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Nanoengineering

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Physics of Low-Dimensional Structures

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Contents

I. Composite fermion theory of quantum Hall effect 4

1. Introduction
2. Integral and fractional quantum Hall effects
3. Numerical study of small systems
4. Chern-Simons transformation and statistics in 2D systems
5. Jain's composite fermion picture
6. Energy scales and pseudopotentials
 - 6.1 Two-fermion problem
 - 6.2 Two-body interaction pseudopotential
 - 6.3 Fractional grandparentage
 - 6.4 Definition of the short range pseudopotential
 - 6.5 Hidden symmetry of the short range repulsion
 - 6.6 Comparison with atomic shells: Hund's rule
 - 6.7 Higher Landau levels
7. Fermi Liquid model of composite fermions
8. Composite fermion hierarchy

II. Electron-hole systems in high magnetic fields 44

1. Excitonic complexes
2. Hidden symmetry
3. Interaction of charged excitons
4. Generalized composite fermion picture for charged excitons
5. Spatially separated electron-hole system
6. Charged excitons in a finite magnetic field

Summary 57

Literature 58

I. Composite fermion theory of quantum Hall effect

The composite Fermion (composite fermion) picture [1] offers a simple intuitive way of understanding many of the surprising properties of a strongly interacting two-dimensional electron fluid in a large magnetic field. In this chapter, the simple way in which the mean field composite fermion picture describes the low lying bands of angular momentum multiplets for any value of the applied magnetic field is illustrated and compared with the results of exact numerical diagonalization of small systems. The reason for the success of the composite fermion approach is discussed in some detail, and a composite fermion hierarchy picture of the incompressible quantum fluid states is introduced. The composite fermion picture is used to understand the energy spectrum and photoluminescence of systems containing both electrons and valence band holes.

1. Introduction

The study of the electronic properties of quasi-two-dimensional (2D) systems has resulted in a number of remarkable discoveries in the past two decades. Among the most interesting of these are the integral [2] and fractional [3] quantum Hall effects. In both of these effects, incompressible states of a 2D electron liquid are found at particular values of the electron density for a given value of the magnetic field applied normal to the 2D layer.

The integral quantum Hall effect (IQHE) is rather simple to understand. The incompressibility results from a cyclotron energy gap,

$$\hbar\omega_c,$$

in the single particle spectrum. When all states below the gap are filled and all states above it are empty, it takes a finite (cyclotron) energy to produce an infinitesimal compression.

Excited states consist of electron-hole pair excitations and require a finite excitation energy.

Both localized [4] and extended single particle states are necessary to understand the experimentally observed behavior of the magneto-conductivity [5].

The fractional quantum Hall effect (FQHE) is more difficult to understand and more interesting in terms of new basic physics. The energy gap that gives rise to the Laughlin [6] incompressible fluid state is completely the result of the interaction between the electrons.

The elementary excitations are fractionally charged Laughlin quasiparticles, which satisfy fractional statistics [7]. The standard techniques of many body perturbation theory are incapable of treating fractional quantum Hall systems because of the complete degeneracy of the single particle levels in the absence of the interactions. Laughlin [6] was able to determine the form of the ground state wave function and of the elementary excitations on the basis of physical insight into the nature of the many body correlations.

Striking confirmation of Laughlin's picture was obtained by exact diagonalization of the interaction Hamiltonian within the subspace of the lowest Landau level of small systems [8].

Jain [9], Lopez and Fradkin [10], and Halperin, Lee and Read [11] have extended Laughlin's approach and developed a composite Fermion (composite fermion) description of the 2D electron gas in a strong magnetic field. This composite fermion description has offered a simple picture for the interpretation of many experimental results. However, the underlying reason for the validity of many of the approximations used with the composite fermion approach was not immediately understood [12].

The object of this chapter is to present a simple and understandable summary of the composite fermion picture as applied to fractional quantum Hall systems. Exact numerical calculations for up to eleven electrons on a spherical surface will be compared with the predictions of the mean field composite fermion picture. The composite fermion hierarchy [13] will be introduced, and its predictions compared with numerical results. It will be shown that sometimes the mean field composite fermion hierarchy correctly predicts Laughlin-like incompressible ground states, and that sometimes it fails.

The composite fermion hierarchy depends on the validity of the mean field approximation. This seems to work well in predicting not only the Laughlin-Jain families of incompressible ground states at particular values of the applied magnetic field, but also in predicting the lowest lying band of states at any value of the magnetic field. The question of when the mean field composite fermion picture works and why [12] will be discussed in some detail. As first suggested by Haldane [8], the behavior of the pseudopotential

$$V(L)$$

describing the energy of interaction of a pair of electrons as a function of their total angular momentum L is of critical importance.

In the following sections some examples of other strongly interacting 2D Fermion systems will be presented, and some problems not yet completely understood will be introduced. The single particle states for electrons confined to a plane in the presence of an applied magnetic field [14] are explained. The integral and fractional quantum Hall effects are discussed briefly. Haldane's idea [8] that the condensation of Laughlin quasiparticles leads to a hierarchy containing all odd denominator fractions is discussed. The numerical calculations for a finite number of electrons confined to a spherical surface in the presence of a radial magnetic field are discussed. Results for a ten electron system at different values of the magnetic field are presented. The ideas of fractional statistics and the Chern-Simons transformation are introduced. Jain's composite fermion approach [1] is outlined. The sequence of Jain condensed states (given by filling factor

$$\nu = \frac{n}{2pn + 1}$$

(where n is any integer and p is a positive integer) is shown to result from the mean field approximation. The application of the composite fermion picture to electrons on a spherical surface is shown to predict the lowest band of angular momentum multiplets in a very simple way that involves only the elementary problem of addition of angular momenta [15]. The two energy scales, the Landau level separation

$$\hbar\omega_c$$

and the Coulomb energy

$$\frac{e^2}{\lambda}$$

(where λ is the magnetic length) are discussed. It is emphasized that the Coulomb interactions and Chern-Simons gauge interactions between fluctuations (beyond the mean field level) cannot possibly cancel for arbitrary values the applied magnetic field. The reason for the success of the composite fermion picture is discussed in terms of the behavior of the pseudopotential $V(L)$ and a kind of ‘‘Hund’s rule’’ for monopole harmonics [12]. A phenomenological Fermi liquid picture is introduced to describe low lying excited states containing three or more Laughlin quasiparticles [16]. The composite fermion hierarchy picture [13] is introduced. Comparison with exact numerical results indicates that the behavior of the quasiparticle pseudopotential is of critical importance in determining the validity of this picture at a particular level of the hierarchy. The systems containing electrons and valence holes are also investigated [17]. Photoluminescence and the role of excitons and negatively charged exciton complexes is also discussed.

2. Integral and fractional quantum Hall effects

The Hamiltonian for an electron confined to the x-y plane in the presence of a perpendicular magnetic field B is

$$H_0 = \frac{1}{2\mu} \left(p + \frac{e}{c} A \right)^2$$

Here μ is the effective mass, $p=(p_x, p_y, 0)$ is the momentum operator and $A(x, y)$ is the vector potential (whose curl gives B). For the ‘‘symmetric gauge,’’

$$A = \frac{B}{2} (y, -x, 0),$$

the single particle eigenfunctions are of the form

$$\psi_{nm}(r, \theta) = e^{-im\phi} u_{nm}(r)$$

The angular momentum of the state ψ_{nm} is $-m$ and its eigenenergy is given by

$$E_{nm} = \frac{1}{2} \hbar \omega_c (2n + 1 + |m| - m)$$

In this equation,

$$\hbar \omega_c = \frac{eB}{\mu c}$$

is the cyclotron frequency, $n=0, 1, 2, \dots$, and $m=0, \pm 1, \pm 2, \dots$

The lowest energy states (lowest Landau level) have $n=0$ and $m=0, 1, 2, \dots$, and energy

$$E_{0m} = \frac{1}{2} \hbar \omega_c$$

It is convenient to introduce a complex coordinate

$$z = r e^{-i\theta} = x - iy$$

and to write the lowest Landau level wave functions as

$$\psi_{0m}(z) = N_m z^m e^{-|z|^2/4}$$

where N_m is a normalization constant.

In this expression we have used the magnetic length

$$\lambda = \sqrt{\frac{\hbar c}{eB}}$$

as the unit of length.

The function

$$|\psi_{0m}(z)|^2$$

has its maximum value at a radius r_m which is proportional to $m^{1/2}$.

All single particle states belonging to a given Landau level are degenerate, and separated in energy from neighboring levels by the cyclotron energy.

If the system has a “finite radial range,” then the m -values are restricted to being less than some maximum value

$$m = 0, 1, 2, \dots, N_\phi - 1.$$

The value of N_ϕ (the Landau level degeneracy) is equal to the total flux through the sample, BC (where C is the area), divided by the quantum of flux

$$\varphi_0 = \frac{hc}{e}.$$

The filling factor ν is defined as the ratio of the number of electrons, N , to N_ϕ . When ν has an integral value, an infinitesimal decrease in the area C requires promotion of an electron across the cyclotron gap to the first unoccupied Landau level, making the system incompressible.

This incompressibility together with the existence of both localized and extended states in the system is responsible for the observed behavior of the magneto-conductivity of quantum Hall systems at integral filling factors [5].

In order to construct a many electron wave function

$$\Psi(z_1, z_2, \dots, z_N)$$

corresponding to a completely filled lowest Landau level, the product function which places one electron in each of the $N=N_\phi$ orbitals

$$\psi_{0m} \quad (m = 0, 1, \dots, N - 1)$$

must be antisymmetrized. This can be done with the aid of a Slater determinant

$$\Psi = \begin{vmatrix} 1 & 1 & 1 & \dots & 1 \\ z_1 & z_2 & z_3 & & z_N \\ z_1^2 & z_2^2 & z_3^2 & & z_N^2 \\ \vdots & & & \ddots & \\ z_1^N & z_2^N & z_3^N & & z_N^N \end{vmatrix} \exp\left(-\frac{1}{4} \sum_{k=1}^N |z_k|^2\right)$$

This is the well-known Vandemonde determinant. It is not difficult to show that it is equal to

$$\Psi_1 = \prod_{i < j} (z_i - z_j) \exp(\dots)$$

Of course, since each of the N_ϕ orbitals is occupied by one electron, the filling factor is $\nu=1$.

Laughlin noticed that if the factor $(z_i - z_j)$ arising from the Vandemonde determinant was replaced by $(z_i - z_j)^{2p+1}$, where p was an integer, the wave function

$$\Psi_{2p+1} = \prod_{i < j} (z_i - z_j)^{2p+1} \exp(\dots)$$

would be antisymmetric, keep the electrons further apart (and therefore reduce the Coulomb repulsion), and correspond to a filling factor

$$\nu = \frac{1}{2p+1}.$$

This results because the highest power of z_i in the polynomial factor in Ψ_{2p+1} is $(2p+1)(N-1)$ and it must be equal to the highest orbital index ($m=N_\phi-1$), giving

$$N_\phi - 1 = (2p+1)(N-1)$$

and

$$\nu = \frac{N}{N_\phi} = \frac{1}{2p+1}$$

in the limit of large systems.

The additional factor

$$\prod_{i < j} (z_i - z_j)^{2p}$$

multiplying Ψ_1 is the Jastrow factor which accounts for correlations between electrons.

It is observed experimentally that states with filling factors $\nu=2/5, 3/5, 3/7$, etc. exhibit fractional quantum Hall behavior in addition to the Laughlin $\nu=(2p+1)^{-1}$ states.

Haldane [8] suggested that a hierarchy of condensed states arose from the condensation of quasiparticles (QP's) of "parent" fractional quantum Hall states.

In his picture, Laughlin condensed states of the electron system occurred when

$$N_\phi = (2p+1)N_e$$

where the exponent $2p+1$ was an odd integer and the symbol N_e denoted the number of electrons. Condensed QP states occurred when $N_e=2qN_{QP}$, because the number of places available for inserting a QP in a Laughlin state was N_e . Haldane required the exponent $2q$ to be even because in his definition of a QP it was a boson.

This scheme gives rise to a hierarchy of condensed states which contains all odd denominator fractions. Haldane cautioned that the validity of the hierarchy scheme at a particular level depended upon the QP interactions which were totally unknown.

3. Numerical study of small systems

Haldane [8] introduced the idea of putting a small number of electrons on a spherical surface of radius R at the center of which is a magnetic monopole of strength $2S\phi_0$.

The single particle Hamiltonian can be expressed as [19]

$$H_0 = \frac{\hbar^2}{2\mu R^2} (\mathbf{L} - S\hat{R})^2$$

where \mathbf{L} is the angular momentum vector operator (in the units of Planck's constant), R is the unit vector in the radial direction, and μ is the mass.

The components of \mathbf{L} satisfy the usual commutation rules

$$[L_\alpha, L_\beta] = i\epsilon_{\alpha\beta\gamma} L_\gamma$$

The eigenstates of H_0 can be denoted by

$$|l, m\rangle;$$

they are eigenfunctions of L^2 and L_z with eigenvalues $l(l+1)$ and m , respectively.

The lowest energy eigenvalue (shell) occurs for $l=S$ and has energy

$$E_0 = \frac{1}{2} \hbar\omega_c$$

The n^{th} excited shell has $l=S+n$, and

$$E_n = \frac{\hbar\omega_c}{2S} [l(l+1) - S^2] = \hbar\omega_c \left[n + \frac{1}{2} + \frac{n(n+1)}{2S} \right]$$

where the cyclotron energy is equal to

$$\hbar\omega_c = S \frac{\hbar^2}{\mu R^2}$$

and the magnetic length is

$$\lambda = \frac{R}{\sqrt{S}}$$

If we concentrate on a partially filled lowest Landau level we have only

$$N_\phi = 2S + 1$$

degenerate single particle states (since the electron angular momentum l must be equal to S and its z-component m can take on values between $-l$ and l).

The Hilbert space \mathcal{H}_{MB} of N electrons in these N_ϕ single particle states contains

$$N_{MB} = \binom{N_\phi}{N} = \frac{N_\phi!}{N!(N_\phi - N)!}$$

antisymmetric many body states.

The single particle configurations

$$|m_1, m_2, \dots, m_N\rangle = c_{m_1}^+ c_{m_2}^+ \dots c_{m_N}^+ |vac\rangle$$

can be chosen as a basis of \mathcal{H}_{MB} . Here c_m^+ creates an electron in the single particle state

$$|m\rangle \equiv |l = S, m\rangle$$

and

$$|vac\rangle$$

is the vacuum state.

The space \mathcal{H}_{MB} can also be spanned by the angular momentum eigenfunctions,

$$|L, M, \alpha\rangle$$

where L is the total angular momentum, M is its z-component, and α is a label which distinguishes different multiplets with the same L .

If

$$\hbar\omega_c \gg \frac{e^2}{\lambda}$$

the diagonalization of the interaction Hamiltonian

$$H_I = \sum_{i < j} \frac{e^2}{r_{ij}}$$

in the Hilbert space \mathcal{H}_{MB} of the lowest Landau level gives an excellent approximation to exact eigenstates of an interacting N electron system.

The single particle configuration basis is particularly convenient since the many body interaction matrix elements in this basis,

$$\langle m'_1, m'_2, \dots, m'_N | H_I | m_1, m_2, \dots, m_N \rangle$$

are expressed through the two body ones,

$$\langle m'_1, m'_2 | H_I | m_1, m_2 \rangle$$

in a very simple way.

On the other hand, using the angular momentum eigenstates

$$|L, M, \alpha\rangle$$

allows the explicit decomposition of the total Hilbert space \mathcal{H}_{MB} into total angular momentum eigensubspaces.

Because the interaction Hamiltonian is a scalar, the Wigner-Eckart theorem tells us that

$$\langle L', M', \alpha' | H_I | L, M, \alpha \rangle = \delta_{LL'} \delta_{MM'} V_{\alpha\alpha'}(L)$$

where the reduced matrix element

$$V_{\alpha\alpha'}(L) = \langle L, \alpha' | H_I | L, \alpha \rangle$$

is independent of M.

The eigenfunctions of L are simpler to find than those of H_I , because efficient numerical techniques exist for obtaining eigenfunctions of operators with known eigenvalues.

Finding the eigenfunctions of L and then using the Wigner-Eckart theorem considerably reduces dimensions of the matrices that must be diagonalized to obtain eigenvalues of H_I .

Some matrix dimensions are listed in Tab. 1, where the degeneracy of the lowest Landau level and the dimensions of the total many body Hilbert space, \mathcal{H}_{MB} , and of the largest M subspace, $N_{\text{MB}}(M=0)$, are given for the Laughlin $\nu=1/3$ state of six to eleven electron systems. Recall that the N electron Laughlin $\nu=1/(2p+1)$ state occurs at $N_\phi=(2p+1)(N-1)$.

Tab. 1 The Landau level degeneracy $N_\phi=2S+1$ and the dimensions of the total N electron Hilbert space, N_{MB} , and of the largest M-subspace, $N_{\text{MB}}(M=0)$, at the filling factor $\nu=1/3$.

N	N_ϕ	N_{MB}	$N_{\text{MB}}(M=0)$
6	16	8008	338
7	19	50388	1656
8	22	319770	8512
9	25	2042975	45207
10	28	13123110	246448
11	31	84672315	1371535

For example, in the eleven electron system at $\nu=1/3$, the $L=0$ block that must be diagonalized to obtain the Laughlin ground state is only 1160×1160 , small compared to the total dimension of 1 371 535 for the entire $M=0$ subspace.

