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Classification of Abelian and non-Abelian multilayer fractional quantum Hall states through the pattern of zeros

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A large class of fractional quantum Hall (FQH) states can be classified according to their pattern of zeros, which describes the way ideal ground-state wave functions go to zero as various clusters of electrons are brought together. In this paper, we generalize this approach to classify multilayer FQH states. Such a classification leads to the construction of a class of non-Abelian multilayer FQH states that are closely related to \( \hat{g}_k \) parafermion conformal field theories, where \( \hat{g}_k \) is an affine simple Lie algebra. We discuss the possibility of some of the simplest of these non-Abelian states occurring in experiments on bilayer FQH systems at \( \nu =2/3, 4/5, 4/7, \text{ etc.} \)

I. INTRODUCTION

One of the most important questions in condensed-matter physics relates to how we should characterize and classify the many different quantum phases of matter. A large part of the story is the theory of symmetry breaking, which tells us that we should classify various phases based on the symmetries of the ground-state wave function. Yet with the discovery of the fractional quantum Hall (FQH) states\(^1\)\(^2\) came also the understanding that there are many distinct and fascinating quantum phases of matter, called topological phases, whose characterization has nothing at all to do with symmetry. How should we characterize and systematically classify the different possible topological phases that may occur in a FQH system?

Let us run through a few obvious possibilities. We know that the FQH states contain topology-dependent and topologically stable degenerate ground states, which allow us to introduce the concept of topological order in FQH states.\(^3\)\(^4\) Such topology-dependent degenerate ground states suggest that the low-energy theories describing the FQH states are topological gauge-field theories in 2 + 1 dimensions, called Chern-Simons theories.\(^5\)\(^-\)\(^8\) So one possibility is that we may try to classify the different FQH phases by classifying all of the different possible Chern-Simons theories. But this is not a satisfactory approach for non-Abelian FQH states because we do not have a good way of knowing which Chern-Simons theories can possibly correspond to a physical system made of electrons and which cannot.

Another possibility comes through the connection between FQH wave functions and conformal field theory (CFT). It was discovered around 1990 that correlation functions in certain two-dimensional conformal field theories may serve as good model wave functions for FQH states.\(^9\) Thus perhaps we may classify FQH states by classifying all of the different conformal field theories that may be used to construct a “valid” FQH wave function. While the connection between FQH wave functions and CFT correlators has been extremely fruitful in both constructing new model wave functions and understanding their topological properties, the CFT approach is incomplete in that there does not exist a complete classification of all possible conformal field theories that may be used to construct FQH wave functions. This precludes the possibility of systematically classifying FQH states. The CFT approach also obscures the essential physics and the essential properties of the conformal field theory that allow one to obtain amenable FQH wave functions.

In an attempt to obtain a systematic classification of FQH states, it was shown recently that a wide class of single-component FQH states and their topological excitations can be classified by their pattern of zeros, which describe the way ideal FQH wave functions go to zero when various clusters of particles are brought together.\(^10\)\(^-\)\(^12\) This analysis led to the discovery of some non-Abelian FQH states whose corresponding CFT has not yet been identified. It also helped elucidate the role of CFT in constructing FQH wave functions. The CFT encodes the way the wave function goes to zero as various clusters of electrons are brought together. The order of these zeros must satisfy certain conditions and the solutions to these conditions correspond to particular CFTs. Thus in classifying and characterizing FQH states, one can bypass the CFT altogether and proceed directly to classifying the different allowed pattern of zeros and subsequently obtaining the topological properties of the quasiparticles from the pattern of zeros. This construction can then even be thought of as a classification of the allowed CFTs that can be used to construct FQH states. Furthermore, these considerations give way to a natural notion of which pattern-of-zeros solutions are simpler than other ones. In this sense, then, one can see that the Moore-Read Pfaffian quantum Hall state\(^9\) is the “simplest” non-Abelian generalization of the Laughlin state.

In this paper, we generalize the pattern-of-zeros classification to multilayer FQH wave functions. Such a systematic classification leads us to the construction of non-Abelian multilayer FQH states and also helps identify the simplest non-Abelian generalizations of the Halperin\(^3\) bilayer FQH states. Therefore, in the same way that the Pfaffian FQH state is the simplest non-Abelian single-layer state and so one of the first non-Abelian states expected to be realized experimentally, our construction allows us to identify the simplest non-Abelian bilayer states and therefore some of the ones that deserve further consideration in future experimental and numerical work.
We would like to point out that the “one-dimensional (1D) charge-density-wave” characterization of single-component FQH states is closely related to the pattern-of-zeros approach. Our pattern-of-zeros approach for multilayer FQH states can also be viewed as a way to generalize the 1D charge-density-wave characterization to multicomponent cases.

This paper is organized as follows. We begin by describing the ideal FQH wave functions that we can characterize by the pattern of zeros and their relation to symmetric holomorphic polynomials. In Sec. IV, we describe the different ways to characterize the pattern of zeros. In Sec. V, we find the conditions that the pattern of zeros must satisfy in order to describe valid FQH wave functions. In Sec. VI, we sketch how one may begin to construct ideal Hamiltonians whose ground states will be FQH wave functions with a given pattern of zeros. After a brief summary of the pattern-of-zeros data and conditions in Sec. VII, we explain in Sec. VIII the relation between the pattern-of-zeros approach and the CFT approach to FQH wave functions. In Sec. IX, we describe some example solutions of this systematic classification of multilayer FQH wave functions, which yields many non-Abelian multilayer states. In Sect. X, we discuss some of the simplest of these non-Abelian bilayer states that may be relevant for experiments on two-component quantum Hall systems and that warrant further numerical study.

II. FRACTIONAL QUANTUM HALL STATES AND SYMMETRIC POLYNOMIALS

The ground-state wave function of a two-dimensional system of electrons in the lowest Landau level can be written in the form

\[ \Phi(z_1, \ldots, z_N) = \prod_{i \neq j} (z_i - z_j) \Phi(z_i) \Phi(z_j), \]

where \( z_i = x_i + iy_i \) are the coordinates of the \( i \)th electron, and \( \Phi(z_1, \ldots, z_N) \) is a holomorphic function of \( z_i \). Since the electrons obey Fermi statistics, \( \Phi \) is antisymmetric under interchange of any two coordinates \( z_i \) and \( z_j \) when all of the electrons are identical. In many physical situations, the electrons may be distinguished by various quantum numbers, such as a spin index (when the Zeeman energy is not too high), a layer index (in a multilayer two-dimensional electron system), a valley index (such as in graphene and SiGe heterostructures), etc. In such cases, the ground-state wave function in the lowest Landau level may instead be written in the form

\[ \Psi = \Phi(z_i) e^{-\sum_{i<j} |z_i - z_j|^2 / 4}, \]

where \( I = 1, \ldots, N_I \) is a flavor index and \( N_I \) is the number of different flavors. \( \Phi \) is then necessarily antisymmetric only under interchange of \( z_i \) and \( z_j \) for any \( i \) and \( j \). Given any antisymmetric polynomial \( \Phi_{\text{antisy}}(z_i) \), we can uniquely construct a symmetric polynomial,

\[ \Phi_{\text{symm}}(z_i) = \prod_{I<j} (z_i - z_j)^{\delta_{ij}}. \]

\[ \Phi_{\text{symm}} \] will also be a polynomial because \( \Phi_{\text{antisy}} \) must vanish when any two identical particles approach each other. Thus the above division by the factor \( \prod_{I<j} (z_i - z_j) \) will never produce any poles in the resulting function.

Therefore to classify FQH phases of electrons, we can restrict our attention mainly to symmetric polynomials \( \Phi(z_i^I) \), where \( \Phi \) is invariant under the interchange of \( z_i^I \) and \( z_j^I \) for any \( i \) and \( j \) but not necessarily invariant under the interchange of \( z_i^I \) and \( z_j^I \) if \( I \neq J \). In this paper, we will often refer to \( I \) as a layer index. In the following, \( \Phi \) will always refer to such a symmetric multilayer polynomial.

We will introduce data, such as \( n, m \), and \( S_a \), to characterize and classify bosonic FQH states (i.e., symmetric polynomials) \( \Phi(z_i^I) \). From the above discussion, we see that the same set of data also characterizes fermionic FQH states whose wave functions are given by \( \Phi_{\text{antisy}}(z_i^I) = \Phi((z_i^I)) \prod_{I<j} (z_i^I - z_j^I). \)

III. IDEAL HAMILTONIANS AND IDEAL WAVE FUNCTIONS

Progress in understanding the various possible topological phases in the FQH effect has occurred largely because of the discovery of various kinds of model wave functions and the ideal Hamiltonians that yield these wave functions as their ground states. In the pattern-of-zeros approach, we classify all of the possible ideal wave functions for FQH phases. In this section we will explain what we mean by “ideal.”

For each topological phase in the FQH system, we want to have a certain representative wave function that captures the topological properties of the phase. The prototypical example of such an ideal wave function is the Laughlin wave function at filling fraction \( \nu = 1/m \),

\[ \Phi_{1/m}(z_i) = \prod_{i < j} (z_i - z_j)^m. \]

At every particle in \( \Phi_{1/m} \), there are \( m \) zeros, and there are no off-particle zeros. The exact ground state of the microscopic Hamiltonian with the Coulomb interaction does not have this simple property. Nevertheless, the ideal Laughlin wave function captures all of the essential topological properties of these phases.

