their collision-assisted PL from metastable e-X_{td}^- pair states (which would then occur at a higher energy than the excitonic recombination).

VI. PHOTOLUMINESCENCE OF FRACTIONALLY CHARGED EXCITONS AT LARGE LAYER SEPARATION

It was first realized by Chen and Quinn[28] that at a large layer separation d, the PL spectrum of the 2DEG near the Laughlin filling factor $\nu\approx(2p+1)^{-1}$ (i.e., at low density of Laughlin quasiparticles) can be understood in terms of annihilation of a well-defined number n of QE's ($0\leqslant n\leqslant 2p+1$) and/or creation of an appropriate number ($2p+1-n$) of QH's. Independently of the actual average value of ν (average over the entire 2DEG), the recombination probes a finite area of the 2DEG (in the vicinity of the annihilated hole) that has the local filling factor of $\nu=(2p+1)^{-1}$ plus a specific number n of QE's bound to the hole to form a well-defined FCX eigenstate $h\mathrm{QE}_n$. For the $\nu=\frac{1}{3}$ state, four possible recombination events involving QE's and QH's are

$$h+n\mathrm{QE}\rightarrow(3-n)\mathrm{QH}+\gamma, \tag{4}$$

where $n=0$, 1, 2, or 3, and γ denotes the emitted photon. We have verified this conjecture numerically for the $\nu=\frac{1}{3}$ state of up to nine electrons. Indeed, if only the ''first-order'' process (4) is allowed, it describes almost all of the total PL oscillator strength of an initial state $h\mathrm{QE}_n$. However, we find that this process is allowed only for some of the $h\mathrm{QE}_n$ complexes because of the translational symmetry of the 2DEG (in the vicinity of the position of the recombination event). As a result of this symmetry, two angular momentum quantum numbers, $\mathcal{M}$ and $\mathcal{K}$, must be simultaneously conserved in PL.[41,42] To study the selection rules following from the (local) 2D translational invariance, it is more convenient to use spherical geometry, in which they take a simpler form of the conservation of L and L_z.[35]

Let us analyze the four processes (4) in detail. The emission energy ω (we set $\hbar=1$) is measured from the exciton energy E_X (recombination energy of a free $k=0$ exciton in the absence of the 2DEG) at the same d. The PL intensity of the process $i\rightarrow f+\gamma$ is defined as

$$\tau^{-1}=|\langle f|\mathcal{P}|i\rangle|^2, \tag{5}$$

so that $\tau^{-1}\equiv 1$ for the free-exciton recombination. Because of the boson-fermion mapping,[54] identical selection rules are obtained using either statistics to describe Laughlin quasiparticles. In the fermionic picture,[55] the angular momenta of a QE in the initial Ne-h state i and of a QH in the final $(N-1)e$ state f (both at the same monopole strength $2S$) are equal, $l_{\rm QE}=l_{\rm QH}=S-N+2$ (but a QE in state f has different angular momentum of $S-N+3$). The hole angular momentum in the initial state is $l_h=S$.

$h\rightarrow 3{\rm QH}+\gamma$. An infinite planar system without any QE's in the vicinity of the hole is (locally) represented by a finite spherical system at $2S=3(N-1)$. This gives $l_h=S=\frac{3}{2}(N-1)$ and $l_{\rm QH}=\frac{1}{2}(N+1)$. The allowed total angular momenta of three QH's in the final state are obtained by the addition of three angular momenta $l_{\rm QH}$ (of three identical fermionic QH's). The ${\rm QH}_3$ molecule (most tightly packed three-QH droplet) has $l_{{\rm QH}_3}=l_{\rm QH}+(l_{\rm QH}-1)+(l_{\rm QH}-2)=\frac{3}{2}(N-1)$. Since $l_h=l_{{\rm QH}_3}$, the $h\rightarrow 3{\rm QH}+\gamma$ optical process is allowed and creates the ${\rm QH}_3$ molecule. It is expected to have rather small intensity τ^{-1}, because the ''optical hole'' (vacancy) created in the 2DEG by annihilation of a valence hole is given by the single-particle wave function $|m\rangle$ of characteristic radius λ and has small overlap with the much larger ${\rm QH}_3$ molecule. Also, the emission energy ω will be low because of the high energy of QH-QH repulsion in the final state (${\rm QH}_3$ is the eigenstate of pair angular momentum with $\mathcal{R}=1$, i.e., maximum QH-QH repulsion[55]).