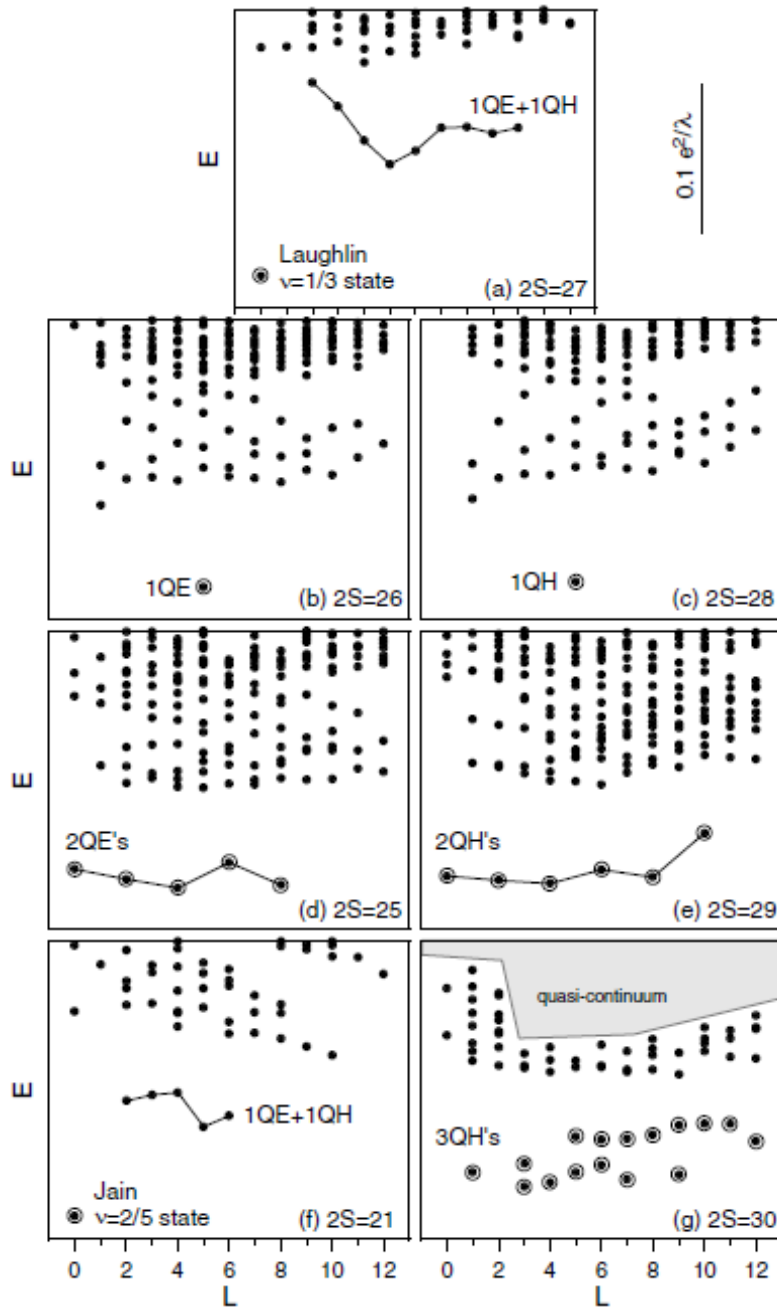


Fig. 1 The energy spectra of ten electrons in the lowest Landau level calculated on a Haldane sphere with $2S$ between 21 and 30. The open circles and solid lines mark lowest energy bands with the fewest composite Fermion quasiparticles.

Typical results for the energy spectrum are shown in Fig. 1 for $N=10$ and for a few different values of $2S$ between 21 and 30. The low energy bands marked with open circles and solid lines will be discussed in detail in the following sections. Frames (a) and (f) show two $L=0$ incompressible ground states: Laughlin state at $\nu=1/3$ and Jain state at $\nu=2/5$, respectively.

In other frames, a number of QP's form the lowest energy bands.

4. Chern-Simons transformation and statistics in 2D systems

Before discussing the Chern-Simons gauge transformation and its relation to particle statistics, it is useful to look at a system of two particles each of charge $-e$ and mass μ , confined to a plane, in the presence of a perpendicular magnetic field

$$\mathbf{B} = (0, 0, B) = \nabla \times \mathbf{A}(\mathbf{r})$$

Because \mathbf{A} is linear in the coordinate $\mathbf{r}=(x,y)$ – for example, in the symmetric gauge

$$\mathbf{A} = \frac{B}{2}(y, -x, 0),$$

the Hamiltonian separates into the center of mass (CM) and relative (REL) coordinate pieces, with

$$\mathbf{R} = \frac{1}{2}(\mathbf{r}_1 + \mathbf{r}_2)$$

and

$$\mathbf{r} = \mathbf{r}_1 - \mathbf{r}_2$$

being the CM and REL coordinates, respectively.

The energy spectra of CM and REL hamiltonians, H_{CM} and H_{REL} are identical to that of a single particle of mass μ and charge $-e$.

We have already seen that for the lowest Landau level

$$\psi_{0m}(r, \varphi) = N_m r^m e^{-im\varphi} e^{-|z|^2/4}.$$

For the relative motion

$$\varphi = \varphi_1 - \varphi_2$$

and an interchange of the pair,

$$\hat{P}\psi(\mathbf{r}_1, \mathbf{r}_2) = \psi(\mathbf{r}_2, \mathbf{r}_1)$$

is accomplished by replacing φ by $\varphi+\pi$.

In 3D systems, where two consecutive interchanges must result in the original wave function, this implies that

$$e^{im\pi}$$

must be equal to either $+1$ (m even; Bosons) or -1 (m odd; Fermions).

It is well-known [20,21] that for 2D systems m need not be an integer.

Interchange of a pair of identical particles can give

$$\hat{P}\psi(\mathbf{r}_1, \mathbf{r}_2) = e^{i\theta\pi}\psi(\mathbf{r}_2, \mathbf{r}_1)$$

where the statistical parameter θ can assume non-integral values leading to anyon statistics.

A Chern-Simons (CS) transformation is a singular gauge transformation [11] in which an electron creation operator

$$\psi_e^+(\mathbf{r})$$

is replaced by a composite particle operator given by

$$\psi^+(\mathbf{r}) = \psi_e^+(\mathbf{r}) \exp\left[i\alpha \int d^2\mathbf{r}' \arg(\mathbf{r} - \mathbf{r}') \psi_e^+(\mathbf{r}') \psi_e(\mathbf{r}')\right].$$

Here

$$\arg(\mathbf{r} - \mathbf{r}')$$

is the angle the vector $\mathbf{r} - \mathbf{r}'$ makes with the x-axis and α is an arbitrary parameter.

The kinetic energy operator can be written in terms of the transformed operator as

$$K = \frac{1}{2\mu} \int d^2\mathbf{r} \psi^+(\mathbf{r}) \left[-i\hbar\nabla + \frac{e}{c} \mathbf{A}(\mathbf{r}) + \frac{e}{c} \mathbf{a}(\mathbf{r}) \right] \psi(\mathbf{r})$$

Here

$$\mathbf{a}_r(\mathbf{r}) = \frac{\alpha\phi_0}{2\pi} \frac{\hat{z} \times (\mathbf{r} - \mathbf{r}')}{|\mathbf{r} - \mathbf{r}'|^2}$$

and

$$\mathbf{a}(\mathbf{r}) = \alpha\phi_0 \int d^2\mathbf{r}' \mathbf{a}_r(\mathbf{r}) \psi^+(\mathbf{r}') \psi(\mathbf{r}')$$

where z is a unit vector perpendicular to the 2D layer.

The Chern-Simons transformation can be thought of as an attachment to each particle of flux tube carrying a fictitious flux $\alpha\phi_0$ (where $\phi_0 = hc/e$ is the quantum of flux) and a fictitious charge

- e which couples in the standard way to the vector potential produced by the flux tubes on every other particle.

The

$$\mathbf{a}_r(\mathbf{r})$$

is interpreted as the vector potential at position \mathbf{r} due to a magnetic flux of strength $\alpha\phi_0$ localized at \mathbf{r}' , and

$$\mathbf{a}(\mathbf{r})$$

is the total vector potential at position \mathbf{r} due to all Chern-Simons fluxes.

The Chern-Simons magnetic field associated with the particle at \mathbf{r}' is

$$\mathbf{b}(\mathbf{r}) = \nabla \times \mathbf{a}_r(\mathbf{r}) = \alpha\phi_0 \delta(\mathbf{r} - \mathbf{r}') \hat{z}$$

Because two charged particles cannot occupy the same position, one particle never senses the magnetic field of other particles, but it does sense the vector potential resulting from their Chern-Simons fluxes.

The classical equations of motion are unchanged by the presence of the Chern-Simons flux, but the quantum statistics of the particles are changed unless α is an even integer.

For the two-particle system, the vector potential associated with the Chern-Simons flux

$$\mathbf{a}_{\mathbf{r}_2}(\mathbf{r}_1)$$

depends only on the relative coordinate

$$\mathbf{r} = \mathbf{r}_1 - \mathbf{r}_2.$$

When $\mathbf{a}(\mathbf{r})$ is added to $\mathbf{A}(\mathbf{r})$, the vector potential of the applied magnetic field, the Schrödinger equation has a solution

$$\tilde{\psi}_m = e^{-i\alpha\phi} \psi_m$$

where ψ_m is the solution with $\alpha=0$ (i.e. in the absence of Chern-Simons flux).

If α is an odd integer, Boson and Fermion statistics are interchanged; if α is even, no change in statistics occurs and electrons are transformed into composite Fermions with an identical energy spectrum.

The Hamiltonian for the composite particle system (charged particles with attached flux tubes) is much more complicated than the original one with $\alpha=0$.

What is gained by making the Chern-Simons transformation?

The answer is that one can use the “mean field” approximation in which

$$\mathbf{A}(\mathbf{r}) + \mathbf{a}(\mathbf{r}),$$

the vector potential of the external plus Chern-Simons magnetic fields, is replaced by

$$\mathbf{A}(\mathbf{r}) + \langle \mathbf{a}(\mathbf{r}) \rangle,$$

where $\langle \dots \rangle$ is the mean field value obtained by simply replacing

$$\rho(\mathbf{r}') = \psi^\dagger(\mathbf{r}') \psi(\mathbf{r}')$$

by its average value ρ_0 .

A mean field energy spectrum can be constructed in which the massive degeneracy of the original partially filled electron Landau level disappears.

One might then hope to treat both the Coulomb interaction and the Chern-Simons gauge field interactions among the fluctuations (beyond the mean field level) by standard many body perturbation techniques (e.g. by the random phase approximation, RPA).

Unfortunately, there is no small parameter for a many body perturbation expansion unless α , the number of Chern-Simons flux quanta attached to each particle, is small compared to unity.

However, a Landau-Silin [22] type Fermi liquid approach can take account of the short range correlations phenomenologically. A number of excellent papers on anyon superconductivity [23] treat Chern-Simons gauge interactions by standard many body techniques.

Halperin, Lee, and Read [11] have treated the half filled Landau level as a liquid of composite Fermions moving in zero effective magnetic field. Their RPA-Fermi-liquid approach gives a surprisingly satisfactory account of the properties of that state.

The vector potential associated with fluctuations beyond the mean field level is given by

$$\delta \mathbf{a}(\mathbf{r}) = \mathbf{a}(\mathbf{r}) - \langle \mathbf{a}(\mathbf{r}) \rangle.$$

The perturbation to the mean field Hamiltonian contains both linear and quadratic terms in $\delta \mathbf{a}(\mathbf{r})$, resulting in both two body - containing

$$\rho(\mathbf{r}_1)\rho(\mathbf{r}_2)$$

and three body - containing

$$\rho(\mathbf{r}_1)\rho(\mathbf{r}_2)\rho(\mathbf{r}_3)$$

- interaction terms. The three body interaction terms are usually ignored, though for α of the order of unity this approximation is of questionable validity.

5. Jain's composite fermion picture

Jain noted that in the mean field approximation, an effective filling factor ν^* of the composite Fermions was related to the electron filling factor ν by the relation

$$\frac{1}{\nu^*} = \frac{1}{\nu} - 2p.$$

Remember that $1/\nu$ is equal to the number of flux quanta of the applied magnetic field per electron, and $2p$ is the (even) number of Chern-Simons flux quanta (oriented opposite to the applied magnetic field) attached to each electron in the Chern-Simons transformation.

The above equation implies that when

$$\nu^* = \pm 1, \pm 2, \pm 3, \dots$$

(negative values correspond to the effective magnetic field B^* seen by the composite fermions oriented opposite to B) and a non-degenerate mean field composite fermion ground state occurs, then

$$\nu = \frac{\nu^*}{1 + 2p\nu^*}$$

This Jain sequence of condensed states ($\nu=1/3, 2/5, 3/7, \dots$ and $\nu=2/3, 3/5, \dots$ for $p=1$) is the set of fractional quantum Hall states most prominent in experiment.

When ν^* is not an integer, QP's of the neighboring Jain state will occur.

It is quite remarkable that the mean field composite fermion picture predicts not only the Jain sequence of incompressible ground states, but the correct band of low energy states for any value of the applied magnetic field. This is very nicely illustrated for the case of N electrons on a Haldane sphere.

When the monopole strength seen by an electron has the value $2S$, the effective monopole strength seen by a composite fermion is

$$2S^* = 2S - 2p(N - 1).$$

This equation reflects the fact that a given composite fermion senses the vector potential produced by the Chern-Simons flux on all other particles, but not its own Chern-Simons flux.

In Tab. 2 the ten particle system is described for a number of values of $2S$ between 29 and 15.

Tab. 2 The effective composite fermion monopole strength $2S^$, the numbers of composite fermion quasiparticles (quasielectrons, N_{QE} , and quasiholes, N_{QH}), the angular momentum of the lowest composite fermion shell l^* , the composite fermion and electron filling factors ν^* and ν , and the angular momenta L of the lowest lying band of multiplets for a ten electron system at $2S$ between 29 and 15.*

$2S$	29	28	27	26	25	21	15
$2S^*$	11	10	9	8	7	3	-3
N_{QH}	2	1	0	0	0	0	0
N_{QE}	0	0	0	1	2	6	6
l_0^*	11/2	5	9/2	4	7/2	3/2	3/2
ν^*			1			2	-2
ν			1/3			2/5	2/3
L	0, 2, 4, 6, 8, 10	5	0	5	0, 2, 4, 6, 8	0	0

The Laughlin $\nu=1/3$ state occurs at

$$2S_3 = 3(N-1) = 27.$$

For values of $2S$ different from this value,

$$2S - 2S_3 = \pm N_{QP}$$

(“+” corresponds to quasiholes, QH, and “-” to quasielectrons, QE).

Let us apply the composite fermion description to the ten electron spectra in Fig. 1. At $2S=27$, we take $p=1$ and attach two Chern-Simons flux quanta each electron. This gives $2S^*=9$ so that the ten composite fermions completely fill the $2S^*+1$ states in the lowest angular momentum shell (lowest Landau level).

There is a gap

$$\hbar\omega_c^* = \frac{eB^*}{\mu c}$$

to the next shell, which is responsible for the incompressibility of the Laughlin state.

Just as $|S|$ played the role of the angular momentum of the lowest shell of electrons, $l^*=|S^*|$ plays the role of the composite fermion angular momentum and $2|S^*|+1$ is the degeneracy of the composite fermion shell. Thus, the states with $2S=26$ and 28 contain a single quasielectron (QE) and quasihole (QH), respectively.

For the QE state, $2S^*=8$ and the lowest shell of angular momentum $l_0^*=4$ can accommodate only nine composite fermions. The tenth is the QE in the $l_1^*=l_0^*+1=5$ shell, giving the total angular momentum $L=5$.

For the QH state, $2S^*=10$ and the lowest shell can accommodate eleven composite fermions each with angular momentum $l_0^*=5$. The one empty state (QH) gives $L=l_0^*=5$.

For $2S=25$ we obtain $2S^*=7$, and there are two QE's each of angular momentum $l_1^*=9/2$ in the first excited composite fermion shell. Adding the angular momenta of the two QE's gives the band of multiplets $L=0, 2, 4, 6, \text{ and } 8$.

Similarly, for $2S=29$ we obtain $2S^*=11$, and there are two QH's each with $l_0^*=11/2$, resulting in the allowed pair states at $L=0, 2, 4, 6, 8, \text{ and } 10$.

At $2S=21$, the lowest shell with $l_0^*=3/2$ can accommodate only four composite fermions, but the other six composite fermions exactly fill the excited $l_1^*=5/2$ shell. The resulting incompressible ground state is the Jain $\nu=2/5$ state, since $\nu^*=2$ for the two filled shells.

A similar argument leads to $\nu^*=-2$ (minus sign means B^* oriented opposite to B) and $\nu=2/3$ at $2S=15$.

At $2S=30$, the addition of three QH angular momenta of $l_0^*=6$ gives the following band of low lying multiplets $L=1, 3^2, 4, 5^2, 6^2, 7^2, 8, 9^2, 10, 11, 12, 13, \text{ and } 15$. As demonstrated on an example in Fig. 1, this simple mean field composite fermion picture correctly predicts the band of low energy multiplets for any number of electrons N and for any value of $2S$.

6. Energy scales and pseudopotentials

The mean field composite Fermion picture is remarkably successful in predicting the low energy multiplets in the spectrum of N electrons on a Haldane sphere.