The Laughlin wave functions are also ground states of certain ideal Hamiltonians that contain interaction potentials that are constructed only from \( \delta \) functions and their derivatives. For example, if the interaction potential is of the form

\[ V_i(z_1, z_2) = \delta(z_1 - z_2), \]

then the following wave function is an exact ground state, with vanishing total potential energy:

\[ \Phi_{1/2} = \prod_{i<j} (z_i - z_j)^2. \]

On the other hand, if the interaction potential between two electrons is of the form
\begin{equation}
V_2(z_1,z_2) = v_0 \delta(z_1 - z_2) + v_2 z_1^2 \delta(z_1 - z_2) \partial_{z_1}^2 \partial_{z_1}^2 (7)
\end{equation}

with \( v_0 > 0 \) and \( v_2 > 0 \), then the zero-energy ground state will be

\begin{equation}
\Phi_{1/4} = \prod_{i \neq j} (z_i - z_j)^4. \tag{8}
\end{equation}

More complicated ground states can be obtained through more complicated interactions. For example, consider the following three-body interaction between electrons:

\begin{equation}
V_{PF} = S(v_0 \delta(z_1 - z_2) \delta(z_2 - z_3) - v_1 \delta(z_1 - z_2) \delta(z_2 - z_3) \partial_{z_3} \partial_{z_3}). \tag{9}
\end{equation}

where \( S \) is the total symmetrization operator between \( z_1, z_2, \) and \( z_3 \). An exact zero-energy ground state of this interaction is the Pfaffian wave function at \( \nu = 1 \),

\begin{equation}
\Phi_{PF} = \mathcal{A}\left(\frac{1}{z_i - z_j}\right) \prod_{i < j} (z_i - z_j), \tag{10}
\end{equation}

where \( \mathcal{A} \) is the total antisymmetrization operator between \( z_1, \ldots, z_N \).

Since the Pfaffian and Laughlin wave functions are exact zero-energy ground states of interaction potentials that are constructed only from \( \delta \) functions and their derivatives, these wave functions can be characterized completely by their pattern of zeros, i.e., by the order of the zeros in the wave function as different numbers of particles are taken together. These wave functions have certain ideal properties; for example, the zeros that are bound to a particle lie precisely at the location of that particle, not slightly away from it. In this paper, we will classify ideal FQH wave functions, which are wave functions that are exact zero-energy ground states of such ideal Hamiltonians.

The existence of these ideal Hamiltonians is crucial. We can write down any arbitrary complex function of \( N \) variables but we cannot know whether it corresponds to a valid topological phase of matter in the limit \( N \to \infty \) unless we also know that it is the ground state of some local, gapped Hamiltonian. If this is the case, then it is conceivable that there exists some physical situation in which the low-energy effective interactions between electrons yields a phase of matter that is in the same universality class as the ideal wave functions that we constructed. Unfortunately, judging whether a many-particle interacting Hamiltonian is gapped in the thermodynamic limit is in general an intractable problem. So here we limit ourselves to classifying ideal wave functions, which at least we believe can be realized as ground states of local Hamiltonians. Whether the corresponding Hamiltonian can be gapped in the thermodynamic limit is a question that we must attempt to answer along the way.

The ideal quasihole wave functions for these states are also zero-energy eigenstates of their corresponding ideal Hamiltonians. The true ground state is distinguished by being the unique translationally invariant state with lowest angular momentum. Thus we can also classify topologically distinct quasiholes by their pattern of zeros, i.e., by the order of the zeros in the quasihole wave function as different numbers of electrons are taken to the quasihole. While the quasihole wave functions are also zero-energy states, the quasiparticles will be gapped and will have the same topological properties as the quasiholes. Therefore in the following, when we discuss the pattern of zeros of the quasiparticles, we are referring to the pattern of zeros of the quasihole wave functions.

Not all FQH phases have such ideal wave functions. The hierarchy states and the composite fermion states, for example, do not have ideal wave functions. While these phases do have their own model wave functions, they are not ideal in the sense that they cannot be directly described by their pattern of zeros or written as a correlation function of conformal primary operators in a CFT. Therefore, the pattern-of-zeros construction does not directly classify these phases. We will however discuss how they are related to the ideal multilayer FQH wave functions in Sec. IX A.

IV. PATTERN-OF-ZEROS CHARACTERIZATION

The spirit of the pattern-of-zeros approach is to consider bringing together \( a_I \) particles of type \( I \), for \( I = 1, \ldots, N_f \), and asking how \( \Phi \) goes to zero under such a procedure. The order of the zero will be denoted as \( S_\vec{a} \), where \( \vec{a} = (a_1, \ldots, a_{N_f}) \). In the following we will more precisely define \( S_\vec{a} \) and discuss some different yet equivalent ways of characterizing the pattern of zeros. This discussion is a straightforward generalization of the discussion in the single-layer case.\(^{10}\)

A. \( S_\vec{a} \) characterization

Consider a set of \( a_I \) coordinates of each type \( I \), and set \( \vec{a} = (a_1, \ldots, a_{N_f}) \). Define \( S_\vec{a} \) as the minimal power of \((\Pi_{I=1}^{N_f} P^a_{\vec{a}_I})\) in the polynomial \( \Phi \). This means that if we set

\begin{equation}
\xi^{(a)}_i = \lambda \xi^{(a)} + z^{(d)} \quad i = 1, \ldots, a_I, \quad \forall I,
\end{equation}

and we take \( \lambda \to 0 \), then

\begin{equation}
\Phi \sim \lambda^{S_\vec{a}} P(\{\xi^{(a)}_i\}, \{z^{(d)}\}, \{\xi^{(a)}_{a_I+1}, \ldots\}) + O(\lambda^{S_\vec{a}+1}), \tag{12}
\end{equation}

where \( P(\{\xi^{(a)}_i\}, \{z^{(d)}\}, \{\xi^{(a)}_{a_I+1}, \ldots\}) \) is a polynomial in \( \{\xi^{(a)}_i\} \) and the remaining coordinates \( \{z^{(d)}\} \) and \( \{\xi^{(a)}_{a_I+1}, \ldots\} \). We refer to \( z^{(d)} \) as the coordinate of an \( \vec{a} \) cluster. We assume that \( S_\vec{a} \) is independent of the choice of \( z^{(d)} \), which must be the case for translationally invariant wave functions. We also assume that \( S_\vec{a} \) is independent of the choice of \( \{\xi^{(a)}_i\} \) and that different polynomials \( P(\{\xi^{(a)}_i\}, \{z^{(d)}\}, \{\xi^{(a)}_{a_I+1}, \ldots\}) \) obtained from different choices of \( \xi^{(a)}_i \) are linearly dependent. This is the assumption of unique fusion.

We can immediately deduce some basic properties of \( S_\vec{a} \). Since \( \Phi \) has no poles, it is clear that \( S_{\vec{a}} \geq 0 \). Since \( \Phi \) must be
single valued under rotating $\lambda$ in the complex plane by an angle $2\pi$, $S_\lambda$ must be an integer. Let $S_{\alpha^I}$ be the minimal power of $\zeta_I$; that is, $(\zeta_I)^{S_{\alpha^I}} = \delta_{I\bar{I}}$. A translationally invariant $\Phi$ will have $S_{\alpha^I} = 0$, otherwise it will vanish everywhere.

Thus, for a translationally invariant polynomial, $S_\alpha$ is a nonnegative integer that characterizes the order of zero that results when the size of an $\alpha$ cluster goes to zero.

B. Derived polynomials and the $D_{\alpha\bar{b}}$ characterization

In the previous section, we introduced the derived polynomials $P((\zeta^I_1, \zeta^I_2, \ldots), (\zeta^I_{a+1}, \ldots))$. As a consequence of the unique fusion condition, these polynomials are actually independent of $(\zeta^I_1)$. We may consider more general derived polynomials by bringing together other sets of coordinates in $P$ to obtain $\tilde{P}(\zeta^I_{\alpha\bar{b}}, \zeta^I_{\beta}, \ldots)$. Then we may consider bringing together an $\alpha$ cluster and a $\bar{b}$ cluster,

$$\tilde{P}(\zeta^I_{\alpha\bar{b}}, \zeta^I_{\beta}, \ldots) \rightarrow (\zeta^I_{\alpha\bar{b}} - \zeta^I_{\beta}) P_{\alpha\bar{b}} \tilde{P}(\zeta^I_{\alpha\bar{b}}, \zeta^I_{\beta}, \ldots) + O((\zeta^I_{\alpha\bar{b}} - \zeta^I_{\beta}) P_{\alpha\bar{b}+1}).$$

(13)

Thus, $D_{\alpha\bar{b}}$ characterizes the order of the zeros in the derived polynomials as a cluster of $\alpha$ electrons are brought close to a cluster of $\bar{b}$ electrons. The unique-fusion condition assumes that the derived polynomials obtained from different ways of fusion are always linearly dependent.