$h+{\rm QE}\rightarrow 2{\rm QH}+\gamma$. One QE in the initial state occurs at $2S=3(N-1)-1$ that gives $l_h=\frac{3}{2}N-2$ and $l_{\rm QE}=l_{\rm QH}=\frac{1}{2}N$. The h-QE pair states have angular momentum L_i given by $l_h-l_{\rm QE}\leqslant L_i\leqslant l_h+l_{\rm QE}$. The state at $l_{h{\rm QE}}=l_h-l_{\rm QE}=N-2$ describes the hQE complex with the smallest average h-QE separation. The two QH's in the final state can have pair angular momenta of $L_f=2l_{\rm QH}-\mathcal{R}=N-\mathcal{R}$ where $\mathcal{R}$ is an odd integer, and the ${\rm QH}_2$ molecule has $l_{{\rm QH}_2}=N-1$. Clearly, $l_{h{\rm QE}}\neq L_f$ for any final two-QH state so that the $h{\rm QE}\rightarrow 2{\rm QH}+\gamma$ optical process is forbidden. The hQE can only recombine through a ''second-order'' process, $h{\rm QE}\rightarrow 3{\rm QH}+{\rm QE}+\gamma$, which will have very small intensity. The only state of an h-QE pair that has $L_i=L_f$ and thus can recombine through a ''first-order'' process (4) is the one with the next larger value of angular momentum, $l_{h{\rm QE}^*}=N-1$. This state (denoted by $h{\rm QE}^*$) is[29] the first excited h-QE pair state at d larger than about λ. The $h{\rm QE}^*$ state may occur at a finite temperature as a result of excitation of the long-lived hQE complex. Because ${\rm QH}_2$ is smaller and has (three times) smaller QH-QH repulsion energy than ${\rm QH}_3$, the $h{\rm QE}^*$ is expected to recombine with higher intensity and at higher energy than an uncorrelated hole.

$h+2{\rm QE}\rightarrow{\rm QH}+\gamma$. Two QE's in the initial state occur at $2S=3(N-1)-2$ that gives $l_h=\frac{1}{2}(3N-5)$ and $l_{\rm QE}=l_{\rm QH}=\frac{1}{2}(N-1)$. The ${\rm QE}_2$ molecule has $l_{{\rm QE}_2}=2l_{\rm QE}-1=N-2$. The h-${\rm QE}_2$ pair states have L_i given by $l_h-l_{{\rm QE}_2}\leqslant L_i\leqslant l_h+l_{{\rm QE}_2}$, and the $h{\rm QE}_2$ ground state has $l_{h{\rm QE}_2}=l_h-l_{{\rm QE}_2}=\frac{1}{2}(N-1)$. Since $l_{h{\rm QE}_2}=l_{\rm QH}$, the $h{\rm QE}_2$ state is optically active. Because of the small size and energy of a single QH, the $h{\rm QE}_2$ will recombine at even higher intensity and higher energy than $h{\rm QE}^*$.

$h+3{\rm QE}\rightarrow\gamma$. Three QE's in the initial state occur at $2S=3(N-1)-3$ that gives $l_h=\frac{3}{2}N-3$ and $l_{\rm QE}=\frac{1}{2}N-1$. The ${\rm QE}_3$ molecule has $l_{{\rm QE}_3}=l_{\rm QE}+(l_{\rm QE}-1)+(l_{\rm QE}-2)=\frac{3}{2}N-6$. The h-${\rm QE}_3$ pair states have L_i given by $l_h-l_{{\rm QE}_3}\leqslant L_i\leqslant l_h+l_{{\rm QE}_3}$, i.e., $L_i\geqslant 3$. The smallest value, $l_{h{\rm QE}_3}=3$, describes the $h{\rm QE}_3$ molecule, and all other $h+3{\rm QE}$ states (not only the h-${\rm QE}_3$ pair states) have $L_i>3$. Since $L_f=0$ and $L_i\geqslant 3$, neither the $h{\rm QE}_3$ state nor its excitations can recombine through a ''first-order'' process (4). Instead, the $h{\rm QE}_3$ recombination must occur through a ''second-order'' process, $h+3{\rm QE}_3\rightarrow{\rm QE}+{\rm QH}+\gamma$, which corresponds to recombination of an optically active $h{\rm QE}_2$ in the presence of the nearby third QE. This turns out to be allowed only for $L_i>3$, and hence $h{\rm QE}_3$ is not only unstable,[29] but nonradiative as well.