It was suggested originally that this success resulted from the cancellation of the Coulomb and Chern-Simons gauge interactions among fluctuations beyond the mean field level.

In Fig. 2, we show the lowest bands of multiplets for eight non-interacting electrons and for the same number of non-interacting mean field composite fermions at $2S=21$.

The energy scale associated with the Chern-Simons gauge interactions which convert the electron system in frame (a) to the composite fermion system in frame (b) is

$$\hbar\omega_c^* \propto B .$$

The energy scale associated with the electron-electron Coulomb interaction is

$$\frac{e^2}{\lambda} \propto \sqrt{B} .$$

Coulomb interaction lifts the degeneracy of the non-interacting electron bands in frame (a).

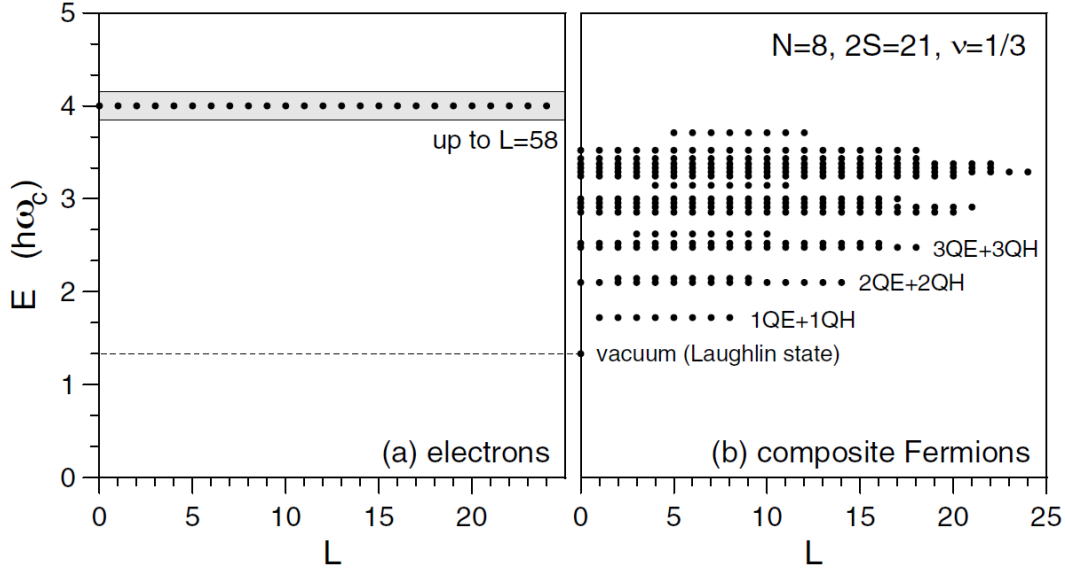


Fig. 2 The energy spectra of eight (a) non-interacting electrons and (b) non-interacting composite Fermions. The characteristic energy of the Coulomb interaction is marked in frame (a) with a shaded rectangle.

However, for very large value of B the Coulomb energy can be made arbitrarily small compared to the Chern-Simons energy (as marked with a shaded rectangle in Fig. 2), i.e. to the separation between the composite fermion Landau levels.

The energy separations in the mean field composite fermion model are completely wrong, but the structure of the low lying states (i.e., which angular momentum multiplets form the low lying bands) is very similar to that of the fully interacting electron system and completely different from that of the non-interacting electron system.

6.1 Two-fermion problem

An intuitive picture of why this occurs can be obtained by considering the two Fermion problem. The relative (REL) motion of a pair of electrons (ij) is described by a coordinate

$$z_{ij} = z_i - z_j = r_{ij} e^{i\phi_{ij}},$$

and for the lowest Landau level its wave function contains a factor

$$z_{ij}^m$$

where $m=1, 3, 5, \dots$

If every pair of particles has identical behavior, the many particle wave function must contain a similar factor for each pair giving a total factor

$$\prod_{i<j} z_{ij}^m.$$

As we have seen, the highest power of z_i in this product is $m(N-1)$.

If $m(N-1)$ is equal to $N\phi-1=2S$, the maximum value of the z-component of the single particle angular momentum, the Laughlin $\nu=1/m$ wave function results.

For electrons, the m^{th} cyclotron orbit, whose radius is r_m , encloses a flux

$$m\phi_0,$$

that is

$$\pi r_m^2 B = m\phi_0.$$

For a Laughlin $\nu=1/m$ state the pair function must have a radius

$$r_m = r_1 \sqrt{m}.$$

Let us describe the composite fermion orbits by radius

$$\rho_{\tilde{m}}$$

and require that the m^{th} orbit enclose m flux quanta. It is apparent that if a flux tube carrying two flux quanta (oriented opposite to the applied magnetic field B) is attached to each electron in the Chern-Simons transformation of the $\nu=1/3$ state, the smallest orbit of radius

$$\rho_{\tilde{m}=1}$$

has exactly the same size as

$$r_{m=3}.$$

Both orbits enclose three flux quanta of the applied field, but the composite fermion orbit also encloses the two oppositely oriented Chern-Simons flux quanta attached to the electrons to form the composite fermions.

In the absence of electron-electron interactions, the energies of these orbits are unchanged, since they still belong to the degenerate single particle states of the lowest Landau level.

In the mean field approximation the Chern-Simons fluxes are replaced by a spatially uniform magnetic field, leading to an effective field $B^*=B/m$.

The orbits for the composite fermion pair states in the mean field approximation are exactly the same as those of the exact Chern-Simons Hamiltonian. The smallest orbit has radius

$$\rho_{\tilde{m}=1}$$

equivalent to the electron orbit

$$r_{m=3}.$$

However, in the mean field approximation, the energies are changed (because

$$\hbar\omega_c^*$$

replaces

$$\hbar\omega_c.$$

This energy change leads to completely incorrect mean field composite fermion energies, but the mean field composite fermion orbitals give the correct structure to the low lying set of multiplets.

In the presence of a repulsive interaction, the low lying energy states will have the largest possible value of m . For a monopole strength $2S=m(N-1)$, where m is an odd integer, every pair can have radius r_m and avoid the large repulsion associated with r_1, r_3, \dots, r_{m-2} .

These ideas can be made somewhat more rigorous by using methods of atomic and nuclear physics for studying angular momentum shells of interacting Fermions.

6.2 Two-body interaction pseudopotential

As first suggested by Haldane [8], the behavior of the interacting many electron system depends entirely on the behavior of the two body interaction pseudopotential, which is defined as the interaction energy V of a pair of electrons as a function of their pair angular momentum.

In spherical geometry, in order to allow for meaningful comparison of the pseudopotentials obtained for different values of $2S$ (and thus different single electron angular momenta l), it is convenient to use the “relative” angular momentum $\mathcal{R}=2l-L_{12}$ rather than L_{12} (the length of

$$\hat{L}_{12} = \hat{l}_1 + \hat{l}_2.$$

The pair states with a given $\mathcal{R}=m$ (an odd integer) obtained on a sphere for different $2S$ are equivalent and correspond to the pair state on a plane with the relative (REL) motion described by angular momentum m and radius r_m .

The pair state with the smallest allowed orbit (and largest repulsion) has $\mathcal{R}=1$ on a sphere or $m=1$ on a plane, and larger \mathcal{R} and m means larger average separation.

In the limit of

$$\frac{\lambda}{R} \rightarrow 0$$

(i.e., either $2S \rightarrow 0$ or $R \rightarrow \infty$), the pair wave functions and energies calculated on a sphere for $\mathcal{R}=m$ converge to the planar ones (ψ_{0m} and its energy).

The pseudopotentials $V(\mathcal{R})$ are plotted in Fig. 3 for a number of values of the monopole strength $2S$.

The open circles mark the pseudopotential calculated on a plane ($\mathcal{R}=m$). At small \mathcal{R} the pseudopotentials rise very quickly with decreasing \mathcal{R} (i.e. separation). More importantly, they increase more quickly than linearly as a function of $L_{12}(L_{12}+1)$. The pseudopotentials with this property form a class of so-called “short range” repulsive pseudopotentials [12].

If the repulsive interaction has short range, the low energy many body states must, to the extent that it is possible, avoid pair states with the smallest values of \mathcal{R} (or m) and the maximum two body repulsion.

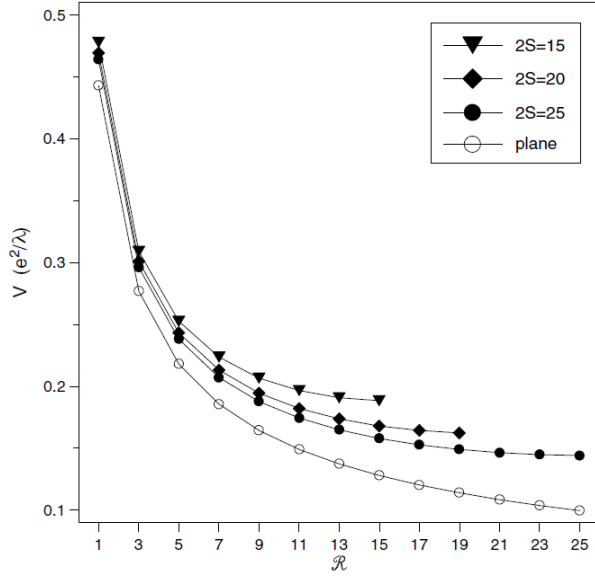


Fig. 3 The pseudopotentials of the Coulomb interaction in the lowest Landau level calculated on a Haldane sphere with $2S=15$, 20, and 25 (solid triangles, diamonds, and circles, respectively), and on a plane (open circles).

6.3 Fractional grandparentage

It is well-known in atomic and nuclear physics that eigenfunction of an N Fermion system of total angular momentum L can be written as

$$|l^N, L\alpha\rangle = \sum_{L_{12}} \sum_{L'\alpha'} G_{L\alpha, L'\alpha'}(L_{12}) |l^2, L_{12}; l^{N-2}, L'\alpha'; L\rangle.$$

Here, the totally antisymmetric state

$$|l^N, L\alpha\rangle$$

is expanded in the basis of states

$$|l^2, L_{12}; l^{N-2}, L'\alpha'; L\rangle$$

which are antisymmetric under permutation of particles 1 and 2 (which are in the pair eigenstate of angular momentum L_{12}) and under permutation of particles 3, 4, ..., N (which are in the $N-2$ particle eigenstate of angular momentum L'). The labels α (and α') distinguish independent states with the same angular momentum L (and L').

The expansion coefficient

$$G_{L\alpha, L'\alpha'}(L_{12})$$

is called the coefficient of fractional grandparentage (CFGP).

For a simple three Fermion system, the expansion reduces to

$$|l^3, L\alpha\rangle = \sum_{L_{12}} F_{L\alpha}(L_{12}) |l^2, L_{12}; l; L\rangle$$

and

$$F_{L\alpha}(L_{12})$$

is called the coefficient of fractional parentage (CFP).

In the lowest Landau level, the individual Fermion angular momentum l is equal to S , half the monopole strength, and the number of independent multiplets of angular momentum L that can be formed by addition of angular momenta of three identical Fermions is given in Tab. 3

Tab. 3 The number of times an L multiplet appears for a system of three electrons of angular momentum l . Top: even values of $2l$; bottom: odd values of $2l$. Blank spaces are equivalent to zeros.

$2l$	$2S = 0$	2	4	6	8	10	12	14	16	18	20	22	24	26
2	1													
4		1		1										
6	<u>1</u>		1	1	1		1							
8		<u>1</u>		2	1	1	1	1		1				
10	<u>1</u>		<u>1</u>	<u>1</u>	2	1	2	1	1	1	1		1	
12		<u>1</u>		<u>2</u>	<u>1</u>	2	2	2	1	2	1	1	1	1
14	<u>1</u>		<u>1</u>	<u>1</u>	<u>2</u>	<u>1</u>	3	2	2	2	2	1	2	1
$2l$	$2S = 1$	3	5	7	9	11	13	15	17	19	21	23	25	27
3		1												
5		1	1		1									
7		<u>1</u>	1	1	1	1		1						
9		<u>1</u>	<u>1</u>	1	2	1	1	1	1		1			
11		<u>1</u>	<u>1</u>	<u>1</u>	2	2	1	2	1	1	1	1		1
13		<u>1</u>	<u>1</u>	<u>1</u>	<u>2</u>	2	2	2	2	1	2	1	1	1

Low energy many body states must, to the possible extent, avoid parentage from pair states with the largest repulsion (pair states with maximum angular momenta L_{ij} or minimum \mathcal{R}).

In particular, we expect that the lowest energy multiplets will avoid parentage from the pair state with $\mathcal{R}=1$. If $\mathcal{R}=1$ (i.e. $L_{12}=2l-1$), then the smallest possible value of the total angular momentum L of the three Fermion system is obtained by addition of vectors \mathbf{L}_{12} (of length $2l-1$) and \mathbf{l}_3 (of length l), and it is equal to

$$|(2l-1)-l| = l-1.$$

Therefore, the three particle states with $L < l-1$ must not have parentage from $\mathcal{R}=1$.

It is straightforward to show that if

$$L < l - (2p-1)$$

where $p=1, 2, 3, \dots$, the three electron multiplet at L has no fractional parentage from

$$\mathcal{R} \leq 2p-1.$$

The multiplets that must avoid one, two, or three smallest values of \mathcal{R} are underlined with an appropriate number of lines in Tab.3 and listed in Tab. 4.

This gives the results in Tab. 4, the values of $2L$ that avoid $\mathcal{R}=1, 3,$ and 5 for various $2l$.

Tab. 4 The allowed values of $2L$ for a three electron system that must have $\mathcal{R} \geq 3$, $\mathcal{R} \geq 5$, and $\mathcal{R} \geq 7$. The listed $2L$ values correspond to the underlined L multiplets in Tab. 3.

$2l$	$2L (\mathcal{R} \geq 3)$	$2L (\mathcal{R} \geq 5)$	$2L (\mathcal{R} \geq 7)$
6	0		
7	3		
8	2		
9	3, 5		
10	0, 4, 6	0	
11	3, 5, 7	3	
12	2, 6 ² , 8	2	
13	3, 5, 7, 9 ²	3, 5	
14	0, 4, 6, 8 ² , 10	0, 4, 6	0

The $L=0$ states that appear at $2l=6$ ($\mathcal{R} \geq 3$), $2l=10$ ($\mathcal{R} \geq 5$), and $2l=14$ ($\mathcal{R} \geq 7$) are the only states for these values of $2l$ that can avoid one, two, or three largest pseudopotential parameters, respectively, and therefore are the non-degenerate ($L=0$) ground states. They are the Laughlin $\nu=1/3, 1/5,$ and $1/7$ states.

If only a single multiplet belongs to an angular momentum subspace, its form is completely determined by the requirement that it is an eigenstate of angular momentum with a given eigenvalue L . The wave function and the type of many body correlations do not depend on the form of the interaction pseudopotential. For interactions that do not have short range, the state that avoids the largest two body repulsion (e.g. the $L=0$ multiplet at $2l=6$) might not have the lowest total three body interaction energy and be the ground state.

If more than one multiplet belongs to a given angular momentum eigenvalue (e.g., two multiplets occur at $L=3$ for $2l=8$), the interparticle interaction must be diagonalized in this subspace (two-dimensional for $2l=8$ and $L=3$).

Whether the lowest energy eigenstate in this subspace has Laughlin type correlations, i.e. avoids as much as possible largest two body repulsion, depends critically on the short range of the interaction pseudopotential.

For the Coulomb interaction, we find that the Laughlin correlations occur and, whenever possible, the coefficient of fractional parentage of the lowest lying multiplets virtually vanishes (it would vanish exactly for an “ideal” short range pseudopotential which increases infinitely quickly with decreasing \mathcal{R}).

For example, for the lower energy eigenstate at $L=3$ and $2l=8$, the coefficient of fractional parentage for $\mathcal{R}=1$ is less than 10^{-3} . A similar thing occurs at $2S=9$ for $L=9/2$, at $2S=10$ for $L=4$ and 6 , at $2S=11$ for $L=9/2, 11/2,$ and $15/2$, at $2S=12$ for $L=5, 6, 7,$ and 9 , at $2S=13$ for $L=11/2, 13/2, 15/2, 17/2,$ and $21/2$, and at $2S=14$ for $L=6^2, 7, 8, 9, 10,$ and 12 .

At $2S=14$ for $L=6$ there are three allowed multiplets. The diagonalization of the Coulomb interaction gives the lowest state that avoids $\mathcal{R}=1$ (coefficient of fractional parentage $\sim 10^{-7}$) and $\mathcal{R}=3$ (coefficient of fractional parentage $< 10^{-2}$), and the next lowest state that avoids $\mathcal{R}=1$ (coefficient of fractional parentage $< 10^{-5}$) but orthogonality to the lowest state requires that it has significant parentage from $\mathcal{R}=3$ (coefficient of fractional parentage ≈ 0.34).

One can see that the set of angular momentum multiplets L that can be constructed at a given value of $2l$ without parentage from pair states with $\mathcal{R}=1$ is identical to the set of all allowed multiplets L at $2l^*=2l-4$.

For a short range repulsion (e.g. the Coulomb repulsion in the lowest Landau level), these multiplets will be (to a good approximation) the lowest energy eigenstates (the appropriate coefficient of fractional parentage for the actual eigenstates will be very small although not necessarily zero).