The fact that $\Phi$ is a single-valued, symmetric polynomial implies

$$D_{\alpha\bar{b}} = D_{\bar{b}\alpha} \in \mathbb{Z}, \quad D_{\alpha\bar{b}} = \text{even}, \quad D_{\alpha\bar{b}} \geq 0.$$  

(14)

We can deduce a relation between $D_{\alpha\bar{b}}$ and $S_\alpha$ as follows. The order of the zero obtained by creating an $(\alpha + \bar{b})$ cluster is $S_{\alpha + \bar{b}}$. One way of creating such a cluster is by first creating an $\alpha$ cluster, then creating a $\bar{b}$ cluster, and finally bringing together the two clusters to create an $(\alpha + \bar{b})$ cluster. The order of zero in this case will be $S_\alpha + S_{\bar{b}} + D_{\alpha\bar{b}}$. Thus $D_{\alpha\bar{b}}$ can be obtained from $\{S_\alpha\}$ through the formula

$$D_{\alpha\bar{b}} = S_{\alpha + \bar{b}} - S_\alpha - S_{\bar{b}}.$$  

(15)

Since $S_{\alpha} = 0$, where recall $(\zeta_I)_{\alpha} = \delta_{I\alpha}$ is the unit vector in the $I$ direction, we also have

$$S_{\alpha + \bar{e}I} = S_\alpha + D_{\alpha, \bar{e}I}.$$  

(16)

From this recursion relation and from the fact that $S_{\bar{e}} = 0$, we may obtain $S_\alpha$ from the sequence $D_{\alpha\bar{b}}$. Therefore we may equivalently label the pattern-of-zeros data using $\{S_\alpha\}$ or $\{D_{\alpha\bar{b}}\}$.

C. Characterization by sequence of highest occupied orbitals

The integer $S_\alpha$ has the following meaning. A polynomial with $a_I$ particles of $I$th kind (i.e., $a_I$ particles in the $I$th layer) has a total order of $S_\alpha$. In other words the total angular momentum of the quantum Hall droplet is $S_\alpha$ if the droplet has $a_I$ particles in the $I$th layer. If we remove one particle from the $I$th layer, then the total angular momentum of the quantum Hall droplet will be reduced to $S_\alpha - a_I$. Thus we can interpret

$$I^\prime = S_\alpha - S_{\alpha - e_I}$$

(17)

as the angular momentum of the highest occupied orbital (HOO) in the $I$th layer for a quantum Hall droplet with $a_I$ particles in the $I$th layer. The $N_I$-dimensional sequence of vectors $I^\prime = (I^\prime_1, \ldots, I^\prime_{N_I})$ will be called the sequence of HOOs.

We see that $I^\prime$ makes sense only when $a_I > 0$. We will set $I^\prime = 0$ when it does not make sense. From Eq. (17), we also see that there is one-to-one correspondence between the sequence $S_\alpha$ and $I^\prime$. Thus we can also use $I^\prime$ to characterize the pattern of zeros in the wave function.

D. Relation to angular momentum on the sphere

A FQH wave function $\Phi((\zeta^I_I))$ defined on a sphere forms a representation of $SU(2)$. In such a case, $z$ represents the stereographic projection onto the plane of a point on the sphere. A single particle in the lowest Landau level can fill any of the $N_{\Phi} + 1$ orbitals, so the representation of $SU(2)$ formed in this case is the one with angular momentum $J = N_{\Phi}/2$. The SU(2) Lie algebra is generated by

$$L^z = z \partial_z - J, \quad L^- = \partial_z, \quad L^+ = z^2 \partial_z + 2Jz.$$  

(18)

In the multilayer case, then, the angular momentum of a particle of type $I$ will be $J_I = N_{\Phi}^I / 2$, where $N_{\Phi}^I$ is the total number of flux quanta through the sphere seen by the particles in the $I$th layer. Note that here we allow the numbers of flux quanta in different layers to be different. The total angular momentum of an $\alpha$ cluster in the $z$ direction will be the eigenvalue of the operator

$$L_{\alpha}^z = \sum_I \left( \sum_{a_{I}} z_{a_{I}} \partial_{z_{a_{I}}} - J_{I} \right).$$  

(19)

The operator $\sum_{a_{I}} z_{a_{I}} \partial_{z_{a_{I}}}$ counts the total power of a polynomial. Since the minimum total power of $\Pi_{I}^I \zeta_{a_{I}} J_{I}$ is $S_\alpha$, the minimum total angular momentum of an $\alpha$ cluster is given by $S_\alpha - \sum_{a_{I}} J_{I}$. This means that the $\alpha$ cluster carries an angular momentum of

$$J_{\alpha} = \alpha \cdot \vec{J} - S_{\alpha} = \frac{1}{2} \alpha \cdot \vec{N}_{\Phi} - S_{\alpha},$$  

(20)

where $\vec{J} = (J_1, \ldots, J_{N_I})$ and $\vec{N}_{\Phi} = (N_{\Phi}^1, \ldots, N_{\Phi}^{N_I})$. We will use this relation later to construct ideal Hamiltonians and to place conditions on the pattern of zeros for when they can correspond to rotationally invariant wave functions on the sphere.

V. CONSISTENCY CONDITIONS

For the pattern of zeros to describe a valid FQH wave function, it must satisfy certain consistency conditions. We already encountered several such conditions above. For instance, we found that $S_\alpha$ is a nonnegative integer, $D_{\alpha\bar{b}} = S_{\alpha + \bar{b}} - S_\alpha - S_{\bar{b}} \geq 0$, and $D_{\alpha\bar{e}I} = S_{\alpha + \bar{e}I} - 2S_\alpha$ is even. In the following we
find additional conditions that the pattern of zeros must satisfy.

### A. Concave condition

One of the most important conditions on the wave function is simply the condition that the wave function has no poles. This condition is remarkably restrictive on the allowed pattern-of-zeros sequences. Consider a derived polynomial \( P(z^{(a)}, z^{(b)}, \ldots) \) and fix all coordinates but \( z^{(a)} \), thus viewing it as a complex function \( f(z^{(a)}) \). \( f(z^{(a)}) \) has zeros at isolated points but no poles anywhere. Some of the zeros are located at \( z^{(a)}, z^{(b)} \), etc. Those zeros are called on-particle zeros. The rest of the zeros are called off-particle zeros.

If we imagine taking \( z^{(a)} \) around \( z^{(a)} \) without enclosing any off-particle zeros, then \( f \) will pick up a phase \( 2\pi D_{\bar{a}b} \). Similarly, if we take \( z^{(a)} \) around \( z^{(c)} \) without enclosing any off-particle zeros, then \( f \) will pick up a phase \( 2\pi D_{\bar{a}c} \). Now consider taking \( z^{(a)} \) around a contour that encloses both \( z^{(a)} \) and \( z^{(c)} \) in the limit that \( z^{(a)} \to z^{(a)} \), the complex function \( f \) must change by a phase that is greater than or equal to \( 2\pi (D_{\bar{a}b} + D_{\bar{a}c}) \). The phase can never be less than this amount because that would require the existence of off-particle poles that get taken to \( z^{(c)} \) in order to diminish the strength of the on-particle zeros. By definition, the phase change in \( f \) under the above procedure is \( 2\pi D_{\bar{a}b+c} \). Therefore, the condition that the wave function have no poles leads directly to the following concavity condition on the integers \( D_{\bar{a}b+c} \):

\[
D_{\bar{a}b+c} \geq D_{\bar{a}b} + D_{\bar{a}c}. \tag{21}
\]

In cases where all of the zeros are located on the particles and there are no off-particle zeros, the above inequality is saturated. This occurs in the Laughlin states \( \Phi = \prod_{i<j} (z_i - z_j)^m \), and their multilayer Abelian generalizations, the Halperin states

\[
\Phi = \prod_{i<j} (z_i' - z_j')^{K_{ij}} \prod_{i<j,i} (z_i - z_j)^{K_{ij}}. \tag{22}
\]

In the following we will rewrite the concave condition as

\[
\Delta_3(\bar{a}, \bar{b}, \bar{c}) \geq 0, \tag{23}
\]

\[
\Delta_3(\bar{a}, \bar{b}, \bar{c}) = D_{\bar{a}b+c} - D_{\bar{a}b} + D_{\bar{a}c} = S_{\bar{a}b+c} - S_{\bar{a}b} - S_{\bar{a}c} - S_{\bar{b}c} + S_{\bar{a}b} + S_{\bar{a}c} + S_{\bar{b}c}. \tag{24}
\]

### B. Cluster condition

The cluster condition is a way to associate some kind of grading to the polynomials that is physically meaningful. Let \( \{ \bar{n}_i \} \) for \( i = 1, \ldots, N_L \) be a set of vectors that generate an \( N_L \)-dimensional lattice, where \( N_L \) is, as before, the number of flavors of particles (or the number of layers). The cluster condition states that the concave condition is saturated, i.e.,

if either \( \bar{a}, \bar{b}, \text{ or } \bar{c} \) lie on the lattice generated by \( \{ \bar{n}_i \} \). That is, if either \( \bar{a}, \bar{b}, \text{ or } \bar{c} \) can be written as a linear combination with integer coefficients of the vectors \( \{ \bar{n}_i \} \). This means that a derived polynomial containing a \( \bar{k} = \sum \bar{k}_p \bar{n}_p \) cluster is nonzero unless \( z^{(a)} \) coincides with the coordinates of another cluster; viewed as a function of the single variable \( z^{(a)} \), it has no off-particle zeros. A consequence of this is that if all of the particles are fused to form \( \bar{n} \) clusters, then the resulting derived polynomial has the Laughlin-Halperin form [see Eq. (22)] and there are no off-particle zeros.

The single-layer Read-Rezayi \( Z_n \) parafermion wave functions satisfy an \( n \)-cluster condition and they are exact ground states of Hamiltonians with \( n+1 \)-body interactions. For a fixed filling fraction, as \( n \) increases, the number of topologically distinct quasiparticles, the ground-state degeneracy on higher genus surfaces and the complexity of interactions necessary to realize the state all increase. This suggests that the energy gap typically decreases with increasing \( n \). Wave functions that do not obey a cluster condition can be thought of as having infinite \( n \) and are not expected to correspond to gapped phases. This intuition also comes from the CFT approach to FQH wave functions; infinite \( n \) corresponds to an irrational conformal field theory, which does not yield a finite number of quasiparticles and a finite ground-state degeneracy on the torus. In the multilayer case, we may use the volume of the unit cell spanned by \( \{ \bar{n}_i \} \) as one way to measure the complexity of a given FQH state.