A. Binding energy and optical properties of $h{\rm QE}_n$ complexes uncoupled from charge excitations of 2DEG

In order to calculate the binding energies Δ, PL energies ω and oscillator strengths τ^{-1} of different $h{\rm QE}_n$ complexes, a finite Ne-h system is diagonalized at the monopole strength $2S=3(N-1)-n$, at which n QE's occur in the $\nu=\frac{1}{3}$ state of N electrons. In this section, the properties of ''isolated'' $h{\rm QE}_n$ complexes are studied. By an isolated $h{\rm QE}_n$ complex we mean one that is uncoupled from additional (other then n QE's) charge excitations of the 2DEG, which is whose wave function involves only the positions of the hole and of n bound QE's. The coupling of the $h{\rm QE}_n$ particles to the underlying 2DEG, as well as its effect on their binding energy and optical properties, will be discussed in Sec. VI B.

To assure that the interaction between the hole and the 2DEG is weak compared to the energy $\varepsilon_{\rm QE}+\varepsilon_{\rm QH}$ ($\approx 0.1e^2/\lambda$ for an infinite system) needed to create additional QE-QH pairs in the 2DEG, the charge of the hole is set to e/ϵ where $\epsilon\gg 1$. This guarantees that the lowest Ne-h states contain exactly n QE's interacting with the hole (even if the large h-QE attraction at a finite d and $\epsilon=1$ induced additional QE-QH pair excitations to screen the hole with additional QE's) and thus that the ground state is the $h{\rm QE}_n$ bound state. If ϵ is sufficiently large, Δ, ω, and τ^{-1} calculated in this way are independent of ϵ and describe the ''ideal'' $h{\rm QE}_n$ wave functions, in which a hole is bound to a ${\rm QE}_n$ molecule [the nQE state with the maximum angular momentum $l_{{\rm QE}_n}=nl_{\rm QE}-\frac{1}{2}n(n-1)$]. The PL intensities τ^{-1} are also independent of d, and the values calculated for the Ne-h systems with $N\leqslant 9$ are listed in Table I. To obtain the

TABLE I. The PL oscillator strength τ_N^{-1} (in the units of the oscillator strength of a free $k=0$ exciton) of fractionally charged excitons $h\mathrm{QE}_n$ calculated in the Ne-h systems ($6\leqslant N\leqslant 9$) on a Haldane sphere.

	MP	h	hQE	hQE*	$h\mathrm{QE}_2$	$h\mathrm{QE}_3$
τ_6^{-1}	0.6154	0.0231		0.0968	0.1144	
τ_7^{-1}	0.6250	0.0187		0.0767	0.0938	
τ_8^{-1}	0.6316	0.0160		0.0649	0.0791	
τ_9^{-1}	0.6364	0.0138		0.0556	0.0680	

dependence of binding energies Δ and PL energies ω on d, the h-QE attraction is multiplied by ϵ. The data obtained for $N=8$ are plotted in Fig. 3.

The MP state in Table I is the lowest energy $L=0$ state at $d=0$ and $2S=3(N-2)$, in which the $k=0$ exciton is decoupled from the $L=0$ Laughlin state of $N-1$ electrons. Its PL oscillator strength equals

$$\tau_{\mathrm{MP}}^{-1}=1-\frac{N-1}{2S+1}\rightarrow 1-\nu \tag{6}$$

for $N\rightarrow\infty$. The he state in Fig. 3(b) is the state that evolves from this MP state when d is increased (it is calculated with full hole charge, $\epsilon=1$), and it has been identified in the $9e$-h spectrum at $d=0.25\lambda$ in Fig. 1(b). Its PL intensity is almost constant at small d (when d increases from 0 to 1, 1.5, and 2λ, then τ_{he}^{-1} decreases by 1%, 6%, and 14%, respectively). Constant τ_{he}^{-1} means almost unchanged wave function, and thus the he state contains a $k=0$ exciton that is only weakly distorted due to interaction with the 2DEG. At $d>2\lambda$, the e-e correlations become dominant and the he state undergoes complete reconstruction (τ_{he}^{-1} drops quickly and Δ_{he} becomes negative). No excitonic recombination is expected in PL spectra at d much larger than 2λ. At $d=0$, the PL energy ω_{he} of the he state equals the energy E_X of a single exciton (because of the hidden symmetry). At $d>0$, ω_{he}