More generally, in the lowest Landau level (remember that $l=S$), the set of multiplets L that can be constructed at given $2S$ without parentage from $\mathcal{R}\leq 2p-1$ (i.e. with $\mathcal{R}\geq 2p+1$ for all pairs; $p=1, 2, \dots$) is identical to the set of all allowed multiplets L at

$$2S^* = 2S - 2p(N-1).$$

The multiplets L forming the lowest Coulomb energy band at a given $2S$ are all multiplets allowed at $2S^*$. But $2S^*=2S-2p(N-1)$ is just the effective magnetic monopole strength in the mean field composite fermion picture!

Thus the composite fermion picture with $2p$ attached flux quanta simply picks the subset of angular momentum multiplets which have no parentage from pair states with $\mathcal{R}\leq 2p-1$, and neglects the long range part of the pseudopotential, $V(\mathcal{R})$ for $\mathcal{R}\geq 2p+1$.

6.4 Definition of the short range pseudopotential

For systems containing more than three Fermions in an angular momentum shell, the simple addition of angular momentum to determine the smallest possible L that has parentage from pair states with $L_{12}=2l-1$ is of no help.

Instead, we make use of the following operator identity

$$\hat{L}^2 + N(N-2)\hat{l}^2 = \sum_{i<j} \hat{L}_{ij}^2$$

Here

$$\hat{L} = \sum_i \hat{l}_i$$

and

$$\hat{L}_{ij} = \hat{l}_i + \hat{l}_j$$

The identity is easily proved by writing out the expression for

$$\hat{L}^2$$

and for

$$\sum_{i<j} \hat{L}_{ij}^2$$

and eliminating

$$\sum_{i<j} \hat{l}_i \cdot \hat{l}_j$$

from the pair of equations.

Taking matrix elements of the operator identity between states

$$|l^N, L\alpha\rangle$$

gives

$$\begin{aligned} L(L+1) + N(N-2)l(l+1) &= \langle l^N, L\alpha | \sum_{i<j} \hat{L}_{ij}^2 | l^N, L\alpha \rangle \\ &= \frac{1}{2} N(N-1) \sum_{L_{12}} \mathcal{G}_{L\alpha}(L_{12}) L_{12}(L_{12}+1) \end{aligned}$$

where

$$\mathcal{G}_{L\alpha}(L_{12}) = \sum_{L'\alpha'} |G_{L\alpha, L'\alpha'}(L_{12})|^2.$$

The coefficients of grandparentage satisfy the relation

$$\sum_{L_{12}} \sum_{L'\alpha'} G_{L\alpha, L'\alpha'}(L_{12}) G_{L\beta, L'\alpha'}(L_{12}) = \delta_{\alpha\beta}$$

Of course, the energy of the multiplet

$$|l^N, L\alpha\rangle$$

is given by

$$E_\alpha(L) = \frac{1}{2} N(N-1) \sum_{L_{12}} \mathcal{G}_{L\alpha}(L_{12}) V(L_{12})$$

where $V(L_{12})$ is the electron pseudopotential.

It is important to make the following observations:

1. The expectation value of

$$\sum_{i<j} \hat{L}_{ij}^2$$

in a many body state

$$|l^N, L\alpha\rangle$$

increases as

$$L(L+1),$$

but it is totally independent of α ;

2. If the pseudopotential

$$V_H(L_{12})$$

were a linear function of

$$\hat{L}_{ij}^2$$

(we refer to V_H as the “harmonic pseudopotential”), all many body multiplets with the same value of L would be degenerate;

3. The difference

$$\Delta V(L_{12}) = V(L_{12}) - V_H(L_{12})$$

between the actual pseudopotential V and its harmonic part V_H lifts this degeneracy. If N_L many body multiplets of V_H occur at angular momentum L , the anharmonic term ΔV in the pseudopotential causes them to “repel one another” and results in a band of N_L non-degenerate multiplets.

Because the expectation value of

$$\sum_{i<j} \hat{L}_{ij}^2$$

in a many body state of angular momentum L increases as $L(L+1)$, a strict Hund's rule holds for harmonic pseudopotentials:

For V_H that increases as a function of L_{12} , the highest energy state is always at the maximum possible value of L equal to

$$L^{MAX} = Nl - \frac{1}{2}N(N-1),$$

and the lowest energy state is at the minimum allowed value of L equal to L^{MIN} .

If V_H decreases as a function of L_{12} , the opposite occurs:

The lowest energy state is at L^{MAX} , and the highest energy state is at L^{MIN} (this is a standard Hund's rule of atomic physics).

Neither of these Hund's rules may remain true in the presence of a large anharmonic term ΔV .

For example, if the number of multiplets N_L at a value slightly larger than L^{MIN} is very large compared to N_L at $L=L^{MIN}$, the strong level repulsion due to ΔV within this L subspace can overcome the difference in the expectation values of V_H , and the lowest eigenvalue of V at L can be lower than that at L^{MIN} .

However, only very few multiplets occur at large values of L :

$N_L=1$ for $L=L^{MAX}$ (for $M=L=L^{MAX}$, the only state is $|l, l-1, \dots, l-N+1\rangle$);

$N_L=0$ for $L=L^{MAX}-1$;

$N_L \leq 1$ for $L=L^{MAX}-2$;

$N_L \leq 1$ for $L = L^{\text{MAX}} - 3$; etc.

As a result, breaking of the Hund's rule that refers to the behavior of energy at large L requires stronger anharmonicity than at small L .

For the Coulomb pseudopotential in the lowest Landau level we always find that the highest energy indeed occurs at L^{MAX} .

However, the ability to avoid parentage from pair states having large L_{ij} often favors many body states at small $L > L^{\text{MIN}}$ with large N_L , as prescribed by the composite fermion picture.

The anharmonicity of the Coulomb pseudopotential in the lowest Landau level (which increases with increasing L_{12}) is critical for the behavior of the fractional quantum Hall systems. We have found that the condition for the occurrence of subbands separated by gaps in the energy spectrum, and, in particular, for the occurrence of non-degenerate incompressible fluid ground states at specific values of the filling factor, is that the anharmonic term $\Delta V(L_{12})$ is positive and increases with increasing L_{12} .

In other words, the total pseudopotential $V(L_{12})$ must increase more quickly than linearly as a function of $L_{12}(L_{12}+1)$.

6.5 Hidden symmetry of the short range repulsion

From our numerical studies we have arrived at the following conjectures:

1. The Hilbert space

$$\mathcal{H}_{NI}$$

of N identical Fermions each with angular momentum l contains subspaces

$$\mathcal{H}_{NI}^{(p)}$$

of states that have no parentage from $\mathcal{R} \leq 2p-1$.

The subspaces

$$\tilde{\mathcal{H}}_{NI}^{(p)} = \mathcal{H}_{NI}^{(p)} \setminus \mathcal{H}_{NI}^{(p+1)}$$

can be defined; they hold states without parentage from $\mathcal{R} \leq 2p-1$, but with some parentage from $\mathcal{R} = 2p+1$.

Then

$$\mathcal{H}_{NI} = \tilde{\mathcal{H}}_{NI}^{(0)} \oplus \tilde{\mathcal{H}}_{NI}^{(1)} \oplus \tilde{\mathcal{H}}_{NI}^{(2)} \oplus \dots$$

2. For an “ideal” short range repulsive pseudopotential V_{SR} , for which

$$V_{SR}(\mathcal{R}) \gg V_{SR}(\mathcal{R} + 2)$$

the huge difference between energy scales associated with different pair states results in the following (dynamical) symmetry:

- (a) subspaces

$\tilde{\mathcal{H}}_{NI}^{(p)}$ are the interaction eigensubspaces,

- (b) p is a good quantum number,
- (c) energy spectrum splits into bands (larger p corresponds to lower energy), and
- (d) energy gap above the p th band scales as

$$V(2p-2) - V(2p)$$

3. For a finite short range pseudopotential V (increasing more quickly than V_H as a function of L_{12}), the above symmetry is only approximate, but the correlation between energy and parentage from highly repulsive pair states persists, and so do the gaps in the energy spectrum.

The mixing between neighboring subbands is weak, although the structure of energy levels within each subband depends on the form of $V(L_{12})$ at $\mathcal{R} \geq 2p+1$.

4. The set of angular momentum multiplets in subspace

$$\mathcal{H}_{NI}^{(p)}$$

is identical to

$$\mathcal{H}_{NI^*}$$

where $I^* = I - p(N-1)$.

A general analytic proof for the last conjecture can be found in [24].

All of the above conjectures can be immediately translated into the planar geometry.

The harmonic pseudopotential $V_H(m)$, used to define the class of short range pseudopotentials, is that of a repulsive interaction potential $V(r)$ which is linear in r^2 .

Then,

$$\mathcal{H}_\nu = \tilde{\mathcal{H}}_\nu^{(0)} \oplus \tilde{\mathcal{H}}_\nu^{(1)} \oplus \tilde{\mathcal{H}}_\nu^{(2)} \oplus \dots$$

where

$$\mathcal{H}_\nu$$

is the Hilbert space of electrons filling a fraction ν of an infinitely degenerate Landau level, and subspaces

$$\tilde{\mathcal{H}}_\nu^{(p)}$$

contain states without parentage from $m \leq 2p-1$, but with some parentage from $m=2p+1$.

The (approximate) dynamical symmetry holds for the Coulomb interaction, and the low energy band

$$\mathcal{H}_\nu^{(p)}$$

contains the same angular momentum multiplets as

$$\mathcal{H}_{\nu^*},$$

with ν^* defined by the composite fermion prescription.

The validity of our conjectures for systems interacting through the Coulomb pseudopotential is illustrated in Fig. 4 for four electrons in the lowest Landau level at $2S=5, 11, 17,$ and 23 . Different symbols mark bands corresponding to (approximate) subspaces

$$\tilde{\mathcal{H}}_{NI}^{(p)}$$

with different p . The same sets of multiplets reoccur for different $2S$ in bands related by

$$\mathcal{H}_{NI}^{(p)} \sim \mathcal{H}_{NI^*}$$

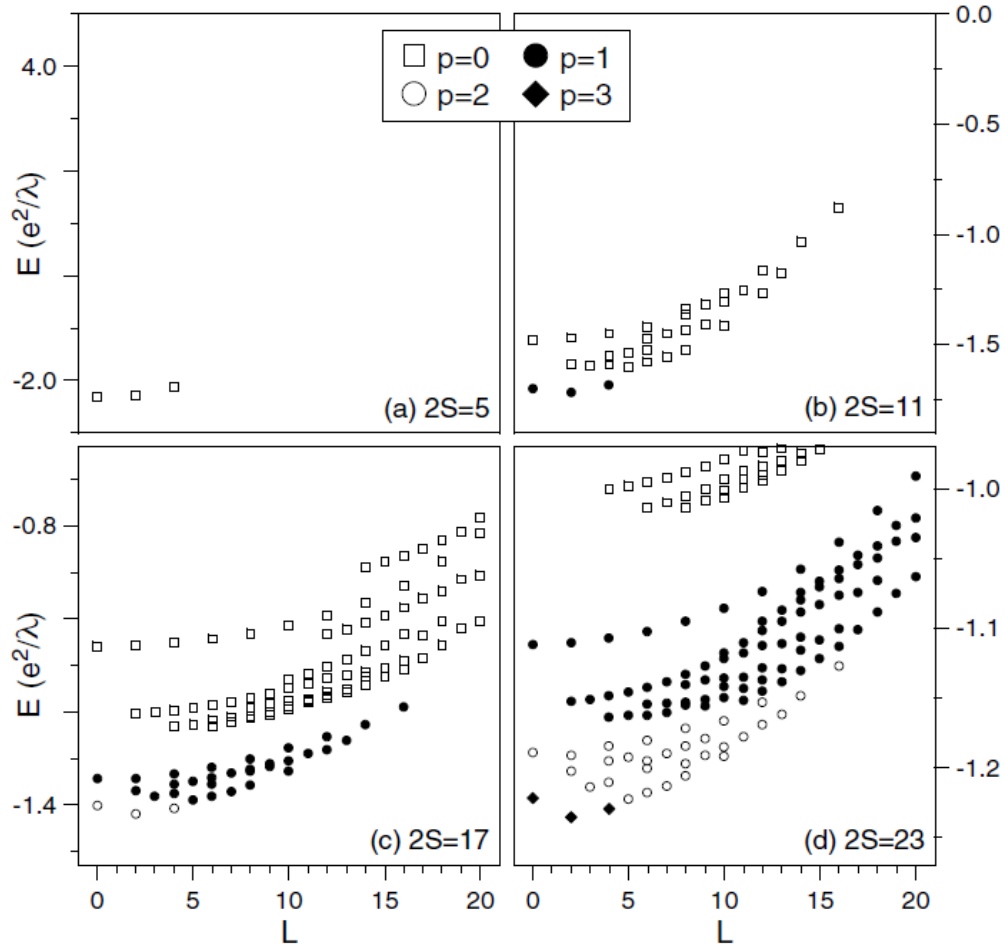


Fig. 4 The energy spectra of four electrons in the lowest Landau level calculated on a Haldane sphere with $2S=5, 11, 17,$ and 23 . All those values of $2S$ are equivalent in the mean field composite Fermion picture (the Chern-Simons transformation with $p=0, 1, 2,$ and $3,$ respectively). Different symbols mark states with different numbers of avoided pair states with highest energy.

6.6 Comparison with atomic shells: Hund's rule

Our conjectures (verified by the numerical experiments) are based on the behavior of systems of interacting Fermions partially filling a shell of degenerate single particle states of angular momentum l . This is a central problem in atomic physics and in nuclear shell model studies of

energy spectra. It is interesting to compare the behavior of the spherical harmonics of atomic physics with that of the monopole harmonics considered here.

For monopole harmonics $l=S+n$, where S is half of the monopole strength (and can be integral or half integral) and n is a non-negative integer. For the lowest angular momentum shell $l=S$. For spherical harmonics $S=0$ and $l=n$.

If in each case electrons are confined to a 2D spherical surface of radius R , one can evaluate the pair interaction energy V as a function of the pair angular momentum L_{12} .

The resulting pseudopotentials, $V(\mathcal{R})$ for the fractional quantum Hall system in the lowest Landau level, and $V(L_{12})$ for atomic shells in a zero magnetic field, are shown in Fig. 5 for a few small values of l .

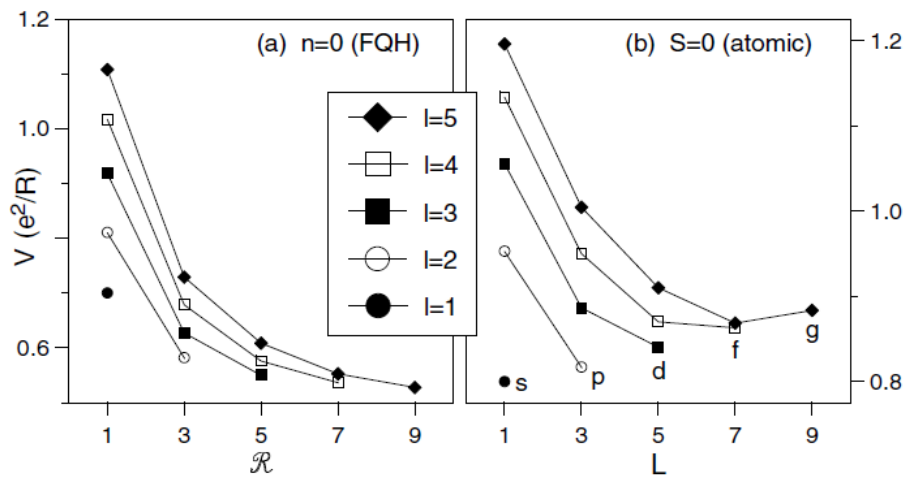


Fig. 5 The pseudopotentials V of the Coulomb interaction for the pair of electrons each with angular momentum l : (a) lowest Landau level on a Haldane sphere, monopole harmonics, $n=0$, $l=S$, V plotted as a function of relative pair angular momentum \mathcal{R} ; (b) atomic shell, spherical harmonics, $S=0$, $l=n$, V plotted as a function of pair angular momentum L .

In obtaining these results we have restricted ourselves to spin-polarized shells, so only orbital angular momentum is considered. It is clear that in the case of spherical harmonics the largest pseudopotential coefficient occurs for the lowest pair angular momentum, exactly the opposite of what occurs for monopole harmonics.

As a consequence of the relation between the total angular momentum L and the average pair angular momentum L_{12} , the standard atomic Hund's rule predicts that the energy of a few electron system in an atomic shell will, on the average, decrease as a function of total angular momentum, which is opposite to the behavior of energy of electrons in the lowest Landau level.

The difference between the energy spectra of electrons interacting through atomic and fractional quantum Hall pseudopotentials of Fig. 5 is demonstrated in Fig. 6, where we plot the result for four electrons in shells of angular momentum $l=3$ and 5.