The cluster condition is extremely powerful and simplifying because it allows us to determine the entire pattern-of-zeros sequence from knowledge of a “small” number of them. To see how this works, first observe using Eq. (25) that

\[
D_{\bar{a}b,c} = \sum D_{\bar{a}b,c} = \sum m_{ij} \bar{n}_i \tag{26}
\]

where we have defined the matrix \( m_{ij} = D_{\bar{a}b,c} \). So for any vector \( \bar{k} = \sum \bar{k}_i \bar{n}_i \), where \( k_i \) is an integer and \( \Sigma k_i \bar{n}_i = 0 \), we have

\[
D_{\bar{a}b,c} = \sum k_i m_{ij} \bar{n}_j. \tag{27}
\]

The above equations imply

\[
D_{\bar{a}b,c} = \sum n_{IA} D_{\bar{a}b,c} = \sum n_{IA} m_{IA} = (mn^T)_{II} = \sum n_{IA} D_{\bar{a}b,c} = \sum n_{IA} D_{\bar{a}b,c} = (mn^T)_{II}, \tag{28}
\]

where we have also defined the matrix \( n_{IA} = (\bar{n}_i)_J \).

In terms of the sequence \( \{ \bar{s}_{IJ} \} \), this implies that for \( \bar{k} = \sum \bar{k}_i \bar{n}_i \), where \( k_i \) is an integer,
Therefore, all of the integers \( S_\alpha \) are specified by the points \( \vec{a} \) within the unit cell spanned by \( \{ \vec{n}_j \} \).

In terms of the HOO sequence \( I_{\vec{a}} \), we have
\[
I_{\vec{a}+\vec{\epsilon}} = S_{\vec{a}+\vec{\epsilon}} - (S_\alpha - \vec{\epsilon}) = S_\alpha - \sum A k_A m_A I_{\vec{a}} + \sum A k_A m_A I_{\vec{a}}.
\]

Finally, note that since \( D_{\vec{a},\vec{\epsilon}} \) is an even integer, we have
\[
even = D_{\vec{a},\vec{\epsilon}} = \sum I n_I D_{\vec{a},\vec{\epsilon}} = \sum I n_I m_I = (nmT)_{IJ}.
\]  

\[ \text{(31)} \]

\section*{C. Equal area layers}

The density profiles of the single-particle states in the lowest Landau level, \( e^{-|z|^2/4l_B^2} \), are in the shape of a ring, with a peak at a radius \( r_m = \sqrt{2m I_B} \), where \( l_B \) is the magnetic length. Such a wave function has an angular momentum \( m \). When many of these orbitals are filled by particles, the total wave function will describe a uniform, rotationally symmetric state that goes to zero at a radius \( r_{max} = \sqrt{2m}\text{max} I_B \), where \( m_{\text{max}} \) refers to the filled orbital with maximum angular momentum. Therefore, a given quantum Hall wave function will describe a QH droplet of area \( 4\pi m_{\text{max}} I_B^2 \), and \( m_{\text{max}} \) is given by the maximum power of \( z_j \) (or \( z_i \) for any other fixed \( i \)).

An important constraint on the multilayer quantum Hall wave functions is that they must describe systems in which each layer occupies the same area, up to small corrections. The requirement that each layer occupies exactly the same area amounts to the requirement that each layer has exactly the same number of flux quanta, \( N_\Phi \).

An important constraint on the multilayer quantum Hall wave functions is that they must describe systems in which each layer occupies the same area, up to small corrections. The requirement that each layer occupies exactly the same area amounts to the requirement that each layer has exactly the same number of flux quanta, \( N_\Phi \). However, it is reasonable to include states in which different layers occupy equal areas only up to \( O(N_\Phi) \) corrections, where \( N_\Phi \) is the number of electrons.

Such a requirement of approximately equal area layers is summarized in the following equation:
\[
\lim_{N_\Phi \to \infty} N_\Phi = 1.
\]  

We wish to see how this condition translates into a condition on the pattern of zeros. The conditions are slightly different depending on whether we ultimately want to characterize gapped FQH phases of fermions or bosons. If we are interested in fermionic phases, we require that \( \Psi((z_I')) = \prod_{I,J} (z_I' - z_J') \Phi((z_I')) \) be a valid FQH wave function of fermions, which does not require that \( \Phi((z_I')) \) be a valid FQH wave function of bosons. In what follows we will explicitly analyze the bosonic case, where we require \( \Phi((z_I')) \) to be a valid multilayer FQH wave function of bosons.

\( N_\Phi \) is equal to the maximal power of \( z_I' \), for the boson wave function \( \Phi \), this is given by
\[
N_\Phi = S_\Phi - S_\Phi - \vec{\epsilon}_I.
\]  

where recall \( N_\Phi \) is the number of particles of type \( I \). Using the cluster condition, we find
\[
N_\Phi = \sum I (n^{-1} m)_{IJ} + S_\Phi - S_\Phi - \vec{\epsilon}_I = m_{IJ}.
\]  

This can be seen most easily by ignoring the \( O(N_\Phi) \) terms in Eq. (34), taking \( N_\Phi/N_\Phi \to 1 \), and inverting the result to obtain \( N_\Phi - N_\Phi \Sigma_{(m^{-1} n)_{IJ}} \), which must be nonnegative. From this analysis, we learn that if \( n^{-1} m \) is not invertible, then the pattern of zeros cannot fix the ratio of particles \( N_I/N_I \) in the different layers. Therefore the corresponding FQH state has a gapless mode corresponding to the relative density fluctuations between the different layers.

As a simple example of this analysis, consider the \( 1,1,1 \) Halperin bilayer state, which is known to have a gapless density mode and for which \( n^{-1} m = (1,1) \), which is not invertible. A macroscopic number of particles can freely go from one layer to the other without changing the area of the quantum Hall droplets, signaling the existence of a gapless relative density mode.

Inverting Eq. (34) yields
\[
\sum I (m^{-1} n)_{IJ} (N_\Phi - S_\Phi + S_\Phi - \vec{\epsilon}_I + m_{IJ}) = N_I.
\]  

from which we can read off the filling fraction in each layer,
\[
\nu_I = \sum I (m^{-1} n)_{IJ}.
\]  

The total filling fraction is the sum of the filling fraction of each layer: \( \nu = \sum \nu_I \).

For fermions, Eq. (33) is modified to
\[
N_\Phi = S_\Phi - S_\Phi - \vec{\epsilon}_I + N_I - 1,
\]  

due to the extra factor \( \prod_{I,J} (z_I' - z_J') \) in \( \Psi((z_I')) \). Note that \( \{ S_I \} \) still describes the pattern of zeros of the symmetric polynomial \( \Phi \). The result for fermions is therefore
\[
\nu_I = \sum I (m^{-1} n)_{IJ} \geq 0,
\]  

where \( I \) is the \( N_I \times N_I \) identity matrix. If \( (1+m^{-1} n) \) is not invertible in the fermionic case, then there are gapless relative density modes, which is why the filling fraction in each layer becomes undefined.

\section*{D. Shift and rotational invariance on the sphere}

Consider a multilayer quantum state with \( N_I \) particles in the \( I \)th layer. We want to put the quantum state on a sphere with \( N_\Phi \) flux quanta in the \( I \)th layer. We would like to know for which set of \( N_\Phi \) can the quantum Hall state completely fill the sphere? Naively, one may expect \( N_\Phi \) and \( N_I \) are re-
lated by the filling fraction in each layer $N^j_i / \nu_j = N^j_i$. However, the precise relation between the number of flux quanta and the number of electrons includes a shift,

$$v^{-1} \sum_i v_j N^j_i = v^{-1} N_e - S,$$  \hspace{1cm} (40)

where $S$ is of order 1 in the large $N_e$ limit [see Eq. (36)].

More precisely, completely filling the sphere means that the quantum Hall state is rotationally invariant with zero total angular momentum. Using Eq. (20), we find that, for a bosonic FQH state characterized by $S_2$ with $N_1$ particles and $N^j_i$ flux quanta in the $i$th layer, the maximum total angular momentum is given by

$$J_\tilde{N} = \frac{1}{2} \tilde{N} \cdot \tilde{N}_\Phi - S_\tilde{N},$$  \hspace{1cm} (41)

where $\tilde{N} = (N_1, \ldots, N_N)$ and $\tilde{N}_\Phi = (N^1_\Phi, \ldots, N^N_\Phi)$. For a fermionic FQH state characterized by $S_2$ with $N_1$ particles and $N^j_i$ flux quanta in the $i$th layer, the maximum total angular momentum is given by

$$J_\tilde{N} = \frac{1}{2} \tilde{N} \cdot \tilde{N}_\Phi - S_\tilde{N} - \frac{1}{2} \sum_i N_j (N_j - 1).$$  \hspace{1cm} (42)

In the above, $\tilde{N}_\Phi$ must satisfy

$$\text{for bosons}$$

$$N^j_i \geq S_N - S_{\tilde{N} - e_i}$$

in order for the wave function to fit into each layer.