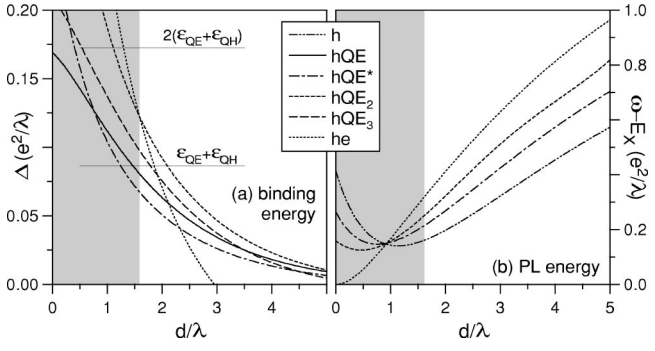

FIG. 3. The binding energy Δ (a) and recombination energy ω (b) of fractionally charged excitons $h\mathrm{QE}_n$ as a function of layer separation d, calculated for the $8e$-h system with a fixed number of Laughlin quasiparticles in the $8e$ electron system ($\epsilon\gg 1$; see text). E_X is the exciton energy and λ is the magnetic length. The he state contains an exciton and originates from the multiplicative states at $d=0$. In the shaded parts of both graphs, the he has the largest binding energy and the $h\mathrm{QE}_n$ complexes do not form.

$>E_X$ because the dipole moment of the $k=0$ exciton (perpendicular to the layers) causes its repulsion by the surrounding electrons.

For FCX's, both intensity and energy behave as predicted in proceeding paragraphs. The occurrence of four possible PL peaks (although not all of them will occur at the same d due to different ranges of stability of different complexes; see Sec. VII) reflects quantization of the total charge $-q$ that can be bound to a hole in the units of the charge of Laughlin quasiparticles: $q/e=1$, $\frac{2}{3}$, $\frac{1}{3}$, and 0 for he, $h\mathrm{QE}_2$, hQE*, and h states, respectively (q does not include the uniform charge density of the underlying Laughlin state).

At $d>\lambda$, all radiative FCX's emit at the energy below ω_{he}. The ordering, $\omega_h<\omega_{h\mathrm{QE}*}<\omega_{h\mathrm{QE}_2}<\omega_{he}$, and almost equal spacing between the PL energies at $d>\lambda$ results from the comparison of the initial- and final-state energies,

$$E_i=N\varepsilon_0+n\varepsilon_{\mathrm{QE}}+\frac{n(n-1)}{2}V_{\mathrm{QE}}-\Delta_{h\mathrm{QE}_n},$$

$$E_f=(N-1)\varepsilon_0+(3-n)\varepsilon_{\mathrm{QH}}+\frac{(3-n)(2-n)}{2}V_{\mathrm{QH}}, \tag{7}$$

where ε_0 is the Laughlin ground-state energy per electron, and $V_{\mathrm{QE}}=V_{\mathrm{QE\text{-}QE}}(1)$ and $V_{\mathrm{QH}}=V_{\mathrm{QH\text{-}QH}}(1)$ are the energies of QE-QE and QH-QH interactions per pair. At $d\gg\lambda$, when the $h\mathrm{QE}_n$ binding energy can be neglected, for the separations between the three FCX peaks we obtain

$$\omega_{h\mathrm{QE}*}-\omega_h=\varepsilon_{\mathrm{QE}}+\varepsilon_{\mathrm{QH}}+2V_{\mathrm{QH}},$$

$$\omega_{h\mathrm{QE}_2}-\omega_{h\mathrm{QE}*}=\varepsilon_{\mathrm{QE}}+\varepsilon_{\mathrm{QH}}+V_{\mathrm{QH}}+V_{\mathrm{QE}}. \tag{8}$$

At smaller d, the separation between peaks decreases because $\Delta_{h\mathrm{QE}_2}>\Delta_{h\mathrm{QE}*}>\Delta_h=0$. The crossing occurs at $d\approx\lambda$. At $d<\lambda$, the ordering of the PL energies in Fig. 3(b) is reversed, but this (shaded) part of the graph has no physical significance (FCX's do not occur). Two other points could be important.

First, the PL oscillator strengths of all radiative FCX's in Table I (h, hQE*, and $h\mathrm{QE}_2$) decrease as a function of N. Hence, the results of our finite-size calculations alone are not conclusive as to whether the recombination of these complexes contributes to the PL spectra of infinite systems. However, the vanishing of τ_h^{-1}, $\tau_{h\mathrm{QE}*}^{-1}$, and $\tau_{h\mathrm{QE}_2}^{-1}$ for $N\rightarrow\infty$ would have to result from an additional, unexpected symmetry recovered in this limit (in analogy to the 2D translational-rotational symmetry that resulted in vanishing of $\tau_{h\mathrm{QE}}^{-1}$ and $\tau_{h\mathrm{QE}_3}^{-1}$). Therefore, it is most likely that h, hQE*, and $h\mathrm{QE}_2$ remain (weakly) optically active in an infinite system, and our data suggests that $\tau_h^{-1}<\tau_{h\mathrm{QE}*}^{-1}<\tau_{h\mathrm{QE}_2}^{-1}$.