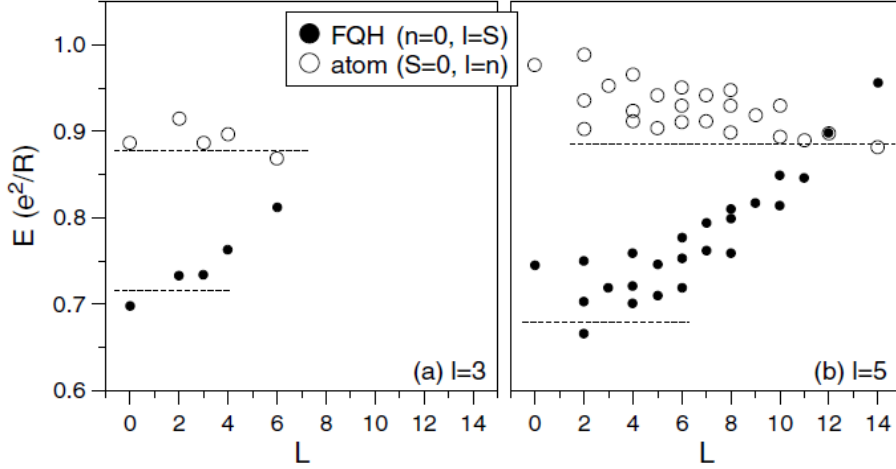


Fig. 6 The energy spectra of four electrons in a degenerate shell of angular momentum $l=3$ (a) and $l=5$ (b), interacting through the pseudopotentials of Fig. 5: open circles - atomic shell ($S=0$ and $l=n$), solid circles - lowest Landau level ($n=0$ and $l=S$).

The solid circles correspond to monopole harmonics and the open ones to spherical harmonics. Note that at L^{MAX} the former give the highest energy and the latter the lowest.

Due to anharmonicity of the pseudopotentials, the behavior of energy at low L does not always follow a simple Hund's rule for either fractional quantum Hall or atomic system.

The fractional quantum Hall ground state for $l=3$ occurs at $L=0$ (this is the Jain $\nu=2/3$ incompressible state). However, for $l=5$, the lowest of the three states at $L=2$ has lower energy than the only state at $L=0$. This ground state at $L=2$ contains one quasihole in the Laughlin $\nu=1/3$ state and it is the only four electron state at this filling in which electrons can avoid parentage from the $\mathcal{R}=1$ pair state.

Exactly opposite happens for the atomic system at $l=5$, where the anharmonicity is able to push the highest of the three $L=2$ states above the high energy state at $L=0$.

6.7 Higher Landau levels

Thus far we have considered only the lowest angular momentum shell (lowest Landau level) with $l=S$. The interaction of a pair of electrons in the n^{th} excited shell of angular momentum $l=S+n$ can easily be evaluated to obtain the pseudopotentials $V(L_{12})$ shown in Fig. 7.

Here we compare $V_n(L_{12})$ as a function of $L_{12}(L_{12}+1)$ for $n=0, 1$, and 2 .

It can readily be observed that $V_{n=0}$ increases more quickly than $L_{12}(L_{12}+1)$ in entire range of L_{12} , while $V_{n=1}$ and $V_{n=2}$ do so only up to certain value of L_{12} (i.e., above certain value of $\mathcal{R}=2l-L_{12}$).

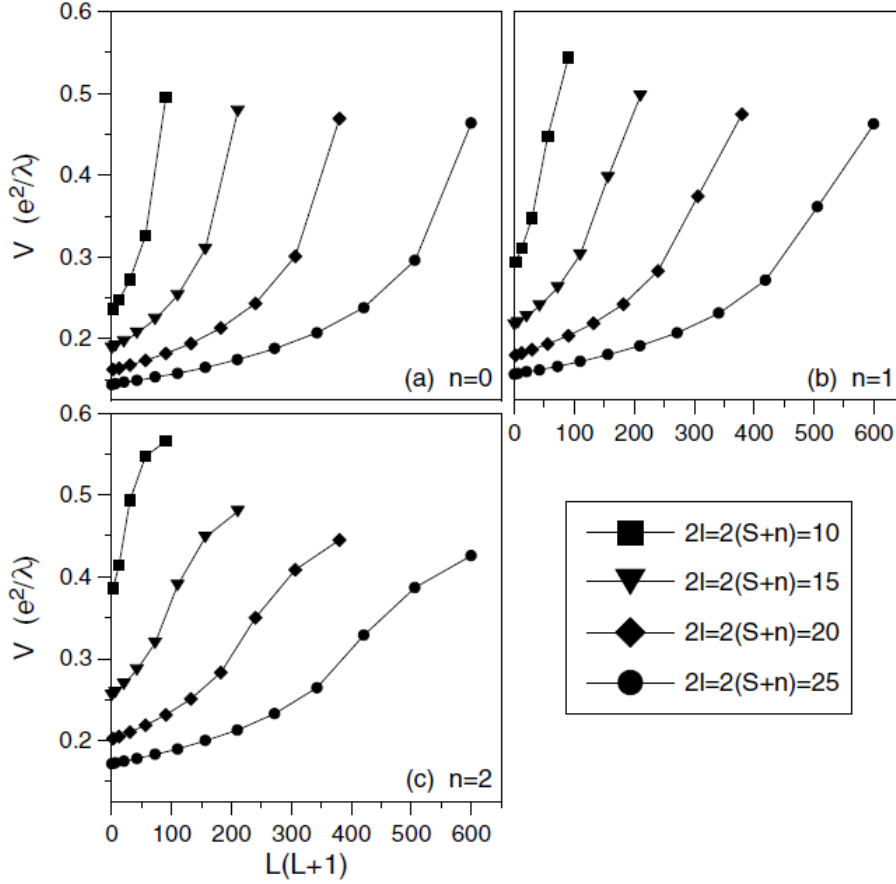


Fig7 The pseudopotentials V of the Coulomb interaction in the lowest (a), first excited (b), and second excited (c) Landau levels, calculated on a Haldane sphere, as a function of squared pair angular momentum $L(L+1)$. Different symbols correspond to different Landau level degeneracy $2l+1$.

For $n=1$, $V_{n=1}$ has short range for $\mathcal{R} \geq 3$ but is essentially linear in $L_{12}(L_{12}+1)$ from $\mathcal{R}=1$ to 5.

For $n=2$, $V_{n=2}$ has short range for $\mathcal{R} \geq 5$ but is sublinear in $L_{12}(L_{12}+1)$ from $\mathcal{R}=1$ to 7.

More generally, we find that the pseudopotential in the n th excited shell (Landau level) has short range for $\mathcal{R} \geq 2n+1$.

Because the conclusions of the composite fermion picture depend so critically on the short range of the pseudopotential, they are not expected to be valid for all fractional fillings of higher Landau levels.

For example, the ground state at $\nu=2+1/3=7/3$ does not have Laughlin type correlations (i.e. electrons in the $n=1$ Landau level do not avoid parentage from $\mathcal{R}=1$) even if it is non-degenerate ($L=0$) and incompressible (as found experimentally [25]).

7. Fermi Liquid model of composite fermions

The numerical results of the type shown in Fig. 1 have been understood in a very simple way using Jain's composite Fermion picture.

For the ten particle system, the Laughlin $\nu=1/3$ incompressible ground state at $L=0$ occurs for $2S=3(N-1)=27$.

The low lying excited states consist of a single QP pair, with the QE and QH having angular momenta $l_{QE}=11/2$ and $l_{QH}=9/2$. In the mean field composite fermion picture, these states should form a degenerate band of states with angular momentum $L=1, 2, \dots, 10$.

More generally,

$$l_{QE} = \frac{1}{2}(N+1)$$

and

$$l_{QH} = \frac{1}{2}(N-1)$$

for the Laughlin state of an N electron system, and the maximum value of L is N .

The energy of this band would be

$$E = \hbar\omega_c^* = \frac{1}{3}\hbar\omega_c$$

which is the effective composite fermion cyclotron energy needed to excite one composite fermion from the (completely filled) lowest to the (completely empty) first excited composite fermion Landau level.

From the numerical results, two shortcomings of the mean field composite fermion picture are apparent.

First, due to the QE-QH interaction (neglected in the composite fermion picture), the energy of the QE-QH band depends on L , and the “magnetoroton” QE-QH dispersion has a minimum at $L=5$. Second, the state at $L=1$ either does not appear, or is part of the continuum (in an infinite system) of higher energy states above the magnetoroton band.

At $2S=27-1=26$ and $2S=27+1=28$, the ground state contains a single quasiparticle (QE or QH, respectively), whose angular momenta

$$l_{QE} = l_{QH} = \frac{1}{2}N = 5$$

result from the Chern-Simons transformation which gives

$$2S^* = 2S - 2(N-1)$$

which is 8 for QE and 10 for QH, and

$$l_{QE} = S^* + 1$$

and

$$l_{QH} = S^*.$$

States containing two identical QP's form lowest energy bands at $2S=25$ (two QE's) and $2S=29$ (two QH's). The allowed angular momenta of two identical composite fermion QP's (which are Fermions) each with angular momentum l_{QP} are $L=2l_{QP}-j$ where j is an odd integer.

Of course, l_{QP} depends on $2S$ in the composite fermion picture, and at $2S=25$ we have

$$l_{QE} = S^* + 1 = S - (N - 1) + 1 = \frac{9}{2}$$

yielding $L=0, 2, 4, 6,$ and $8,$ while at $2S=29$ we have

$$l_{QH} = S^* = S - (N - 1) = \frac{11}{2}$$

and $L=0, 2, 4, 6, 8, 10.$

More generally,

$$l_{QE} = \frac{N - 1}{2}$$

and

$$l_{QH} = \frac{N + 1}{2}$$

in the 2QE and 2QH states of an N electron system, and the maximum values of L are $N-2$ for QE's and N for QH's.

Similarly as for the magnetoroton band at $2S=27,$ the composite fermion picture does not account for QP interactions and incorrectly predicts the degeneracy of the bands of 2QP states at $2S=25$ and $27.$

The energy spectra of states containing more than one composite fermion quasiparticle can be described in the following phenomenological Fermi liquid picture.

The creation of an elementary excitation, QE or QH, in a Laughlin incompressible ground state requires a finite energy, ϵ_{QE} or $\epsilon_{QH},$ respectively.

In a state containing more than one Laughlin quasiparticle, QE's and/or QH's interact with one another through the appropriate QE-QE, QH-QH, and QE-QH pseudopotentials.

An estimate of the QP energies can be obtained by comparing the energy of a single QE (for the $N=10$ electron system, the energy of the ground state at $L=N/2=5$ for $2S=27-1=26$) or a single QH ($L=N/2=5$ ground state at $2S=27+1=28$) with the Laughlin $L=0$ ground state at $2S=27.$ There can be finite size effects here, because the QP states occur at different values of $2S$ than the ground state, but using the correct magnetic length

$$\lambda = \frac{R}{\sqrt{S}}$$

(R being the radius of the sphere) in the unit of energy

$$\frac{e^2}{\lambda}$$

at each value of $2S,$ and extrapolating the results as a function of $1/N$ to an infinite system should give reliable estimates of ϵ_{QE} and ϵ_{QH} for a macroscopic system.

The QP pseudopotentials $V_{\text{QP-QP}}$ can be obtained by subtracting from the energies of the 2QP states obtained numerically at $2S=25$ (2QE), $2S=27$ (QE-QH), and $2S=29$ (2QH), the energy of the Laughlin ground state at $2S=27$ and two energies of appropriate non-interacting QP's.

As for the single QP, the energies calculated at different $2S$ must be taken in correct units of

$$\frac{e^2}{\lambda} = \sqrt{S} \frac{e^2}{R}$$

to avoid finite size effects. This procedure was carried out in references [16,26].

In Fig. 8 we plot the QE-QE and QH-QH pseudopotentials for Laughlin $\nu=1/3$ and $1/5$ states.

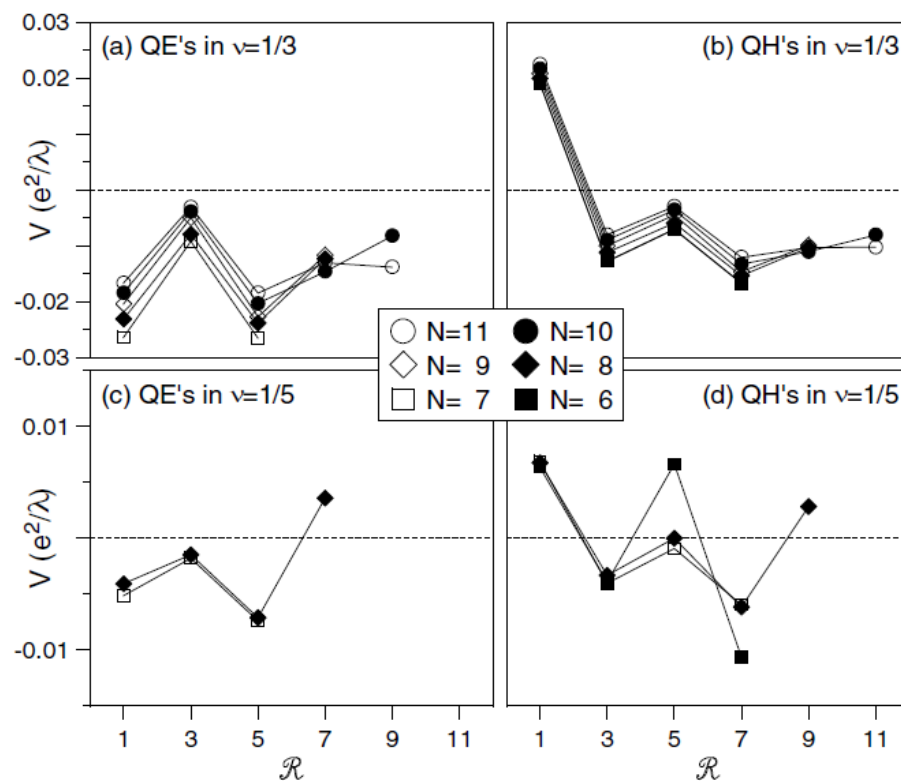


Fig. 8 The pseudopotentials of a pair of quasielectrons (left) and quasiholes (right) in Laughlin $\nu=1/3$ (top) and $\nu=1/5$ (bottom) states, as a function of relative pair angular momentum \mathcal{R} . Different symbols mark data obtained in the diagonalization of between six and eleven electrons.

As we have seen for two electrons (see Fig. 3), the angular momentum L_{12} of a pair of identical Fermions in an angular momentum shell (or a Landau level) is quantized, and the convenient quantum number to label the pair states is $\mathcal{R}=2l_{\text{QP}}-L_{12}$ (on a sphere) or relative (REL) angular momentum m (on a plane). When plotted as a function of \mathcal{R} , pseudopotentials calculated for systems containing between six to eleven electrons (and thus for different QP angular momenta l_{QP}) behave similarly and, for $N \rightarrow \infty$ (i.e., $2S \rightarrow \infty$), they seem to converge to the limiting pseudopotentials $V_{\text{QP-QP}}(\mathcal{R}=m)$ describing an infinite planar system.

In Fig. 9 we plot the QE-QH pseudopotentials for Laughlin $\nu=1/3$ and $1/5$ states.

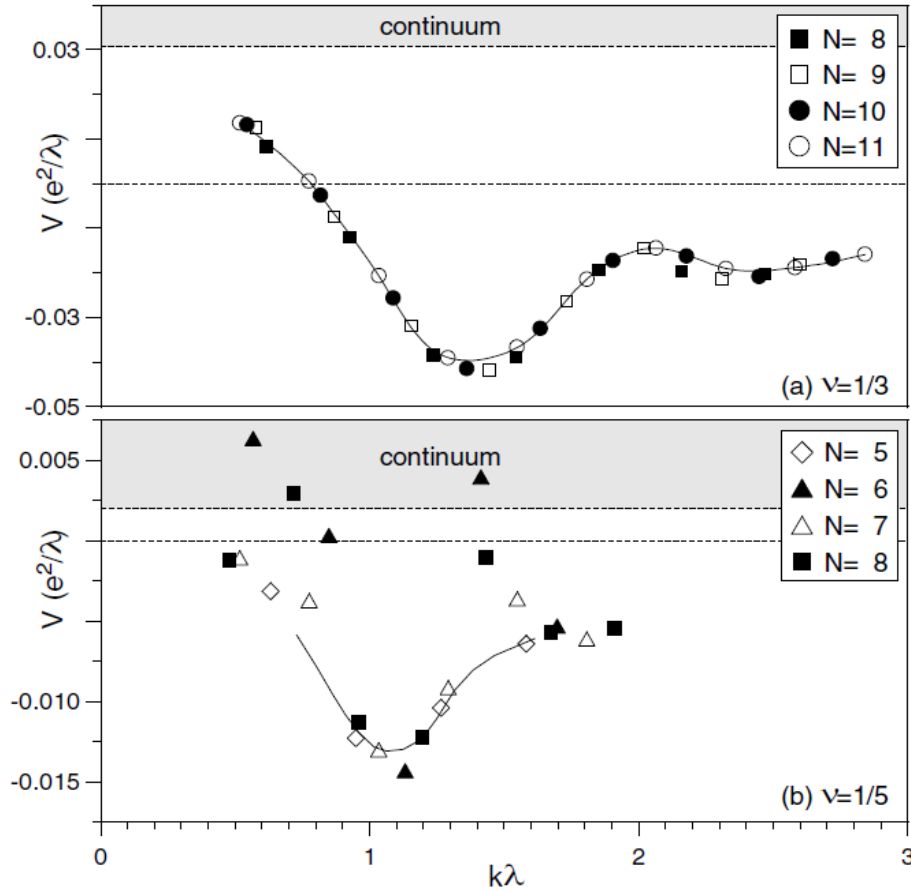


Fig. 9 The pseudopotentials of a quasielectron-quasihole pair in Laughlin $\nu=1/3$ (a) and $\nu=1/5$ (b) states as a function of wavevector k . Different symbols mark data obtained in the diagonalization of between five and eleven electrons.