 Completely filling the sphere requires that

$$N^j_i = S_N - S_{\tilde{N} - e_i} + N_j - 1$$

for fermions. (43)

and $J_\tilde{N} = 0$. We see that $\tilde{N}$ and $\tilde{N}_\Phi$ must satisfy

$$\tilde{N} \cdot \tilde{N}_\Phi = \begin{cases} 2S_\tilde{N} & \text{for bosons} \\ 2S_\tilde{N} + \sum_i N_j (N_j - 1) & \text{for fermions} \end{cases}$$

which implies (for both bosons and fermions)

$$\sum_i N_j (S_N - S_{\tilde{N} - e_i}) = 2S_\tilde{N}.$$  \hspace{1cm} (45)

If a given $\tilde{N}$ does not satisfy Eq. (46), then the corresponding quantum Hall state (with $N_1$ particles on the $i$th layer) cannot completely fill the sphere. For $\tilde{N}$ that satisfies Eq. (46), the corresponding quantum Hall state can completely fill the sphere and has zero total angular momentum if $\tilde{N}_\Phi$ is given by Eq. (44). Equation (46) can generally be satisfied only if $\tilde{N}$ lies on the lattice spanned by $\{e_i\}$.

We would like to remark that it is easy to have different numbers of flux quanta on different layers in numerical calculations. The pattern of $(\tilde{N}, \tilde{N}_\Phi)$ where the quantum Hall state has zero total angular momentum on the sphere can be used as a fingerprint to identify different quantum Hall states through numerical calculations (for example, see Tables IV and V).

E. Additional constraints: $\Delta_3 = \text{even}$

The analysis of the single-layer case in Ref. 10 has suggested an additional condition,

$$\Delta_3(a,b,c) = \text{even}.$$  \hspace{1cm} (47)

There, it was found that allowing $\Delta_3(a,b,c) = \text{odd}$ allows for certain pattern-of-zeros sequences that either do not correspond to single-valued wave functions (such as the square root of the Pfaffian) or could not correspond to translationally invariant wave functions. It was suggested that one way to rule out such possibilities is to impose Eq. (47). How should this condition be generalized to the multilayer situation?

One natural generalization is to impose $\Delta_3(\tilde{a}, \tilde{b}, \tilde{c}) = \text{even}$ for all $\tilde{a}$, $\tilde{b}$, and $\tilde{c}$. However, we find that this condition is too restrictive. It rules out certain known FQH wave functions, such as the $su(3)/u(1)^2$ non-Abelian spin singlet states.19,20

The need to relax this condition while still having it remain compatible with the single-layer situation suggests that we should impose Eq. (47) only for choices of $\tilde{a}$, $\tilde{b}$, and $\tilde{c}$ that are collinear through the origin.

While it was found that allowing $\Delta_3(a,b,c) = \text{odd}$ allows for pattern-of-zeros sequences that do not seem to correspond to valid translationally invariant, single-valued wave functions, there are known cases of CFTs with $\Delta_3(a,b,c) = \text{odd}$ that do seem to yield translationally invariant, single-valued wave functions. One such example is the so-called Gaffnian wave function, which has $\Delta_3(1,1,1) = \text{odd}$ and which can be constructed using the minimal model CFT $\mathcal{M}(5,3)$.21 This CFT however is nonunitary. It has been suggested that FQH wave functions constructed using nonunitary CFTs correspond to gapless phases;22 whether this is always necessarily the case is currently an important open question in FQH theory.

These considerations suggest that in order to restrict ourselves to pattern-of-zeros sequences that have a corresponding unitary CFT, we should impose $\Delta_3(\tilde{a}, \tilde{b}, \tilde{c}) = \text{even}$ for those $\tilde{a}$, $\tilde{b}$, and $\tilde{c}$ that are collinear through the origin. In our search for pattern-of-zeros solutions, we will impose this condition and analyze the resulting states. The precise connection, if any, between this condition and valid FQH wave functions that correspond to unitary CFTs remains to be clarified.

VI. IDEAL HAMILTONIANS

Given a pattern-of-zeros sequence, it is important to be able to construct a local, gapped Hamiltonian whose ground-state wave function has the given pattern of zeros. If this is possible, then we know that the corresponding pattern-of-zeros sequence describes a topological phase of matter. Whether this particular phase is realized in an experiment then depends on the particular types of low-energy effective
interactions between the electrons in the unfilled Landau levels.

We can go about constructing such a Hamiltonian by noticing that on a sphere, the integers \( S_q \) are directly related to the angular momentum of the \( \tilde{a} \) cluster. For an electron system on a sphere with \( N_e \) flux quanta for the \( l \)th layer, an electron of type \( I \) will carry an angular momentum \( J' = N_e / 2 \). For an \( \tilde{a} \) cluster, the maximum angular momentum is therefore \( \tilde{a} \cdot J \). However, for a polynomial \( \Phi(\{ L_l \}) \) described by a pattern of zeros \( \{ S_q \} \), the maximum allowed angular momentum of the \( \tilde{a} \) cluster is only \( J_q = \tilde{a} \cdot J - S_q \). The pattern of zeros forbids the appearance of angular momentum \( \tilde{a} \cdot J - S_q + 1, \tilde{a} \cdot J - S_q + 2, \ldots, \tilde{a} \cdot J \) for any \( \tilde{a} \) clusters in \( \Phi(\{ L_l \}) \).

Such a condition can be easily enforced by writing the Hamiltonian as a sum of projection operators, \( P_S^{(a)} \). Let \( P_S^{(a)} \) be a projection operator that acts on the \( \tilde{a} \)-cluster Hilbert space. \( P_S^{(a)} \) projects onto the subspace of \( \tilde{a} \) clusters that have total angular momenta greater than \( \tilde{a} \cdot J - S \). Now consider the Hamiltonian

\[
H(\{ S_q \}) = \sum_{\tilde{a}} \sum_{\tilde{a} \text{ clusters}} P_S^{(a)},
\]

where \( \sum_{\tilde{a} \text{ clusters}} \) sums over all of the \( \tilde{a} \) clusters for a fixed \( \tilde{a} \). The wave function described by \( \{ S_q \} \) will clearly be a zero-energy ground state of the above \( H(\{ S_q \}) \). In many cases, there is only one unique ground-state wave function with minimal total angular momentum but in general there can be many independent polynomials with the same pattern of zeros. In such a situation, the Hamiltonian would need to be modified further to select for a particular polynomial with the given pattern of zeros.

In order for the above Hamiltonian to be local, \( \Sigma_q \) must be limited to a small, finite number of \( \tilde{a} \) clusters. But as a result, we cannot guarantee in general that the ground-state wave functions will all be described by the sequence \( \{ S_q \} \) for every \( \tilde{a} \), or even that they will obey the cluster condition. In many of the known cases, such as the Laughlin, Moore-Read, and Read-Rezayi wave functions, the sum over \( \tilde{a} \) clusters can indeed be terminated after the first few clusters while still yielding a unique zero-energy ground-state wave function with minimal angular momentum which is described by \( \{ S_q \} \).

The above construction for \( H(\{ S_q \}) \) should therefore be viewed as a starting point for constructing an ideal Hamiltonian that is local, gapped, and whose unique minimal angular momentum ground-state wave function is described by \( \{ S_q \} \). In some of the simplest cases, we know that this construction suffices.

**VII. SUMMARY: PATTERN-OF-ZEROS DATA AND CONDITIONS**

We have found that the polynomials \( \Phi(\{ L_l \}) \), \( l = 1, \ldots, N_f \) that may correspond to stable FQH states are described by the following data:

where \( n \) and \( m \) are \( N_f \times N_f \) matrices with integer entries that satisfy

\[
m_{ij} \geq 0, \quad n_{ij} \geq 0, \quad \det n \neq 0,
\]

\[
nm^T = nm^T, \quad (nm^T)_{ij} = \text{even}.
\]

The above implies that \( n^{-1}m \) is a symmetric matrix. Furthermore, for the pattern of zeros to fix the relative densities of particles in each layer, we have

\[
(n^{-1}m)_{ij} \quad \text{is invertible(for bosons)},
\]

\[
(n^{-1}m)_{ij} + \delta_{ij} \quad \text{is invertible(for fermions)}.
\]

Otherwise, there are gapless relative density fluctuations. We also have

\[
\nu_l \geq 0,
\]

\[
\nu_l = \begin{cases} 
\sum_j (n^{-1}m)_{lj} & \text{for bosons} \\
\sum_j (1 + n^{-1}m)_{lj} & \text{for fermions}.
\end{cases}
\]

The total filling fraction is \( \nu = \sum \nu_l \).

Note that the \( \{ S_q \} \) need to be specified only for values of \( \tilde{a} \) that are contained in the unit cell spanned by \( \{ \tilde{n}_l \} \), where \( \tilde{n}_l \) corresponds to the \( l \)th row of the matrix \( n : (\tilde{n}_l)_{ij} = n_{ij} \).