Second, Eq. (6) implies $\tau_{\mathrm{MP}}^{-1}\rightarrow 0$ for $\nu\rightarrow 1$, in complete disagreement with experiments that show strong excitonic recombination at $\nu=1$ even at the highest available magnetic fields. This means that the description of the experimentally observed excitonic recombination in terms of the ''hidden symmetry'' of the lowest LL fails completely. Since the LL mixing is more important for the excitonic state he than for

the FCX complexes (due to larger interaction energy compared to the cyclotron energy), one can expect enhancement of the he binding at finite B compared to the FCX binding energies. Although this enhancement depends on a particular system (B, QW width, etc.), we have checked that for parameters of Ref. 35 (symmetric 11.5 nm GaAs QW), inclusion of excited LL's lowers the energy of a free exciton by 0.25, 0.12, 0.035, and 0.015 e^2/λ at $B=5$, 10, 30, and 50 T, respectively. Hence, at high magnetic fields ($B\geqslant 10$ T) it can be assumed that even though our he energy obtained in the lowest LL approximation is not very accurate, the error of this approximation is smaller than the peak splittings in Fig. 3(b) and the ordering of peaks is predicted correctly. However, at low fields ($B\leqslant 5$ T) the excitonic state he will probably remain bound up to much larger d than predicted in Fig. 3(a), and its recombination could occur below that of FCX complexes.

B. Binding energy and optical properties of $h\mathrm{QE}_n$ complexes coupled to charge excitations of 2DEG

In this section we calculate the optical properties of the $h\mathrm{QE}_n$ complexes coupled to the underlying 2DEG, that is of actual complexes that occur in an e-h system at a finite d. To do so, the finite Ne-h spectra similar to those in Fig. 1 are calculated including both e-e and e-h interactions (i.e., with $\epsilon=1$). The $h\mathrm{QE}_n$ complexes are identified in these spectra as low energy states with appropriate angular momentum. The binding energy Δ, PL recombination energy ω, and PL oscillator strength τ^{-1} are calculated for these states and compared with the values obtained for $\epsilon\gg 1$ in Sec. VI A. Small difference between the values obtained for $\epsilon\gg 1$ and $\epsilon=1$, as well as the convergence of the two in the $d\to\infty$ limit, confirms the identification of $h\mathrm{QE}_n$ states in the Ne-h spectra.

Figure 4 shows the data calculated for an $8e$-h system. We have checked that the curves plotted here for $N=8$ are very close to those obtained for $N=7$ or 9, so that all important properties of an extended system can be understood from a rather simple $8e$-h computation. In four frames, for each $h\mathrm{QE}_n$ we plot (a) the excitation gap E^*-E above the $h\mathrm{QE}_n$ ground state; (b) the binding energy Δ; (c) the recombination energy ω; and (d) the recombination intensity (PL oscillator strength) τ^{-1}. The excitation gaps and the recombination energies and intensities are obtained from the spectra at $2S=3(N-1)-n$ in which the $h\mathrm{QE}_n$ complexes occur. The binding energy Δ is defined in such a way that $E_{h\mathrm{QE}_n}=E_{\mathrm{QE}_n}+V_{h-\mathrm{LS}}-\Delta$, where $E_{h\mathrm{QE}_n}$ is the energy of the Ne-h system in state $h\mathrm{QE}_n$ calculated at $2S=3(N-1)-n$, E_{QE_n} is the energy of the Ne system in state QE_n calculated at the same $2S=3(N-1)-n$, and $V_{h\text{-LS}}$ is the self-energy of the hole in Laughlin $\nu=\frac{1}{3}$ ground state at $2S=3(N-1)$. As described in Sec. VI A, $V_{h\text{-LS}}$ is calculated by setting the hole charge to a very small fraction of $+e$ so that it does not perturb the Laughlin ground state. The he curves in Fig. 4 are identical to those in Fig. 3. The PL intensity of the he state (which is the MP state at $d=0$) is too large (see Table I) to fit in Fig. 4(d).