As for a conduction electron and a valence hole pair in a semiconductor (an exciton), the motion of a QE-QH pair which does not carry a net electric charge is not quantized in a magnetic field. The appropriate quantum number to label the states is the continuous wavevector k (or momentum), which on a sphere is given by

$$k = \frac{L}{R} = \frac{L}{\lambda\sqrt{S}}$$

(remember that L is given in units of the Planck's constant).

When plotted as a function of k , the pseudopotentials calculated for systems containing between six to eleven electrons fall on the same curve that describes a continuous magnetoroton dispersion $V_{\text{QE-QH}}(k)$ of an infinite planar system (the lines in Fig. 9 are only to guide the eye). Only the energies for $L \geq 2$ are shown in Fig. 9, since the single QE-QH pair state at $L=1$ is either disallowed (hard core) or falls into the continuum of states above the magnetoroton band.

The magnetoroton minima for the Laughlin $\nu=1/3$ and $1/5$ states occur at about

$$k_0 = \frac{1.4}{\lambda}$$

and

$$k_0 = \frac{1.1}{\lambda},$$

respectively.

The magnoroton band at $\nu=1/3$ is well decoupled from the continuum of higher states because the band width $\sim 0.05e^2/\lambda$ is much smaller than the energy gap $\varepsilon_{QE}+\varepsilon_{QH}=0.1e^2/\lambda$ for additional QE-QH pair excitations.

At $\nu=1/5$, the band width $\sim 0.015 e^2/\lambda$ is closer to the single particle gap $\varepsilon_{QE}+\varepsilon_{QH}=0.021e^2/\lambda$ and the state of two magnorotons each with $k\approx k_0$ can couple to the highest energy QE-QH pair states at $k\leq 2k_0$.

Knowing the QP-QP pseudopotentials and the bare QP energies allows us to evaluate the energies of states containing three or more QP's. Typical results are shown in Fig. 10.

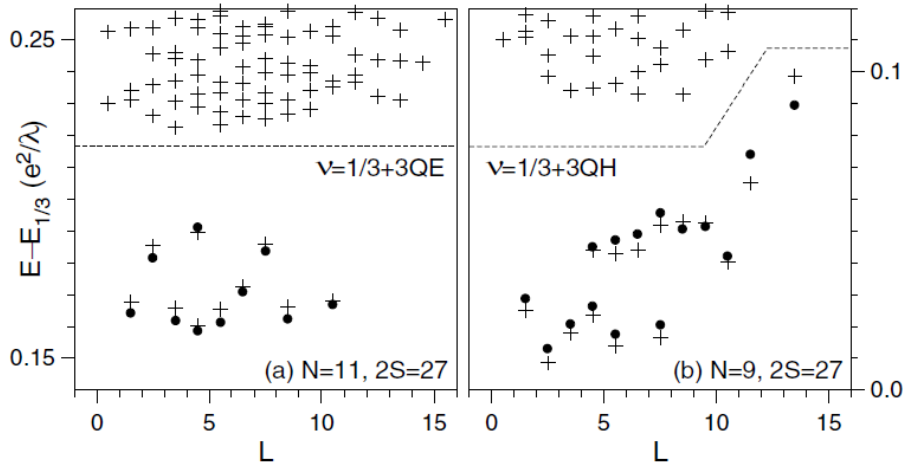


Fig. 10 The energy spectra of three quasielectrons (a) and three quasiholes in the Laughlin $\nu=1/3$ state. The crosses correspond to the Fermi-liquid calculation using pseudopotentials from Fig. 8(a,b); the solid circles give exact spectra obtained in full diagonalization of the Coulomb interaction of eleven (a) and nine (b) electrons.

In frame (a) we show the energy spectra of three QE's in the Laughlin $\nu=1/3$ state of eleven electrons. The spectrum in frame (b) gives energies of three QH's in the nine electron system at the same filling. The exact numerical results obtained in diagonalization of the eleven and nine electron systems are represented by plus signs and the Fermi liquid picture results are marked by solid circles. The exact energies above the dashed lines correspond to higher energy states that contain additional QE-QH pairs.

It should be noted that in the mean field composite fermion picture which neglects the QP-QP interactions, all of the 3QP states would be degenerate and the energy gap separating the 3QP states from higher states would be equal to

$$E = \hbar\omega_c^* = \frac{1}{3}\hbar\omega_c .$$

Although the fit in Fig. 10 is not perfect, it is quite good and justifies the use of the Fermi liquid picture to describe (compressible) states at

$$\nu \neq \frac{1}{2p+1} .$$

8. Composite fermion hierarchy

The sequence of Laughlin-Jain states with filling factor ν given by

$$\nu = \frac{\nu^*}{1 + 2p\nu^*}$$

where $p=1, 2, \dots$, and the composite fermion filling factor ν^* is any non-zero integer, is the most prominent set of condensed states observed experimentally. However, this sequence (together with the conjugate “hole” states, $\nu \rightarrow 1-\nu$) does not contain all odd denominator fractions the way the Haldane hierarchy scheme does.

The question arises quite naturally of how to treat the composite fermion values of ν^* which are not integers. The answer leads to the composite fermion approach to the hierarchy of incompressible quantum fluid ground states [13].

Consider a state of N_0 electrons at a monopole strength $2S_0$ with a filling factor ν_0 . The Chern-Simons transformation that attaches to each electron $2p_0$ flux quanta oriented opposite to the applied magnetic field results in the composite fermion system at an effective filling factor ν_0^* given by

$$\frac{1}{\nu_0^*} = \frac{1}{\nu_0} - 2p_0$$

and an effective monopole strength

$$2S_0^* = 2S_0 - 2p_0(N_0 - 1)$$

The procedure for handling non-integral values of composite fermion filling factor ν_0^* is to set it equal to

$$\nu_0^* = n_1 + \nu_1$$

where n_1 is an integer and ν_1 is the fractional filling of the composite fermion quasiparticle level (same sign as n_1 for QE's and opposite for QH's).

The problem is then that of placing N_1 quasiparticles into $2l_1+1$ available states of a composite fermion shell (Landau level) of angular momentum l_1 :

the QE's into the lowest empty shell with

$$l_1 = |S_0^*| + n_1 + 1$$

or the QH's into the highest filled shell with

$$l_1 = |S_0^*| + n_1$$

We now ignore all completely filled and completely empty composite fermion shells, and reapply the Chern-Simons transformation by setting $S_1=l_1$ and attaching $2p_1$ flux quanta to each of the N_1 quasiparticles in the partially filled composite fermion shell.

This produces a new type of QP's and a new QP filling factor ν_1^* given by

$$\frac{1}{\nu_1^*} = \frac{1}{\nu_1} - 2p_1.$$

If ν_1^* is an integer, we obtain a daughter states in the hierarchy. If it is not, we write

$$\nu_1^* = n_2 + \nu_2$$

where ν_2 represents the partial filling of the new QP shell, and repeat the mean field composite fermion procedure.

This leads to the set of equations:

$$\frac{1}{\nu_l} = 2p_l + \frac{1}{n_{l+1} + \nu_{l+1}}$$

where ν_l is the QP filling factor and $2p_l$ is the number of flux quanta attached to each Fermion at the l^{th} level of the composite fermion hierarchy.

As an example, consider a system of $N_0=12$ electrons at $2S_0=30$.

We apply the mean field composite fermion approximation by attaching to each electron $2p_0=2$ flux quanta. This gives the effective composite fermion monopole strength $2S_0^*=30-2(12-1)=8$. The lowest composite fermion shell is filled with nine particles, and there are $N_1=3$ quasielectrons in the first excited ($n_1=1$) composite fermion shell of angular momentum $l_1=5$. The filling factor at this level of hierarchy is $\nu_0^*=1+\nu_1$.

We now reapply the composite fermion transformation by attaching $2p_1=4$ flux quanta to each of $N_1=3$ QE's at $2S_1=10$ and obtain $2S_1^*=10-4(3-1)=2$. The lowest composite fermion shell of $l_1=1$ is now completely filled yielding $\nu_1^*=1$.

Using the appropriate above filling factor relations we obtain

$$\frac{1}{\nu_1} = 4 + \frac{1}{1} = 5$$

and

$$\frac{1}{\nu_0} = 2 + \frac{1}{1+1/5} = \frac{17}{6}$$

If the mean field composite fermion picture worked on all levels of hierarchy, the twelve electron system at $2S=30$ should have an incompressible $L=0$ ground state corresponding to the filling factor $\nu=6/17$.

In Fig. 11(a) we show the low energy sector of the spectrum calculated for this system using the Fermi liquid picture (only the lowest energy states containing three QE's in the Laughlin $\nu=1/3$ state are calculated).

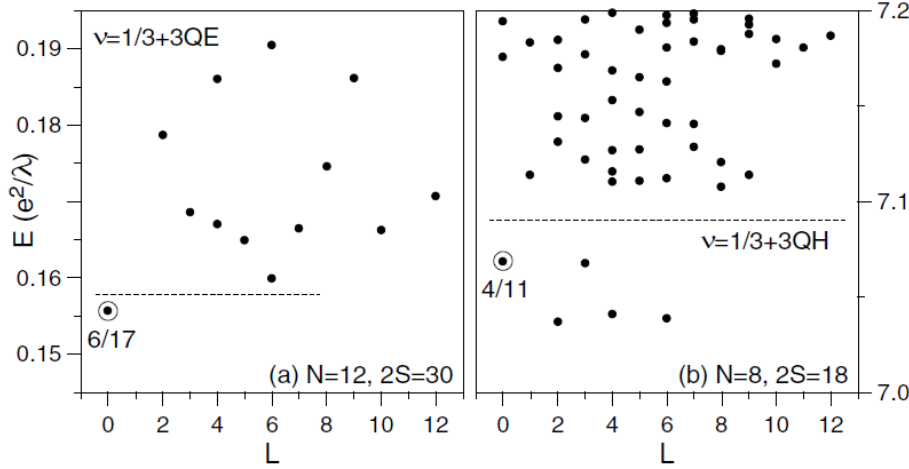


Fig. 11 (a) The low energy spectrum of three quasielectrons in the Laughlin $\nu=1/3$ state of twelve electrons calculated using quasielectron pseudopotential from Fig. 8(a). (b) The energy spectrum of three quasiholes in the Laughlin $\nu=1/3$ state of eight electrons obtained in full diagonalization of electron-electron Coulomb interaction.

Indeed, the $\nu=6/17$ hierarchy ground state at $L=0$ is separated from higher states by a small gap in the twelve electron spectrum (although it is not clear that this small gap will persist in the thermodynamic limit [26]).

Though the composite fermion hierarchy picture seems to work in some cases, there are others where it is clearly in complete disagreement with numerical results. For example, a composite fermion transformation with $2p_0=2$ applied to an $N_0=8$ electron system at $2S_0=18$ gives $2S_0^*=18-2(8-1)=4$, $n_1=1$, and $N_1=3$ QE's left in the shell with $l_1=3$.

Adding the three QE angular momenta of $l_1=3$ gives a low energy band at $L=0, 2, 3, 4$, and 6 . Reapplication of the composite fermion transformation with $2p_1=2$ gives $2S_1^*=6-2(3-1)=2$, i.e. the completely filled lowest shell, $\nu_1^*=1$ ($n_2=1$ and $\nu_2=0$).

From this we get $\nu_1=1/3$ and $\nu_0=4/11$.

In Fig. 11(b) we show the spectrum obtained by exact numerical diagonalization of an eight electron system at $2S=18$. It is apparent that the set of multiplets at $L=0, 2, 3, 4$, and 6 form the low energy band. However the reapplication of the mean field composite fermion transformation to the three QE's in the $l_1=3$ shell (which predicts an $L=0$ incompressible ground state corresponding to $\nu=4/11$) is definitely wrong.

The reason why the composite fermion hierarchy picture does not always work is not difficult to understand. The electron (Coulomb) pseudopotential in the lowest Landau level $V_e(\mathcal{R})$ satisfies the “short range” criterion (i.e. it increases more quickly with decreasing \mathcal{R} than the harmonic pseudopotential V_H) in the entire range of \mathcal{R} , which is the reason for the incompressibility of the principal Laughlin $\nu=1/(2p+1)$ states. However, this does not generally hold for the QP pseudopotentials on higher levels of the hierarchy.

In Fig. 8 we plotted $V_{QE-QE}(\mathcal{R})$ and $V_{QH-QH}(\mathcal{R})$ for the $\nu=1/3$ and $\nu=1/5$ Laughlin states of six to eleven electrons. Clearly, the QE and QH pseudopotentials are quite different and neither one decreases monotonically with increasing \mathcal{R} . On the other hand, the corresponding pseudopotentials in $\nu=1/3$ and $1/5$ states look similar, only the energy scale is different.

The convergence of energies at small \mathcal{R} obtained for larger N suggests that the maxima at $\mathcal{R}=3$ for QE's and at $\mathcal{R}=1$ and 5 for QH's, as well as the minima at $\mathcal{R}=1$ and 5 for QE's and at $\mathcal{R}=3$ and 7 for QH's, persist in the limit of large N (i.e. for an infinite system on a plane).

Consequently, the only incompressible daughter states of Laughlin $\nu=1/3$ and $1/5$ states are those with $\nu_{QE}=1$ or $\nu_{QH}=1/3$ and (maybe) $\nu_{QE}=1/5$ and $\nu_{QH}=1/7$.

It is clear that no incompressible daughter states of the parent Laughlin $\nu=1/3$ state will form at e.g. $\nu=4/11$ ($\nu_{QE}=1/3$) or $4/13$ ($\nu_{QH}=1/5$), but that they will form (at least, in finite systems [26]) at $\nu=6/17$ ($\nu_{QE}=1/5$) or $6/19$ ($\nu_{QH}=1/7$).

From the composite fermion hierarchy scheme we find the Jain-Laughlin states when the Chern-Simons transformation is applied directly to electrons (or to holes in a more than half-filled level). These states occur at integral values of ν^* , the effective composite fermion filling factor, and correspond to completely filling a QP shell.

For example, the $\nu=2/5$ state occurs when $\nu^*=2$, and the composite fermions in the first excited shell (which are Laughlin QE's of the $\nu=1/3$ state) have $\nu_{QP}=1$.

The angular momenta of the two lowest composite fermion shells are

$$l_0^* = |S^*|$$

and

$$l_1^* = |S^*| + 1$$

so they contain $2l_0^*+1$ and $2l_1^*+1$ states, respectively.

Since $\nu_{QP}=1$, there are $N_{QP}=2l_1^*+1$ composite fermion quasiparticles. The total number of states filled by the N Fermions is

$$(2l_0^*+1) + (2l_1^*+1) = 2N_{QP} - 2$$

giving

$$N = 2N_{QP} - 2.$$

For an infinite system this is just Haldane's relation between the number of quasiparticles and the number of electrons,

$$N = 2qN_{QP}$$

for the integer $q=1$.

This demonstrates that integrally filled composite fermion shells correspond to $\nu_{QP}=1$, a completely filled shell of Laughlin QP's.

Adding new Fermions to a system with $\nu_{QP}=1$ requires creating a new type of QP's, and the counting of available QP states turns out to be exactly the same in the composite fermion hierarchy and Haldane's Boson hierarchy pictures.

Integral composite fermion filling (i.e., $\nu_{QP}=1$) gives a valid mean field picture independent of QP-QP interactions provided that the gap for creating new QP's is positive.

When ν^* is non-integral, the mean field picture is valid only at values of L for which the “short range” requirement on the pseudopotential $V_{\text{QP-QP}}(L)$ is satisfied.

The form of the QP-QP interactions obtained from our numerical calculations makes it clear that the mean field approximation is not valid at certain quasiparticle fillings (e.g. for $\nu_{\text{QP}}=1/3$ filling of the quasielectron levels of the electron $\nu=1/3$ state).

II. Electron-hole systems in high magnetic fields

1. Excitonic complexes

There has been a great deal of interest in photoluminescence (PL) of 2D systems in high magnetic fields. An important ingredient in understanding PL is the negatively charged exciton, alternatively also called a trion (X^-).

The trion consists of two electrons bound to a valence band hole. If the total spin of the pair of electrons, J_e , is zero, the trion is said to be a singlet

$$X_s^-;$$

if $J_e=1$ the trion is called a triplet

$$X_t^-.$$

Only the singlet is bound in the absence of a magnetic field, but in infinite magnetic field (so that only a single Landau level is relevant) only the triplet is bound in a 2D system.

It often occurs that the photoexcited hole is separated from the plane of the electron system by a small distance (this can happen, e.g., in wide GaAs quantum wells when the electron gas is confined to one GaAs/AlGaAs interface by remote ionized donors, and the photoexcited holes reside close to the other GaAs/AlGaAs interface).

Several remarkable effects associated with electron-hole systems and charged excitons can be understood using the composite Fermion picture.

2. Charged exciton and the hidden symmetry in the lowest Landau level

First let us consider the idealized 2D system at so large a magnetic field that only the lowest electron and hole Landau levels need be considered. The energy spectrum for a two-electron-one-hole system at $2S=20$ is shown in Fig. 12.

The triplet trion with angular momentum

$$l_{X^-} = S - 1$$

is the only bound state, with binding energy $\sim 0.05e^2/\lambda$. A pair of (unbound) singlet and triplet states occur at the energy equal exactly to the exciton energy E_X . In these so-called “multiplicative” states a neutral exciton X in its ground state is decoupled from the second electron.