Using the cluster condition, we can determine \( S_q \) for all \( \tilde{a} \) simply from \( m \), the fact that the \( S_q \) is zero, and from the values of \( S_q \) for which \( \tilde{b} \) lies in the unit cell spanned by \( \{ \tilde{n}_l \} \),

\[
S_{\tilde{a} \tilde{b}} = S_{\tilde{a}} + \sum_l k_l S_{\tilde{n}_l} + \sum_{l \neq j} k_{lm} S_{\tilde{n}_l} + \frac{1}{2} \sum_{l \neq j} (nm^T)_{lj}(k_{lj} - \delta_{lj} k_{lj}),
\]

where \( \tilde{b} = \sum k_l \tilde{n}_l \), \( S_{\tilde{a}} \) must satisfy

\[
\Delta_2(\tilde{a}, \tilde{a}) = \text{even},
\]

\[
\Delta_2(\tilde{a}, \tilde{b}) \geq 0, \quad \Delta_3(\tilde{a}, \tilde{b}, \tilde{c}) \geq 0,
\]

where

\[
\Delta_2(\tilde{a}, \tilde{b}) = S_{\tilde{a} \tilde{b}} - S_{\tilde{a}} - S_{\tilde{b}},
\]

\[
\Delta_3(\tilde{a}, \tilde{b}, \tilde{c}) = S_{\tilde{a} \tilde{b} \tilde{c}} - S_{\tilde{a} \tilde{b}} - S_{\tilde{b} \tilde{c}} - S_{\tilde{a} \tilde{c}} + S_{\tilde{a}} + S_{\tilde{b}} + S_{\tilde{c}}.
\]

Finally we impose

\[
\Delta_3(\tilde{a}, \tilde{b}, \tilde{c}) = \text{even}
\]

for those \( \tilde{a}, \tilde{b}, \) and \( \tilde{c} \) that are collinear through the origin.

**Product of symmetric polynomials and primitive solutions**

Consider two polynomials of \( \tilde{a} \)-cluster form, \( \Phi \) and \( \Phi' \), and consider their product: \( \Phi = \Phi \Phi' \). The pattern of zeros of
Thus we see that the cluster condition implies that an 

tions.

primitive and nonprimitive solutions. Primitive solutions are 
of-zeros are linear. Thus the pattern of zeros of \( \Phi \) is also 

valid. However, notice that the condition for filling the 
sphere is not linear. There may be two FQH wave functions 
\( \Phi \) and \( \Phi' \) that can fill the sphere but whose product \( \Phi \) cannot 

fill the sphere.

Thus, we can divide the pattern-of-zeros solutions into 

primitive and nonprimitive solutions. Primitive solutions are 
solutions that cannot be written as a sum of two other solu-

tions.

VIII. RELATION TO CONFORMAL FIELD THEORY

The pattern-of-zeros approach is closely related to the 

conformal field theory approach to constructing FQH wave 

functions. In the CFT approach, the symmetric polynomial 

\( \Phi(z_{1}, \ldots, z_{N}) \) that describes a multilayer FQH state can be written 
as a correlation function of a set of electron operators \( V_{c}z_{i} \) in a 

CFT,23

\[
\Phi(z_{1}, \ldots, z_{N}) = \lim_{z_{a} \to \infty} z_{a}^{2\tilde{h}_{a}} \langle V(z_{a}) \rangle \prod_{i \neq a} V_{c}(z_{i}).
\]

(58)

The operators \( V_{c}z_{i} \) are written in the form

\[
V_{c}(z) = \psi_{c}^{\dagger}(z) e^{i \Sigma M_{ij} \Phi(z_{j})},
\]

where \( e^{i \Sigma M_{ij} \Phi(z_{j})} \) is a vertex operator in a \( U(1)^{N} \) CFT. It has 

scaling dimension \( \Sigma M_{ij} \). \( \psi_{c} \) is a simple current operator; 

that is, it satisfies the following fusion relation:

\[
\psi_{a} \psi_{b} = \psi_{a+b, \delta}.
\]

(60)

This Abelian fusion rule is the CFT version of the unique-

fusion condition. The cluster condition implies that \( \psi \) satis-

fies

\[
\psi_{\delta_{1}} \sim 1,
\]

(61)

where \( (\delta_{1}) = n_{1} \). An \( n_{1} \) cluster of electrons will be described 

by the operator

\[
V_{a} = \prod_{l} V_{c}(z_{l}) = \psi_{a} e^{i \Sigma M_{ij} \Phi(z_{j})}.
\]

(62)

Thus we see that the cluster condition implies that an \( \delta_{1} \) 

cluster is described by a vertex operator \( e^{i \Sigma \delta_{m} M_{ij} \Phi(z_{j})} \). If all 
of the particles are grouped into \( n_{1} \) clusters, then the cor-

responding derived polynomial will be just a correlation func-
tion of vertex operators in a \( U(1)^{N} \) theory, which will have 

no off-particle zeros and will be of the Laughlin-Halperin 

form.

Let us denote the scaling dimension of the operator \( V_{a} \) as

\[
h_{a} = h_{a}^{ae} + h_{a}^{a},
\]

(63)

where \( h_{a}^{ae} \) is the scaling dimension of the simple current \( \psi_{a} \) 

and \( h_{a}^{a} \) is the scaling dimension of the vertex operator. No-
tice that since \( \psi_{\delta_{1}} \sim 1 \), the simple current scaling dimensions 

satisfy \( h_{a}^{ae} = h_{a}^{ae} \). The scaling dimension of the Gaussian part 
is given by

\[
h_{a}^{a} = \frac{1}{2} a_{W} (M T M^{T})_{a}.
\]

(64)
The pattern of zeros is related to the scaling dimensions 

through the relation

\[
D_{a, \delta} = h_{a}^{a} - h_{a} = S_{\delta_{1}} - S_{a} - S_{\delta},
\]

(65)

This allows us to obtain the scaling dimensions from the 

pattern of zeros. Using the cluster condition, some algebra 

shows that \( M T = n^{-1} m \) and so Eq. (64) becomes

\[
h_{a}^{a} = \frac{a_{W} n^{-1} m d}{2}.
\]

(66)
The scaling dimensions of the simple-current part can also be 
determined from the pattern of zeros by using the fact that 

\( h_{\delta_{1}} = \frac{(n m)^{\delta_{1}}}{2} \), \( h_{0} = 0 \), and applying Eq. (65) iteratively. This yields

\[
\frac{(n m)^{\delta_{1}}}{2} = \sum_{a} n_{A} h_{\delta_{1}} + S_{\delta_{1}}.
\]

(67)

Multiplying both sides by \( n^{-1} \) gives

\[
h_{\delta_{1}} = \frac{(n m)^{\delta_{1}}}{2} - \frac{\sum_{a} n_{A} h_{\delta_{1}} + S_{\delta_{1}}}{2}.
\]

(68)

In a similar manner, one can obtain

\[
h_{a} = \sum_{l} a_{l} (n^{-1} m d_{l}) - \sum_{\delta_{1}} \frac{(n m)^{\delta_{1}}}{2} - \sum_{a} S_{a}.
\]

(69)

which determines \( h_{a}^{a} \) in terms of the pattern of zeros.

Note that the correlation function of the Gaussian part is, 

leaving the background charge implicit,

\[
\langle \prod_{l} e^{i \Sigma \delta_{m} M_{ij} \Phi(z_{j})} \rangle = \prod_{l} \prod_{l<j} (z_{l} - z_{j})^{(M T M^{T})_{l}} (z_{l} - z_{j})^{(M T M^{T})_{l}}
\]

\[
= \prod_{l} \prod_{l<j} (z_{l} - z_{j})^{(n^{-1} m d_{l})} \prod_{l} \prod_{l<j} (z_{l} - z_{j})^{(n^{-1} m_{l})}.
\]

(70)

Thus the FQH wave function is of the form

\[
\Phi(z_{1}, \ldots, z_{N}) = \Phi_{ae}(z_{1}) \Phi_{ga}(z_{1}),
\]

(71)

where

\[
\Phi_{ga}(z_{1}) = \prod_{l} \prod_{l<j} (z_{l} - z_{j})^{(n^{-1} m)} \prod_{l} \prod_{l<j} (z_{l} - z_{j})^{(n^{-1} m_{l})}.
\]

(72)

\( \Phi_{ae} \) arises from the correlation function of the simple current 

sector and is the “non-Abelian part” of the wave function.

In this manner, each pattern-of-zeros solution corresponds 

to the current algebra of a rational CFT. The connection be-
tween the pattern-of-zeros approach and the CFT approach can be thought of in the following way. The pattern of zeros describes the essential properties of the CFT that yield valid FQH wave functions. So in order to classify ideal FQH wave functions, one can bypass the CFT altogether and go directly to the heart of the matter: characterizing the allowed pattern-of-zeros solutions. Furthermore, since each pattern-of-zeros solution corresponds to a CFT, the pattern-of-zeros classification can be viewed as a classification of the allowed CFTs that can be used to construct FQH wave functions.

In this formulation, the pattern of zeros classifies all those ideal FQH wave functions that can be formulated as a correlation function of conformal primary fields. There are also many FQH wave functions, such as the hierarchy states and the Jain series, that cannot be written in this way. These wave functions are outside of the pattern-of-zeros classification. However, they may be closely related to the pattern-of-zeros construction. We comment on this connection elsewhere.

### A. Alternate labeling

Using Eq. (69), we can derive a formula for \( S_\tilde{a} \) in terms of \( h_{\tilde{a}}^\gamma \), \( n \), and \( m \).

\[
S_\tilde{a} = h_\tilde{a} - \sum_i a_i h_{\tilde{a}}^\gamma_i .
\]  

(73)

Thus there is a one-to-one correspondence between the simple-current scaling dimensions and the sequence \( \{ S_\tilde{a} \} \). This means that there is yet another way to label the pattern of zeros. Earlier, we found that one convenient labeling of the pattern of zeros is with an \( N_f \times N_f \) matrix \( n \), an \( N_f \times N_f \) matrix \( m \), and the value of the non-negative integers \( S_\tilde{a} \) for \( \tilde{a} \) lying inside the unit cell spanned by the rows of \( n \). An alternative, equivalent labeling of the pattern of zeros is by specifying the following data:

\[
n, \ m, \ \{ h_{\tilde{a}}^\gamma \}. \]

(74)

for \( \tilde{a} \) lying inside the unit cell spanned by \( n \). This labeling is convenient because it makes close contact with the corresponding CFT description: if \( h_{\tilde{a}}^\gamma \neq 0 \) for some \( \tilde{a} \), then the CFT has a nontrivial simple-current structure and therefore generally also has some form of non-Abelian statistics.