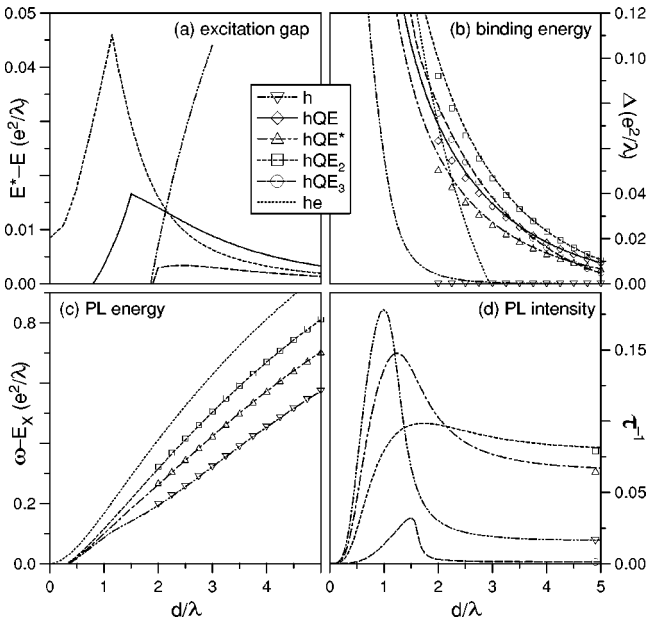

FIG. 4. The excitation gap E^*-E (a), binding energy Δ (b), recombination energy ω (c), and oscillator strength τ^{-1} (d) of fractionally charged excitons $h\mathrm{QE}_n$ as a function of layer separation d, calculated for the $8e$-h system. E_X is the exciton energy and λ is the magnetic length. The he state contains an exciton and originates from the multiplicative states at $d=0$.

The lines in Fig. 4 show data obtained from the spectra similar to those in Fig. 1, which is including all effects of e-h interactions. For comparison, with symbols we have replotted the data from Fig. 3 obtained for $\epsilon\gg 1$ to assure that, at any d, the obtained low-energy eigenstates are given exactly by the $h\mathrm{QE}_n$ wave functions. At $d>\lambda$, very good agreement between binding energies and PL energies calculated for $\epsilon=1$ (lines) and $\epsilon\gg 1$ (symbols) confirms our identification of $h\mathrm{QE}_n$ states in low-energy Ne-h spectra. The PL intensities τ^{-1} calculated for $\epsilon=1$ (lines) converge to those obtained for $\epsilon\gg 1$ and listed in Table I (symbols). The good agreement between the lines and symbols at $d>2\lambda$ shows that the $h\mathrm{QE}_n$ states identified in that Ne-h spectra are indeed described by exact $h\mathrm{QE}_n$ wave functions. At $d<\lambda$, the two calculations give quite different results, which confirms that the description of actual Ne-h eigenstates in terms of the hole interacting with Laughlin quasiparticles of the 2DEG is inappropriate (the correct picture is that of a two-component e-X^- fluid). The formation of $h\mathrm{QE}_n$ complexes at d larger than about 1.5λ can be seen most clearly in the $\tau^{-1}(d)$ curves. For example, while it is impossible to detect the point of transition between the $X^-\mathrm{QH}_2$ and $h\mathrm{QE}_2$ complexes in the dependence of energy spectrum in Fig. 1 on d (because $l_{X^-\mathrm{QH}_2}=l_{h\mathrm{QE}_2}$), it is clearly visible at $d\approx 1.5\lambda$ in Fig. 4(d).

The analysis of the characteristics of $h\mathrm{QE}_n$ complexes plotted in Figs. 3 and 4 leads to the conclusion that the bound complex most important for understanding PL in the weak-coupling regime ($d>1.5\lambda$) is $h\mathrm{QE}_2$, which has the largest binding energy Δ, and significant excitation energy E^*-E and PL oscillator strength τ^{-1}. The $h\mathrm{QE}$ is also a strongly bound complex with large excitation energy, but it is nonra-

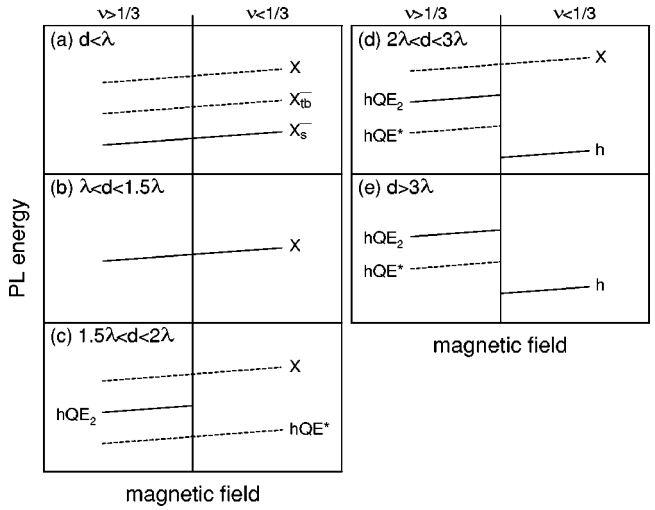

FIG. 5. The schematic PL spectra (PL energy vs magnetic field) near the filling factor $\nu=\frac{1}{3}$ at different layer separations d. Solid and dashed lines mark recombination from ground and excited states, respectively. λ is the magnetic length.

diative (at least, in the absence of scattering or disorder). Although hQE is dark, its first excited state, hQE*, is both bound and radiative and can contribute to the PL spectrum. The charge neutral ''anyon exciton'' $h\mathrm{QE}_3$ suggested by Rashba *et al.*[27] is neither bound nor radiative. Finally, the radiative excitonic state ($k=0$ charge neutral e-h pair weakly coupled to the 2DEG) breaks apart at $d>2\lambda$.