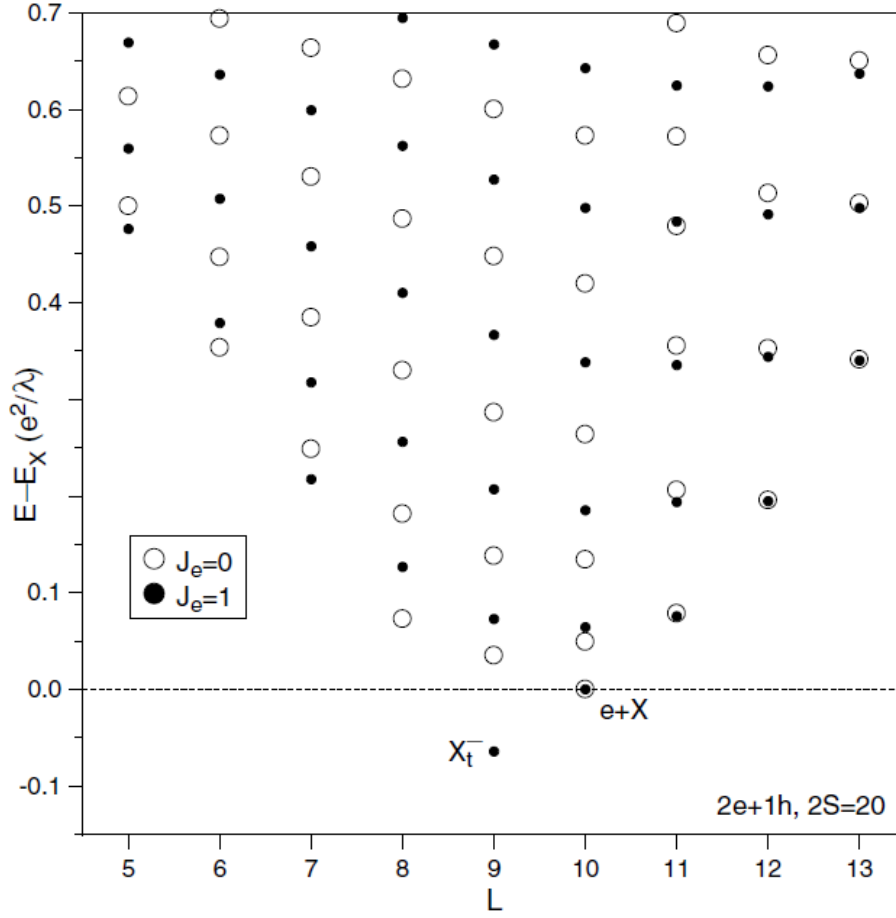


Fig12 The energy spectrum (binding energy vs. angular momentum) of a two-electron-one-hole system in the lowest Landau level at $2S=20$. Open and solid circles mark singlet and triplet spin configurations, respectively.

Addition of exciton and electron angular momenta $L_X=0$ and $l_e=S$ gives a state of total angular momentum $L=S$, and addition of two electron spins of $1/2$ gives both $J_e=0$ and 1 spin configurations.

3. Hidden symmetry

The occurrence of unbound $2e+1h$ states at $E=E_X$ and $L=S$ is a manifestation of the following “hidden symmetry:”

Because of the exact overlap of electron and hole orbitals in the lowest Landau level (scaled with the same magnetic length λ), and thus independence of the strength of interaction of the type of particles involved, the commutator of an operator

$$d_X^+$$

which creates an exciton in its $L_X=0$ ground state (on a sphere,

$$d_X^+ = \sum_m (-1)^m c_m^+ h_m^+,$$

where c_m^+ and h_m^+ are electron and hole creation operators), with the interaction Hamiltonian H is

$$[H, d_X^+] = E_X d_X^+.$$

As a result, if Ψ is an eigenstate of N_e electrons and N_h holes with an eigenenergy E and angular momentum quantum numbers L and M , then the multiplicative state

$$d_X^+ \Psi$$

N_e+1 electrons and N_h+1 holes is also an eigenstate with energy $E+E_X$ and the same L and M .

A good quantum number conserved due to the “hidden symmetry” is the number of decoupled excitons, N_X .

In particular, the ground state for $N_e=N_h=N$ is the totally multiplicative state

$$(d_X^+)^N |vac\rangle$$

with $N_X=N$; for an infinite system this ground state can be viewed as a Bose condensate of non-interacting excitons.

It can be readily found that the application of the PL operator that annihilates an optically active exciton (d_X) reduces its N_X by one, and therefore that only the multiplicative electron-hole states with $N_X>0$ are optically active (have non-vanishing PL intensity).

In Fig. 12, the two multiplicative states at $E=E_X$ and $L=S$ have $N_X=1$, and all others have $N_X=0$.

It is essential to realize that two independent symmetries forbid the recombination of a triplet trion ground state in Fig. 12:

1. Due to the 2D translational/rotational space invariance, the PL operator d_X conserves two angular momentum quantum numbers. On a sphere, these are L and M , and the resulting optical selection rule allows only a state with $L=S$ to decay by e-h recombination. On a plane, these are the projection of total angular momentum M and an additional angular momentum quantum number K associated with partial decoupling of the center-of-mass motion of a charged system in a homogeneous magnetic field, and the radiative channel for a trion is that of $M+K=0$.

This (geometrical) symmetry can be broken by collisions, but persists in systems with a finite quantum well width, finite electron and hole layer separation, or Landau level mixing.

2. Due to the equal strength of e-e, h-h, and e-h interactions, N_X is a good quantum number. Since N_X is decreased in a PL process, only the multiplicative ($N_X>0$) states are radiative.

This (dynamical) symmetry is not broken by collisions, and requires breaking electron-hole orbital symmetry.

Since a number of independent factors are needed to allow for the recombination of a triplet trion, this complex (in narrow and symmetric quantum wells and in high magnetic fields) is expected to be a well defined long-lived quasiparticle. The correlations, optical properties,

etc. are expressed more easily in terms of this quasiparticle than in terms of individual electrons and holes.

The finite angular momentum of a trion in spherical geometry (partial decoupling of the center-of-mass excitations from the relative motion on a plane) can be viewed as the formation of a degenerate Landau level of this (charged) quasiparticle.

As will be shown later, the interaction of trion quasiparticles with one another and with electrons can be described using the ideas familiar in the context of fractional quantum Hall systems (Laughlin correlations, composite Fermions, parentage, etc.).

4. Interaction of charged excitons

The simplest system in which to study trion-trion interaction contains four electrons and two holes. Its energy spectrum at $2S=17$ is shown in Fig. 13.

The low energy spectrum is characterized by four bands which we have identified as follows:

1. The lowest band taking on all even values between $L=0$ and 12 consists of a pair of charged excitons, each with angular momentum

$$l_{X^-} = S - 1;$$

2. The next band contains an electron with $l_e=S$ and a negatively charged biexciton (a bound state of an exciton and a trion)

$$X_2^- = X + X^-$$

with angular momentum

$$l_{X_2^-} = S - 2;$$

the allowed angular momenta go from $L = |l_e - l_{X_2^-}| = 2$ to $L = l_e + l_{X_2^-} - 1 = 14$

3. A band of multiplicative states containing an exciton, a trion, and an electron; it begins at $L = |l_e - l_{X^-}| = 1$ and goes to $L = l_e + l_{X^-} - 1 = 15$.
4. A band of multiplicative states containing two neutral excitons and two free electrons; it takes on all even values of L between zero and $2l_e-1=16$.

One interesting feature of Fig. 13 is that it gives us the effective pseudopotential $V_{AB}(L)$ for the interaction of the pair of Fermions AB (where A and B can be an electron, charged exciton, charged biexciton, etc.) as a function of angular momentum.

As for electrons, it is convenient to use the relative pair angular momentum

$$\mathcal{R} = l_A + l_B - L$$

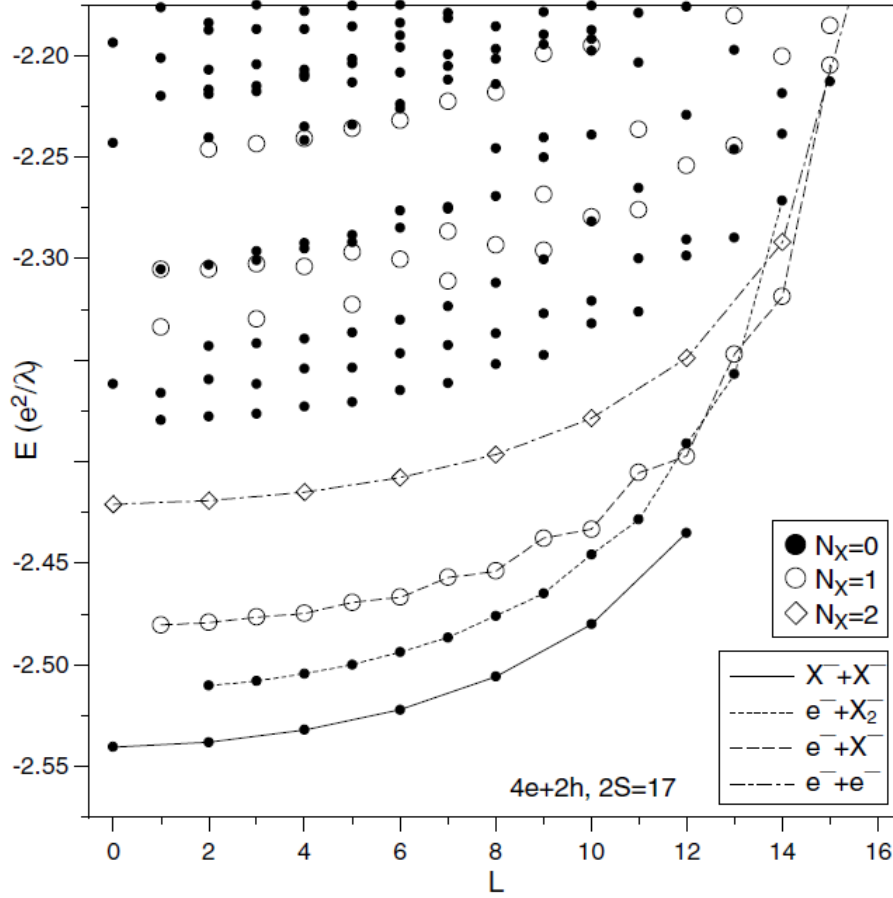


Fig. 13 The energy spectrum of a four-electron-two-hole system in the lowest Landau level calculated on a Haldane sphere with $2S=17$. Different symbols mark states with zero, one, or two decoupled excitons. The lines connect states identified as pseudopotentials of different pairs of bound charged complexes.

For identical Fermions with angular momentum l , the allowed values of L are $2l-j$, where j is an odd integer, i.e., $\mathcal{R}=1, 3, 5, \dots$, and $\mathcal{R} \leq 2l$.

For distinguishable Fermions A and B , all values of L between $|l_A-l_B|$ and $|l_A+l_B|$ are expected, i.e., $\mathcal{R}=0, 1, 2, \dots$, and $\mathcal{R} \leq 2\min(l_A, l_B)$.

However, our numerical results display a “hard core” repulsion for composite particles, and one or more of the pair states with the largest values of L (smallest \mathcal{R}) are forbidden (i.e. the corresponding pseudopotential parameters are effectively infinite).

For

$$A = X_n^-$$

and

$$B = X_m^-$$

the smallest allowed value of \mathcal{R} is given by

$$\mathcal{R}_{AB}^{MIN} = 2 \min(n, m) + 1$$

The identification of pair states AB in Fig. 13 (as marked with lines) was possible by comparing the displayed 4e-2h spectrum with the pseudopotentials of point charge particles with appropriate angular momenta l_A and l_B and binding energies ε_A and ε_B [17].

The appropriate values of angular momenta l_A and l_B , and of the binding energies ε_A and ε_B are obtained by diagonalizing smaller systems (e.g. the 2e-1h system in Fig. 12 for the trion), and the point charge pseudopotentials are used to approximate the AB interaction.

The approximate AB energies obtained in this way are rather close to the exact 4e-2h energies. This implies that, due to different energy scales, the internal dynamics of charged excitons is weakly coupled to their scattering off one another or off electrons, and allows for the interpretation of an electron-hole system in terms of well-defined charged excitonic quasiparticles interacting with one another and with excess electrons through Coulomb like forces.

Slight difference between the actual pseudopotentials in Fig. 13 and the pseudopotentials of point charge particles comes from the larger size of charged excitons and their (nearly frozen) internal degrees of freedom. The latter can be accounted for phenomenologically by attributing each type of composite particles a finite electric polarizability to describe their induced electric dipole moment in the presence of an electric field of other charged particles. Due to an increased charge isotropy, the polarization effects are expected to be greatly reduced in larger systems, and disappear completely in the fluid type states discussed in the following paragraphs.

5. Generalized composite fermion picture for charged excitons

Suppose we have a system of different (distinguishable) charged Fermions (A, B, ...). They can be distinguished either because they are different species (e.g., electrons and charged excitons) or because they are confined to different, spatially separated layers.

If all particles in such system repel one another through short range pseudopotentials (as defined for the electron fractional quantum Hall systems), one can think of many body states with Laughlin-type correlations [6,7] given by a generalized Laughlin-Jastrow factor

$$\prod_{i,j} (z_i^{(A)} - z_j^{(B)})$$

where z 's are the complex coordinates for the positions of the i^{th} Fermion of type A and j^{th} Fermion of type B, and the product is over all pairs.

The restrictions on the integers m_{AB} are that m_{AA} and m_{BB} must be odd, $m_{BA}=m_{AB}$, and m_{AB} must not be smaller than certain minimum values $\mathcal{R}_{AB}^{\text{MIN}}$ to avoid the infinite hard cores for all pairs.

In a state with Laughlin correlations, a number of pair states with largest repulsion are avoided for each pair,

$$\mathcal{R}_{AB} \geq m_{AB}$$

This is equivalent to saying that the high energy collisions (in which any pair of particles would come very close to one another) are forbidden in such state. This intuitive property of

the Laughlin fluid states will be very useful in the discussion of collision assisted trion recombination.

A generalized composite fermion picture can be constructed for a system with Laughlin correlations. In this picture, fictitious flux tubes carrying an integral number of flux quanta ϕ_0 are attached to each particle.

In the multi-component system, each particle of type A carries flux $(m_{AA}-1)\phi_0$ that couples only to charges on all other particles of the same type A, and fluxes $m_{AB}\phi_0$ that couple to charges on all particles of other types B (A and B are any of the types of Fermions).

On a sphere, the effective monopole strength seen by a composite fermion of type A (composite fermion-A) is

$$2S_A^* = 2S - \sum_B (m_{AB} - \delta_{AB})(N_B - \delta_{AB})$$

For different multi-component systems we expect generalized Laughlin incompressible states (for two components denoted as $[m_{AA}, m^{BB}, m^{AB}]$) when all the hard core pseudopotentials are avoided and composite fermions of each kind fill completely an integral number of their composite fermion shells (e.g. $N_A=2l_A^*+1$ for the lowest shell).

In other cases, the low lying multiplets are expected to contain different kinds of composite fermion quasiparticles (generalized QE's or QH's), QP-A, QP-B, ..., in the neighboring incompressible state.

It is interesting to realize that the effective monopole strengths $2S_A^*$, i.e. the effective magnetic fields B_A^* seen by particles of different type are not generally equal. One can think of effective Chern-Simons charges and fluxes of different colors, but the resulting number of different effective composite fermion magnetic fields of different color can no longer be regarded as physical reality, and no cancellation between gauge and Coulomb interactions is possible.

The multi-component (multi-color) composite fermion picture can be applied to electrons and charged excitons in an electron-hole system. We have checked that the pseudopotentials describing interaction of identical composite particles in Fig. 13 all satisfy the short range criterion in the entire range of \mathcal{R} .

For a pair of different particles, the pseudopotential may increase sufficiently quickly for some values of \mathcal{R} but not the others and, for example, for e- and X- only the correlations described by exponents

$$m_{e^-,X^-} = 1 \quad \text{or even}$$

are expected to occur.

As an example, let us consider the 12e-6h system. In Fig. 14 we present its low energy spectrum at $2S=17$, calculated by diagonalizing systems of different combinations of electrons and composite particles interacting through effective pseudopotentials determined in Fig. 13.