### B. Relevant CFTs for multilayer FQH states: \( \hat{g}_k/u(1)^r \) parafermions

In the single-layer case, many of the pattern-of-zeros solutions were found to be closely related to the \( Z_4 \) parafermion CFTs of Zamalodchikov and Fateev. What are the relevant CFTs to expect in the multilayer case? The answer is that some of the corresponding CFTs in the multilayer case will be closely related to the \( \hat{g}_k/u(1)^r \) parafermion CFTs that were constructed by Gepner, where \( \hat{g}_k \) is a simple affine Lie algebra at level \( k \) and \( r \) is the rank of the Lie algebra \( g \).

The case \( g = su(2) \) is equivalent to the \( Z_4 \) parafermion CFTs of Zamalodchikov and Fateev.

The simple-current algebra of the \( \hat{g}_k/u(1)^r \) parafermion CFT has the following structure. For every element \( \tilde{a} \) of the \( r \)-dimensional root lattice of \( g \), associate a simple current operator \( \psi_\tilde{a} \). The simple currents will have the fusion rules

\[
\psi_\tilde{a} \psi_\tilde{b} = \psi_{\tilde{a} \tilde{b}} .
\]  

Furthermore, \( \psi_\tilde{a} = \psi_{-\tilde{a}} \) if \( \tilde{a} - \tilde{b} \) is an element of \( k \) times the long root lattice of \( g \). The scaling dimension of \( \psi_\tilde{a} \) is given by

\[
h_{\tilde{a}}^\gamma = -\frac{a^2}{2k} + n(\tilde{a}) ,
\]  

(76)

where \( n(\tilde{a}) \) is an integer equal to the minimum number of roots from which \( \tilde{a} \) is composed. The inner product \( a^2 = (\tilde{a}, \tilde{a}) \) is defined with respect to the quadratic form matrix of \( g \).

If we are considering quantum Hall states with \( N_f \) layers, then we would expect to see the appearance of the \( \hat{g}_k/u(1)^2 \) parafermion CFTs where \( g = \text{simple Lie algebra of rank } 2 \). There are only three simple Lie algebras of rank 2: \( su(3) \), \( so(5) \), and \( G_2 \). Of these, only \( su(3) \) is simply laced, so the long root lattice is the same as the root lattice. This means that for the pattern-of-zeros solutions that correspond to \( n = (0^k) \), we expect to see solutions that correspond to \( su(3)_k/u(1)^2 \) CFTs. In Appendix B, we will describe the \( su(3)_2/u(1)^2 \) CFT in more detail.

The parafermion CFTs for \( g = so(5) \) and \( g = G_2 \), on the other hand, are more complicated because the long root lattice is different from the root lattice. For example, \( so(5)_k/u(1)^2 \) CFTs will generically be relevant for \( n = (2k 2m) \) while \( (G_2)_k/u(1)^2 \) CFTs will generically be relevant for \( n = (3k 0 3k) \).

### IX. EXAMPLES OF PATTERN-OF-ZEROS SOLUTIONS

In this section, we examine explicitly several simple bilayer pattern-of-zeros solutions. We first fix the cluster structure \( n \) to have a simple form. Then we try to find all the solutions \( \{ S_\tilde{a} \} \) that satisfy the conditions listed in Sec. VII.

The simplest non-Abelian states can be obtained from the simplest cluster structures \( n = (1 1^k) \) and \( n = (2 0^k) \). These are the simplest non-Abelian generalizations of the bilayer Laughlin-Halperin states.

Note that by definition the ordering of the rows in \( n \) is arbitrary; we choose it so that \( n_{11} \neq 0 \). Interchanging the layers yields the same physical system but corresponds to interchanging \( n_{11} \) with \( n_{22} \) and \( n_{12} \) with \( n_{21} \), so two matrices \( n \) and \( n' \) that are related by such an interchange are regarded as equivalent.

To list the solutions for those simplest cases, we may use some known CFTs to construct the appropriate simple-current algebra that corresponds to the pattern-of-zeros solutions. Using this known CFT, we can then write down the wave function explicitly. However, the wave function that we write down may not be unique in some cases; there may be several independent polynomials that have the same pattern of zeros. This corresponds to there being several distinct CFTs whose simple-current algebra possesses the same pattern of zeros.
In the following examples, we will make this choice when necessary so that we can explicitly write down a wave function with a certain pattern of zeros.

A. \( n=(1,0) \)

These states are all Abelian and correspond to the Halperin \((m,n,l)\) states. The unit cell spanned by the \( \vec{n}_i \) consists only of the points \((1,0)\) and \((0,1)\); By translation invariance \( S_{(1,0)}=S_{(0,1)}=0 \). Thus these states are completely characterized by the matrix \( m \) and are of the form

\[
\Phi = \prod_{i<j} (z_i-z_j)^{m_{12}} \prod_{i<j} (w_i-w_j)^{m_{23}} \prod_{i,j} (z_i-w_j)^{m_{12}} \tag{77}
\]

The \( m \) matrix here is exactly the \( K \) matrix that describes all Abelian FQH states. We also have \( h^+_n=0 \) for all of these states—the CFT has no simple-current part and corresponds to a \( U(1)^2 \) Gaussian CFT.

Although the single-layer hierarchy states, such as the \( \nu=2/5 \) hierarchy state, do not have ideal single-layer wave functions, there are ideal multilayer states that have the same topological orders as those single-layer hierarchy states. For example, the topological order in the \( \nu=2/5 \) hierarchy state is described by the \( K \) matrix \( K=\left( \frac{1}{2}, \frac{1}{2} \right) \) (in the symmetric basis). Such a topological order can be represented by the ideal bilayer state with \( n=(1,0) \) and \( m=\left( \frac{1}{3}, \frac{2}{3} \right) \). So although the pattern-of-zeros construction does not directly classify those single-layer hierarchy states, their topological orders can still be described by the pattern-of-zeros approach.

B. \( n=(1,1) \)

This choice of \( n \) requires that the electron operators in the CFT must take the form

\[
\begin{align*}
V_{c1} &= \psi_a e^{i\xi_z \tilde{M}_{13} \phi_1}, \\
V_{c2} &= \psi_a e^{i\xi_z \tilde{M}_{23} \phi_2},
\end{align*}
\tag{78}
\]

where \( \psi^2=1 \) and \( \psi \) has scaling dimension that is integer or half integer. This latter fact can be obtained from the condition \( \Delta S_{\delta}(1,0,1,0,1,0)=\text{even} \). One general way of constructing such a simple-current operator is by expressing it as a set of Majorana fermions from several copies of the Ising CFT,

\[
\psi = \psi_1^{(1)} \cdots \psi_1^{(a)},
\tag{79}
\]

where \( \psi_1^{(i)} \) is the Majorana fermion from the \( i \)th Ising CFT. Such an operator has scaling dimension \( h^o_\psi =a/2 \) and gives rise to the following FOHQ wave function:

\[
\Phi(z_i,w_j) = \Pr\left[ \frac{1}{x_i-x_j} \right]^a \times \Phi_{ga}.
\tag{80}
\]

\( x_i \) represents the coordinates in both layers

\[
x_i = \begin{cases} 
  z_i & 1 \leq i \leq N_1 \\
  w_{i-N_1} & N_1 < i \leq N_2 
\end{cases}
\tag{81}
\]

\( \Phi_{ga} \) is defined in Eq. (72). Note that the simple-current algebra in this case implies \( \psi_1^{(1)} \psi_1^{(0)} = 1 \). That is, these states can all also be viewed as satisfying the cluster condition for \( n=\left( \frac{2}{3}, \frac{1}{3} \right) \) but with a different choice \( m \). For every pattern-of-zeros solution found here, there is an equivalent one for \( n=\left( \frac{2}{3}, \frac{1}{3} \right) \).

Later, we will list the solutions with the \( n=\left( \frac{2}{3}, \frac{1}{3} \right) \) cluster structure. Some of those solutions actually have \( n=\left( \frac{1}{3}, \frac{1}{3} \right) \).

C. \( n=(1,0) \)

Here, the electron operators must be of the form

\[
V_{c1} = \psi_a e^{i\xi_z \tilde{M}_{13} \phi_1},
V_{c2} = e^{i\xi_z \tilde{M}_{23} \phi_2},
\tag{82}
\]

The fact that \( \vec{n}_1=(1,1) \) however also forces \( \psi=1 \). Thus in fact the corresponding CFTs do not have a simple-current part; they all correspond to a \( U(1)^2 \) Gaussian CFT. All of these states are therefore Abelian and correspond to the Halperin \((m,n,l)\) states. Thus, all pattern-of-zeros solutions where \( n=\left( \frac{1}{3}, \frac{1}{3} \right) \) also satisfy the cluster condition for \( n=\left( \frac{2}{3}, \frac{1}{3} \right) \).

Using the cluster condition and the fact that \( S_{(1,0)}=S_{(0,1)}=0 \), it is easy to see that the pattern of zeros for these states are completely characterized by the matrix \( m \) and therefore that these states are identical to the Halperin states.