VII. STABILITY AND EMISSION OF DIFFERENT BOUND STATES: PL SPECTRA AT DIFFERENT LAYER SEPARATIONS

The information presented in Figs. 3 and 4 and in Table I allows understanding of anomalies observed in the PL spectra of the 2DEG near $\nu=\frac{1}{3}$. The crucial observations are (i) the most strongly bound complexes at small layer separation d are the $k=0$ state of a charge-neutral exciton X and different states of charged excitons X^-; (ii) at larger d, the most stable complexes are the bright $h\mathrm{QE}_2$ and dark hQE (whose weakly excited state hQE* is bright); (iii) no charge-neutral ''dressed exciton'' states at $k\neq 0$ occur; and (iv) the charge-neutral ''anyon exciton'' $h\mathrm{QE}_3$ is neither stable nor radiative. Depending on the layer separation d and on whether ν is larger or smaller than $\frac{1}{3}$, the following behavior is expected (see the schematic PL spectra in Fig. 5; it should be understood that the PL spectrum changes continuously as a function of d but discontinuously as a function of ν).

$d<\lambda$. The holes bind one or two ''whole'' electrons to form $k=0$ neutral excitons or various charged exciton states (the relative numbers of X_s^-, X_{td}^-, and X_{tb}^- depend on B, temperature, etc.). No ''dressed exciton'' states with $k\neq 0$ (in-plane dipole moment) occur. The $k=0$ excitons weakly couple to the 2DEG, and the X^-'s are effectively isolated from the 2DEG because of Laughlin e-X^- correlations. As a result, neither the recombination of a $k=0$ exciton and radiative X^- states (X_s^- and X_{tb}^-) nor the lack of recombination of the dark X_{td}^- state are significantly affected by the 2DEG. Only the X^-'s will occur in the absolute ground state of the system. However, because the exciton has shorter optical lifetime than all the X^- states, the PL spectrum at finite temperatures contains peaks corresponding to both exciton (in our notation: he) and X^- recombination. At $d>0$, the he recombination energy is larger than the bare ($\nu=0$) exciton energy E_X at the same d due to the e-X repulsion.

$\lambda\leqslant d<1.5\lambda$. The X^-'s unbind but the neutral excitons still exist. The FCX complexes (hQE and $h\mathrm{QE}_2$) also occur, as the QE-QH pairs are spontaneously created in the 2DEG to screen the charge of each hole. However, the exciton has both the largest binding energy and the largest PL oscillator strength, and its recombination dominates the PL spectrum. A similar electric-field-induced ionization of X^-'s in a QW has been demonstrated at $B=0$ by Shields *et al.*[56]

$1.5\lambda\leqslant d<2\lambda$. The excitons still exist but they no longer have maximum binding energy. To screen the charge of each hole, one QE-QH pair is spontaneously created in the 2DEG to form the FCX complex hQE ($h\rightarrow h\mathrm{QE}+\mathrm{QH}$). Since $h\mathrm{QE}_2$ has larger binding energy than hQE, it can also be formed in the presence of excess QE's ($h\mathrm{QE}+\mathrm{QE}\rightarrow h\mathrm{QE}_2$) but it will be destroyed in the presence of excess QH's ($h\mathrm{QE}_2+\mathrm{QH}\rightarrow h\mathrm{QE}$). Therefore, a discontinuity is expected when ν crosses $\frac{1}{3}$: At $\nu>\frac{1}{3}$, the dark hQE and the bright $h\mathrm{QE}_2$ coexist and the hQE recombination can occur either through binding of the second QE to form a bright $h\mathrm{QE}_2$ (note that $\nu_{\mathrm{QE}}=1$ occurs at $\nu=\frac{2}{5}$ and thus, except for ν almost equal to $\frac{1}{3}$, the QE density is larger than the hole density) or through excitation to a bright hQE*. At $\nu<\frac{1}{3}$, the hQE is the only stable complex and its dominant recombination channel is through excitation to the bright hQE* state that emits at similar rate but lower energy than $h\mathrm{QE}_2$ (by about $\varepsilon_{\mathrm{QE}}+\varepsilon_{\mathrm{QH}}$). The strongly radiative $k=0$ excitons (he) are also visible at finite temperatures. Clearly, different temperature dependence of the emission from the ground state $h\mathrm{QE}_2$ and from the excited states hQE* and he is expected.