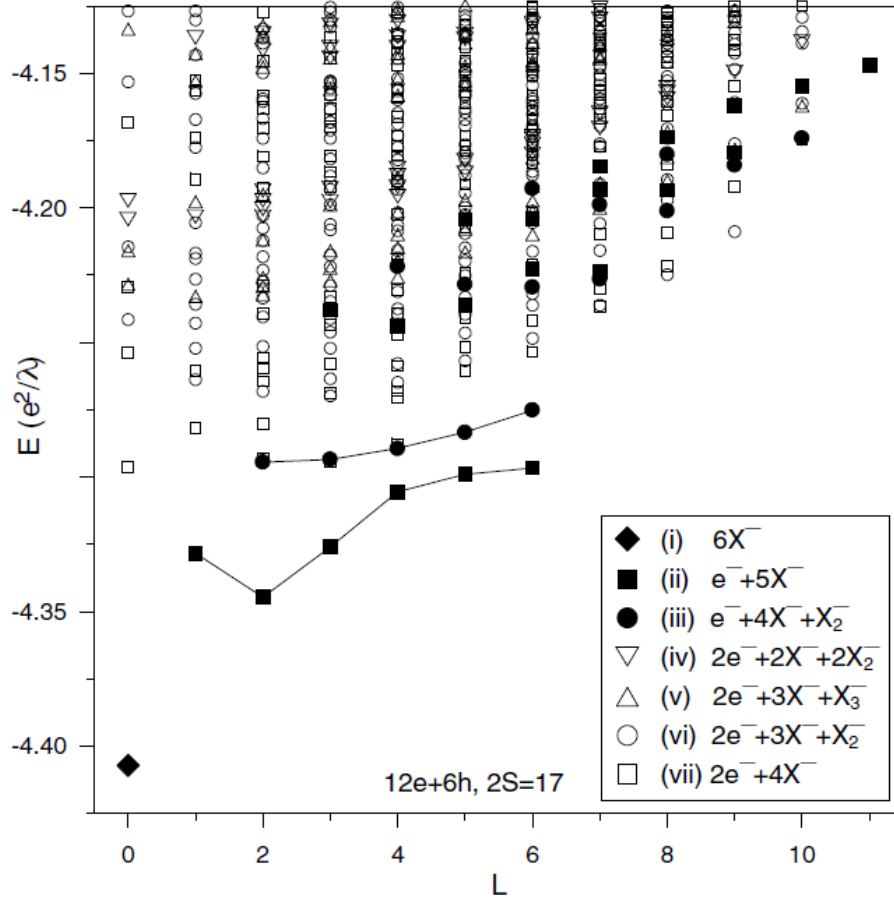


Fig. 14 The approximate lowest energy bands corresponding to different combinations of six bound charged complexes interacting through appropriate pseudopotentials, in the twelve-electron-six-hole spectrum in the lowest Landau level, calculated on a Haldane sphere with $2S=17$. The lines mark lowest subbands of two lowest excited bands.

The following combinations (groupings of 12e and 6h into bound complexes) have the highest total binding energy and thus form the lowest energy bands in the 12e-6h spectrum:

- (i) $6X^-$,
- (ii) $e^- + 5X^-$,
- (iii) $e^- + 4X^- + X_2^-$,
- (iv) $2e^- + 2X^- + 2X_2^-$,
- (v) $2e^- + 3X^- + X_3^-$,
- (vi) $2e^- + 3X^- + X_2^-$
- (vii) $2e^- + 4X^-$

Groupings (ii), (vi), and (vii) also contain neutral excitons that however do not interact with charged particles due to the hidden symmetry.

For each of these groupings, the composite fermion transformation can be applied to determine correlations and identify number and type of quasiparticles that occur in the lowest energy states. For example, for groupings (i)-(iii) the generalized composite fermion picture makes the following predictions.

1. For $m_{X^-X^-} = 3$ we obtain the Laughlin $\nu=1/3$ state with total angular momentum $L=0$.

Because of the hard core of $V_{X^-X^-}$, this is the only state of this grouping.

2. We set $m_{X^-X^-} = 3$ and $m_{e^-X^-} = 1, 2, 3$.

For $m_{e^-X^-} = 1$ we obtain $L=1, 2, 3^2, 4^2, 5^3, 6^3, 7^3, 8^2, 9^2, 10$, and 11.

For $m_{e^-X^-} = 2$ we obtain $L=1, 2, 3, 4, 5$, and 6.

For $m_{e^-X^-} = 3$ we obtain $L=1$.

3. We set $m_{X^-X^-} = 3$, $m_{e^-X_2^-} = 1$, $m_{X^-X_2^-} = 3$, and $m_{e^-X^-} = 1, 2, 3$.

For $m_{e^-X^-} = 1$ we obtain $L=2, 3, 4^2, 5^2, 6^3, 7^2, 8^2, 9$, and 10.

For $m_{e^-X^-} = 2$ we obtain $L=2, 3, 4, 5$, and 6.

For $m_{e^-X^-} = 3$ we obtain $L=2$.

In groupings (ii) and (iii), the sets of multiplets obtained for higher values of $m_{e^-X^-}$

are subsets of the sets obtained for lower values, and we would expect them to form lower energy bands since they avoid additional small values of

$$\mathcal{R}_{e^-X^-}.$$

However, note that the (ii) and (iii) states predicted for

$$m_{e^-X^-} = 3$$

(at $L=1$ and 2, respectively) do not form separate bands in Fig. 14. This is because

$$V_{e^-X^-} = 3$$

increases more slowly than linearly as a function of $L(L+1)$ in the vicinity of

$$\mathcal{R}_{e^-X^-} = 3$$

(see Fig. 13). In such case the composite fermion picture fails [12,17].

We conclude that different kinds of long-lived Fermions (electrons and different charged excitonic complexes) formed in an electron-hole plasma in high magnetic fields can exhibit generalized incompressible fractional quantum Hall ground states with Laughlin-type correlations, and that these states can be described using a generalized composite fermion model.

6. Spatially separated electron-hole system

Even in very high magnetic fields (in the lowest Landau level), an asymmetry between e-e, h-h, and e-h interactions can be introduced by spatially separating 2D electron and hole layers. Such separation, which occurs for example in asymmetrically doped wide quantum wells, breaks the hidden symmetry and allows for a rich photoluminescence (PL) spectrum, which (unlike that for a co-planar system) can be therefore used as a probe of the low lying electron-hole states.

Let us consider an ideal system, in which electrons and holes occupy 2D parallel planes separated by a distance d . The interaction potentials are

$$V_{ee}(r) = V_{hh}(r) = \frac{1}{r}$$

and

$$V_{eh}(r) = -\frac{1}{\sqrt{r^2 + d^2}}.$$

The energy spectrum of a seven-electron-one-hole system is shown in Fig. 15 for $2S=15$ and values of d going from 0 to 5 (measured in units of the magnetic length λ).

For $d=5\lambda$, the e-h interaction is weak and, as a first approximation, we can say that the lowest band of states will consist of the lowest composite fermion band of the electron system plus the (constant) hole energy.

The allowed angular momenta will be given by L_e , the angular momenta of the low lying electron states, added to the hole angular momentum L_h of length $l_h=S=15/2$.

At $2S=15$, the composite fermion picture for the electrons gives

$$2S^* = 2S - 2p(N-1) = 15 - 2 \cdot (7-1) = 3.$$

The seven electrons fill the $l_0^*=3/2$ shell plus three of the QE states in the shell $l_{QE}=5/2$. The resulting electron angular momenta are $L_e=3/2, 5/2, \text{ and } 9/2$. This gives three bands of low lying states, with total angular momenta $6 \leq L \leq 9, 5 \leq L \leq 10, \text{ and } 3 \leq L \leq 12$, respectively. These three bands can be clearly distinguished at $d=5\lambda$ and the states within each band become nearly degenerate at $d \sim 10\lambda$.

For $d=0$, it is more useful to consider bound excitonic complexes (X and X^-) and Laughlin quasiparticles of the e^-X^- fluid.

First consider the multiplicative state with a single X and six electrons. At $2S=15$ six electrons have the Laughlin $\nu=1/3$ ground state since

$$2S^* = 2S - 2p(N-1) = 15 - 2 \cdot (6-1) = 5$$

gives a composite fermion shell which accommodates all six composite fermions. This is the lowest state at $L=0$, marked with a circle in frame (a). For a charge configuration containing one trion and five electrons, we can use the generalized composite fermion model with

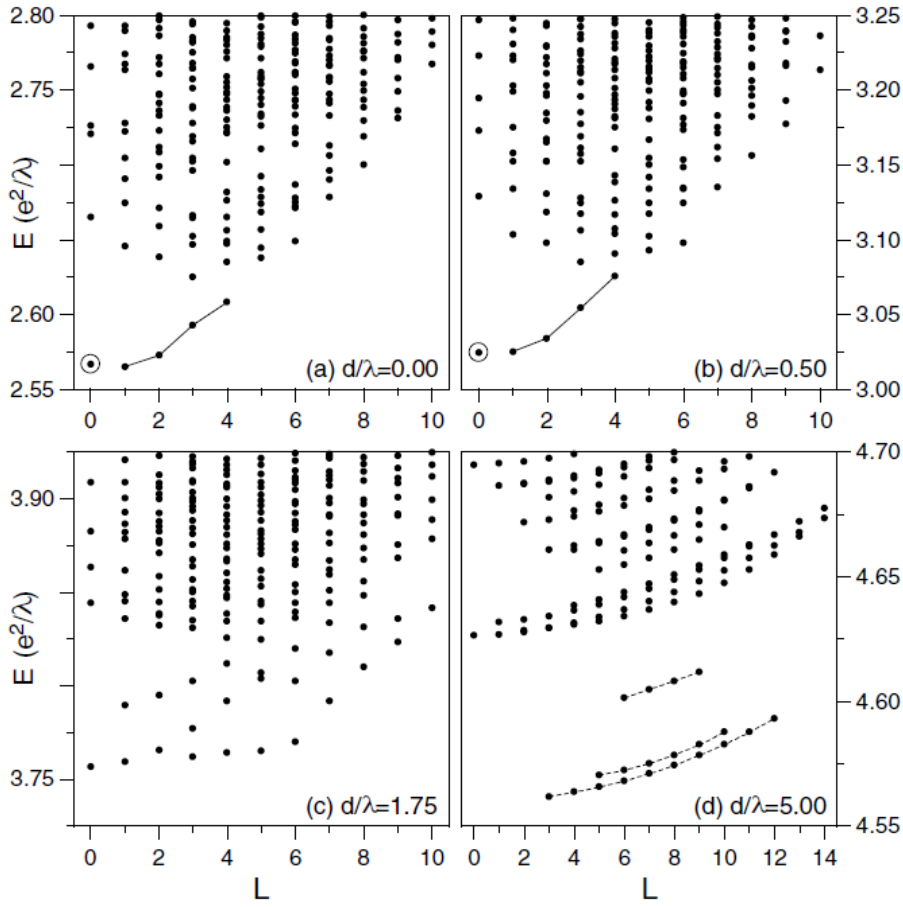


Fig. 15 The energy spectrum of a seven-electron-one-hole system in the lowest Landau level calculated on a Haldane sphere at $2S=15$, for different values of the separation d between electron and hole planes. In frames (a) and (b), the circle marks a multiplicative state and solid lines mark states containing a charged exciton X . In frame (d), the dashed lines mark three lowest bands.

$$m_{e^-X^-} = 3$$

$$m_{e^-X^-} = 2$$

This gives

$$2S_{e^-}^* = 2S - (m_{e^-e^-} - 1)(N_{e^-} - 1) - m_{e^-X^-}N_{X^-} = 15 - (3 - 1)(5 - 1) - 2 \cdot 1 = 5$$

and

$$2S_{X^-}^* = 2S - m_{e^-X^-}N_{e^-} = 15 - 2 \cdot 5 = 5$$

and the angular momenta

$$l_{e^-}^* = S_{e^-}^* = \frac{5}{2}$$

and

$$l_{X^-}^* = S_{X^-}^* - 1 = \frac{3}{2}.$$

There is one empty state in the lowest composite-fermion-electron shell giving $L_e=5/2$, and the composite-fermion-trion has $L_{X^-}=3/2$. Adding these two angular momenta gives $L=1, 2, 3$, and 4 as the lowest band of $5e^-X^-$ states.

The multiplicative state at $L=0$ (open circle) and the band of four multiplets containing a trion at $L=1$ to 4 (line) can clearly be seen at $d=0$ in frame (a).

Although the hidden symmetry is only approximate at $d>0$, these bands can be easily identified at $d=0.5\lambda$ in frame (b).

At an intermediate separation of $d=1.75\lambda$ in frame (c), neither description used for $d<\lambda$ nor $d>>\lambda$ is valid, and it seems that a low energy band occurs at $L=0, 1, 2, 3^2, 4, 5$, and 6.

Most likely, the trion unbinds but the hole is still able to bind one electron, forming an exciton with a significant electric dipole moment. This dipole moment results in repulsion between the exciton the remaining six electrons, so that the correlations are quite different than at $d=0$, where the exciton decouples.

The PL spectrum can be evaluated from the eigenfunctions obtained in the numerical diagonalization of finite systems. For $d>>0$, between one and three peaks are observed in the PL spectrum [27]. Their separations are related to the Laughlin gap (for creation of a QE-QH pair) and to the energy of interaction between the valence band hole and the electron system.

7. Charged excitons in a finite magnetic field

One final point is worth mentioning. The numerical calculations described so far were performed for an idealized model in which electrons and holes were confined to infinitely thin 2D layers, and only the lowest Landau level was considered. For realistic systems, effects due to spin, finite width of the quantum well, and Landau level mixing are very important. The energy spectra of the simple 2e-1h system calculated at $2S=20$ for parameters appropriate to a 11.5nm GaAs/AlGaAs quantum well are shown in Fig. 16.

Two frames correspond to the magnetic field of $B=13T$ and $68T$. We used five electron and hole Landau levels ($n\leq 4$) in the calculation, with the realistic magnetic field dependence of the hole cyclotron mass and the appropriate Zeeman splittings. The interaction matrix elements included finite (and different) effective widths of electron and hole quasi-2D layers.

There are a number of bound trion states in both frames, in contrast to only one singlet bound state at $B=0$ or only one triplet bound state predicted for an idealized system at infinite B . Three of these bound states are of particular importance.

$$X_s^-$$

and

$$X_{tb}^-$$

(b for “bright”), the lowest singlet and triplet states at $L=S$, are the only well bound radiative states, while X_{td}^- (d for “dark”) has by far the lowest energy of all non-radiative ($L\neq S$) states.

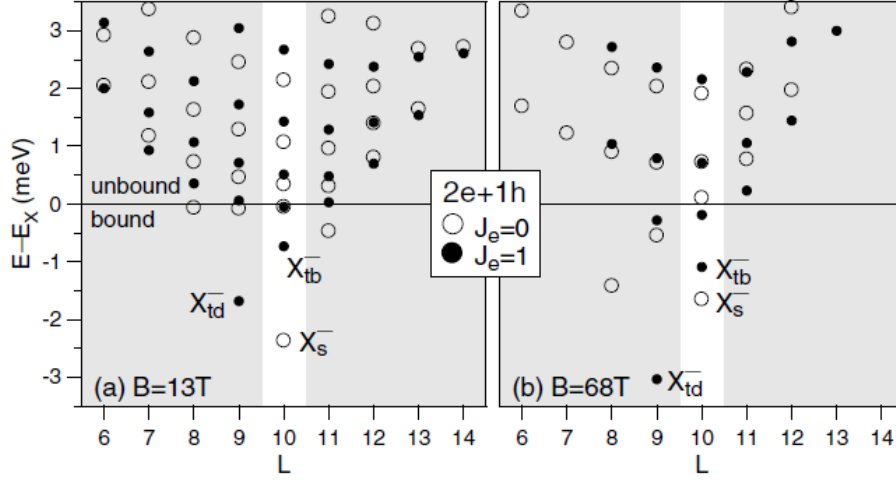


Fig. 16 The energy spectra (binding energy vs. angular momentum) of the two-electron-one-hole system calculated on a Haldane sphere with the Landau level degeneracy of $2S+1=21$. Five electron and hole Landau levels are included, and the parameters are appropriate for the 11.5nm GaAs quantum well in the magnetic field of $B=13T$ (a) and $68T$ (b).

The dark triplet trion state is the state discussed in the preceding sections; it is the only bound state in the lowest Landau level, but unbinds at low magnetic fields. The bright singlet trion state is the only bound state at $B=0$, but unbinds at very high fields due to the hidden symmetry. These states cross at $B\approx 30T$, as predicted in an earlier calculation [28].

The bright triplet trion has been discovered more recently [29]. It occurs only at intermediate fields and crosses neither the singlet trion nor the dark triplet trion. It also has a larger PL intensity than the singlet trion state.

Although an isolated dark triplet trion is non-radiative because of the angular momentum selection rule, its collisions with other trions or with electrons (which break the translational symmetry) could be expected to allow for its recombination.

However, the Laughlin correlations limit high energy collisions at low filling density ($\nu\approx 1/5$ or less) and the PL intensity of a dark triplet trion remains very low also in a presence of other particles [29]. In consequence, the dark triplet dark triplet trion state is not seen in PL, and there is no contradiction between experiment [30], which sees recombination of a triplet state at the energy above the singlet state up to 50T, and theory [29], which predicts that the lowest triplet state crosses the singlet at roughly 30T.

Summary

We have introduced the Jain composite fermion mean field picture and shown how the low lying states can be understood by simple addition of angular momentum. The mean field composite fermion picture gives the correct spectral structure not because of some cancellation between Chern-Simons and Coulomb interactions beyond the mean field level, but because it selects a low angular momentum subset of the allowed multiplets that avoids the largest pair repulsion. The Laughlin correlations, which describe incompressible quantum fluid states, depend critically on the electron pseudopotential being of “short range” (by which we mean that $V(L_{12})$ increases more quickly than $L_{12}(L_{12}+1)$).

The validity of Jain's picture also depends upon $V(L_{12})$ being of short range. Pseudopotentials describing quasiparticles of a Laughlin condensed state display short range behavior only at certain values of L_{12} . We have used this fact to explain why only certain states in the composite fermion hierarchy give rise to incompressible states of the quasiparticle fluid (or daughter states in the hierarchy). The pseudopotentials $V_n(L_{12})$ for higher Landau levels ($n > 0$) do not display short range behavior at all values of L_{12} , implying that Laughlin-like correlations will not necessarily result at $\nu' = 2p + \nu$, where p is an integer and ν is a Laughlin-Jain filling factor.

The composite fermion ideas have been applied successfully to multicomponent plasmas containing different types of Fermions with the prediction of possible incompressible fluid states for these systems. Finally, the energy spectrum and PL of electron-hole systems can be interpreted in terms of composite fermions and Laughlin correlations.

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