$2\lambda\leqslant d<3\lambda$. The excitons still exist but they have very small binding energy. No QE-QH pairs are spontaneously created and the holes can only bind existing QE's, which leads to discontinuity when ν crosses $\frac{1}{3}$: At $\nu>\frac{1}{3}$, the relative numbers of hQE, hQE*, and $h\mathrm{QE}_2$ depend on the hole and QE densities and temperature. However, because the QE density can be assumed larger than the hole density and the hQE is long-lived, both hQE* and $h\mathrm{QE}_2$ are expected to show in the PL spectrum, emitting at energies different by about $\varepsilon_{\mathrm{QE}}+\varepsilon_{\mathrm{QH}}$. At $\nu<\frac{1}{3}$, there are no QE's to bind, and the holes repel the existing QH's. In the ground state, there is no response of the 2DEG to the hole, whose recombination occurs at the local filling factor $\nu=\frac{1}{3}$ (and probes the spectral function of an electron annihilated in an undisturbed Laughlin $\nu=\frac{1}{3}$ state). Although the optical lifetime of an unbound hole is fairly long, no bound radiative FCX's are expected at low temperatures since the recombination of $h\mathrm{QE}_2$ or hQE* must occur through the formation of an unstable hQE ($h\mathrm{QE}+\mathrm{QH}\rightarrow h$) followed by either binding of a second QE to form the $h\mathrm{QE}_2$ or an excitation to form the hQE*. Although weakly bound, the neutral exciton (he) has much

larger oscillator strength than an uncorrelated hole, and it might also be observed in PL at a finite temperature. The exciton binding strongly depends on the LL mixing, so it is more likely to exist at lower B (in the lower density samples).

$d \geq 3\lambda$. No excitons (he) occur, and the recombination can only occur from the $h\mathrm{QE}_2$, $h\mathrm{QE}^*$, or h states, with a discontinuity at $\nu=\frac{1}{3}$.

VIII. CONCLUSION

We have studied PL from a 2DEG in the fractional quantum Hall regime as a function of the separation d between the electron and valence hole layers. Possible bound states in which the hole can occur have been identified and characterized in terms of such single-particle quantities as the angular momentum, binding energy, recombination energy, and oscillator strength. The strict optical selection rules for these bound states have been formulated, following from the (local) 2D translational symmetry of each state. Only some of the bound states turn out radiative, and their relative oscillator strengths are predicted from a rather simple analysis. The discussion is illustrated with the results of exact numerical calculations in Haldane's spherical geometry for a hole interacting with up to nine electrons at filling factors $\nu \sim \frac{1}{3}$.

Different response of the 2DEG to the optically injected hole in the strong- and weak-coupling regime results in a complete reconstruction of the PL spectrum at d of the order of the magnetic length λ. At $d<\lambda$, the hole binds one or two electrons to form a neutral exciton state X or various charged exciton states X^-. The PL spectrum in this regime depends on the lifetimes and binding energies of the X and X^- states, rather than on the original correlations of the 2DEG. No anomaly occurs in PL at the Laughlin filling factor $\nu=\frac{1}{3}$, at which the FQH effect is observed in transport experiments.

At d larger than about 2λ, the Coulomb potential of the distant hole becomes too weak and its range becomes too large to bind individual electrons and form the X or X^- states. Instead, the hole interacts with charge excitations of the 2DEG, namely, repels QH's and attracts QE's of the Laughlin incompressible $\nu=\frac{1}{3}$ fluid. The resulting states in which the hole can occur are the uncorrelated state h (in which the free hole moves in the rigid electron Laughlin fluid at a local filling factor $\nu=\frac{1}{3}$) and the fractionally charged excitons $h\mathrm{QE}$ and $h\mathrm{QE}_2$ (in which the hole binds one or two QE's). Different states have very different optical properties (recombination lifetimes and energies) and which of them occur depends critically on whether QE's are present in the 2DEG. Therefore, discontinuities occur in the PL spectrum at $\nu=\frac{1}{3}$.

Our results invalidate two suggestive concepts proposed to understand the numerical Ne-h spectra and the observed PL of a 2DEG. First, the ''dressed exciton'' states[25,26] with finite momentum ($k\neq 0$) do not occur in the low-energy spectra of e-h systems at small d. Second, the charge neutral ''anyon exciton'' states[27] are neither stable nor radiative at any value of d.

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