D. \( n=(1,0) \)

In this case, the electron operators in the CFT must take the form

\[
\begin{align*}
V_{c1} &= e^{i\xi_z \tilde{M}_{13} \phi_1}, \\
V_{c2} &= \psi_a e^{i\xi_z \tilde{M}_{23} \phi_2},
\end{align*}
\tag{83}
\]

where \( \psi^2=1 \). From \( \Delta S_{\delta}(0,1,1,0,1,0)=\text{even} \), it follows that the scaling dimension of \( \psi \) is integer or half integer. This means that we can in general write it as a product of \( a \) Majorana fermion operators from \( a \) independent copies of the Ising CFT. Thus the pattern-of-zeros solution for this choice of \( n \) includes only two classes of states: the Abelian Halperin states and the following version of the Pfaffian states:

\[
\Phi((z_i,w_j)) = \Pr\left[ \frac{1}{z_i-w_j} \right]^a \times \Phi_{ga}.
\tag{84}
\]

Such a state spontaneously breaks the discrete \( Z_2 \) symmetry associated with interchanging the two layers.

E. \( n=(2,0) \): Non-Abelian bilayer states

For this choice of \( n \), the corresponding CFTs have two simple current operators, \( \psi_{(0,1)} \) and \( \psi_{(1,0)} \), which each square to the identity: \( \psi_{(0,1)}^2 = \psi_{(1,0)}^2 = 1 \). Thus there are a total of three distinct primary currents: \( \psi_{(1,0)}, \psi_{(0,1)}, \) and \( \psi_{(1,1)} = \psi_{(0,0)} \). The pattern of zeros can be fully specified by specifying the scaling dimensions of these simple current operators, \( h^o_{(1,0)}, h^o_{(0,1)}, \) and \( h^o_{(1,1)} \) and the matrix \( m \), which specifies the \( U(1)^2 \) part of the electron operators in the CFT.
Applying $\Delta_s(\bar{a}, \bar{a}, \bar{a}) = e \text{ven}$ for this choice of $n$ shows that the simple-current scaling dimensions are all integer or half integer: $2h^i_n \in \mathbb{Z}$. One general way of encoding these fusion relations and the associated scaling dimensions is to write the electron operators in the CFT using simple-current operators that consist of multiple copies of the $su(3)/2u(1)^2$ parafermion CFT. The most general choice for the electron operators can be written in the form

$$V_{c,1} = \psi^{(a)}_{\alpha_1} \ldots \psi^{(a)}_{\alpha_n} e^{i \Sigma_M J_{1M} \phi^c},$$

$$V_{c,2} = \psi^{(a-b-c)}_{\beta_1} \ldots \psi^{(a-b-c)}_{\beta_n} e^{i \Sigma_M J_{2M} \phi^b},$$

where $\psi^{(a)}_{\alpha}$ is a simple current from the $\alpha$th copy of the $su(3)/2u(1)^2$ parafermion CFT and $\alpha$ and $\beta$ refer to either $(1,0)$, $(0,1)$, or $(1,1)$. Some explicit forms for such correlators of simple-current operators in the $su(3)/2u(1)^2$ parafermion CFT were discussed in Ref. 20. Computing these correlation functions provides one way—not necessarily unique—of constructing a wave function with the desired pattern of zeros.

These pattern-of-zeros solutions can naturally be grouped into two distinct classes. In the first class, $V_{c,1}$ and $V_{c,2}$ contain only one kind of simple current, either $\psi^{(0,1)}_1$, $\psi^{(1,0)}_2$, or $\psi^{(1,1)}_3$, from each copy of the $su(3)/2u(1)^2$ CFT. For example, if $\psi^{(1,0)}_1$ appears in $V_{c,1}$, then $V_{c,2}$ cannot contain $\psi^{(1,1)}_3$ or $\psi^{(0,1)}_2$. In such a situation, we can think of $\psi^{(0)}_1$ as being the Majorana fermion from the Ising CFT. This means that these states can be written in terms of multiple copies of the Ising CFT. This class of multilayer states can be written by choosing

$$V_{c,1} = \psi^{(1)} \ldots \psi^{(a)} e^{i \Sigma_M J_{1M} \phi^c},$$

$$V_{c,2} = \psi^{(a-b-c)} \ldots \psi^{(a-b-c)} e^{i \Sigma_M J_{2M} \phi^b},$$

where $\psi^{(a)}$ is the Majorana fermion from the $a$th copy of the Ising CFT. The wave function for this class of states is therefore,

$$\Phi(\{z_i, w_j\}) = \text{Pr} \left( \frac{1}{z_i - z_j} \right)^{a-c} \text{Pr} \left( \frac{1}{w_i - w_j} \right)^{b-c} \text{Pr} \left( \frac{1}{x_j - x_i} \right)^c \prod_{i<j} (z_i - z_j)^{(m(i/2)} \prod_{i<j} (w_i - w_j)^{(m(1/2)} \prod_{i<j} (z_i - w_j)^{(m(1/2)}. \tag{87}$$

$x_i$ represents the coordinates in both layers,

$$x_i = \begin{cases} z_i & 1 \leq i \leq N_1 \\ w_i-N_1 & N_1 < i \leq N_2. \end{cases} \tag{88}$$

This is the simplest generalization of the single-layer non-Abelian states to a class of non-Abelian bilayer states: the interlayer Pfaffian states. The simplest version of this, with $a=b=c$, is closely related to (and identical to for certain choices of $m$) the spin-charge separated non-Abelian spin-singlet wave function proposed in Ref. 26.

The second class of states cannot be written in terms of multiple copies of the Ising CFT—the full $su(3)/2u(1)^2$ parafermion CFT is necessary. The first class of states, which can be written only using the Ising CFT, have the property that their pattern of zeros satisfies $\Delta_s(\bar{a}, \bar{b}, \bar{c}) = e \text{ven}$ for all choices of $\bar{a}$, $\bar{b}$, and $\bar{c}$. The second class of states, for which the full $su(3)/2u(1)^2$ CFT is necessary, satisfies $\Delta_s(\bar{a}, \bar{b}, \bar{c}) = e \text{odd}$ for certain choices of $\bar{a}$, $\bar{b}$, and $\bar{c}$. Let us compare the simple-current algebra of the $su(3)/2u(1)^2$ parafermion CFT to what one would obtain using two copies of the Ising CFT. If we used two copies of the Ising CFT, we could have $\psi^{(1,0)} = \psi^{(1)}$ with scaling dimension $1/2$, $\psi^{(0,1)} = \psi^{(2)}$ with scaling dimension $1/2$, and $\psi^{(1,1)} = \psi^{(1)} \psi^{(2)}$ with scaling dimension 1. This satisfies $\Delta_s[(1,0), (0,1), (1,1)] = e \text{ven}$. On the other hand, in the $su(3)/2u(1)^2$ parafermion CFT, the only difference is that $\psi^{(1,1)}$ also has scaling dimension 1/2. Thus in this latter theory, two fermions combine to give another fermion. This yields $\Delta_s[(1,0), (0,1), (1,1)] = e \text{odd}$. The fact that there are valid single-valued translationally invariant FQH wave functions that arise from unitary CFTs and that have $\Delta_s(\bar{a}, \bar{b}, \bar{c}) = e \text{odd}$ for certain choices of $\bar{a}$, $\bar{b}$, $\bar{c}$ suggests (see Sec. V E) that we should impose $\Delta_s(\bar{a}, \bar{b}, \bar{c}) = e \text{odd}$ even not in general but only if $\bar{a}$, $\bar{b}$, and $\bar{c}$ are collinear through the origin.

Let us examine the pattern of zeros for a few of the simplest examples of these non-Abelian bilayer states. There is a fermionic $\nu = 2/3$ state with

$$m = \begin{pmatrix} 2 & 2 \\ 2 & 2 \end{pmatrix} \quad \{S_{(2,0)} = 0, S_{(1,1)} = 0, S_{(0,2)} = 0\},$$

$$\nu = 2/3 \begin{pmatrix} h_{(1,0)}^c = 1/2, h_{(1,1)}^c = 0, h_{(0,1)}^c = 1/2 \end{pmatrix}. \tag{89}$$

This is the pattern of zeros for the interlayer Pfaffian state, which is of the form

$$\Psi(\{z_i, w_j\}) = \text{Pr} \left( \frac{1}{x_j - x_i} \right) \Phi_{(2,2,1)}(\{z_i, w_j\}). \tag{90}$$

We use the notation

$$\Phi_{(a,b,c)} = \prod_{i < j} (z_i - z_j)^a \prod_{i < j} (w_i-w_j)^b \prod_{i,j} (z_i-w_j)^c \tag{91}$$

There are also fermionic states at $\nu = 4/5$ and $\nu = 4/7$. These have the following pattern of zeros:
The state at $\nu=4/7$ is the non-Abelian spin singlet state that was proposed in Ref. 19.

Note once again that the pattern of zeros $m$ and $\{S_{ij}\}$ refer to the pattern of zeros of the symmetric polynomial, $\Phi = \Pi_{\nu\neq\nu} (x-x')$.

X. DISCUSSION OF RESULTS AND RELATION TO EXPERIMENT

In single-layer quantum Hall samples, a quantum Hall plateau is seen at $\nu=5/2$, but not at $\nu=1/2$. The reason is that even though in all of these cases there is a single half-filled Landau level, the existence of the two filled extra Landau levels modifies the effective interactions between the electrons in the unfilled level. In the $\nu=5/2$ case, numerical calculations suggest that these effective interactions are modified in such a way that a non-Abelian quantum Hall state may be realized.

Experiments on multicomponent quantum Hall systems should be able to probe an even wider variety of regimes with distinct effective interactions. For example, for a two-component FQH system, we can study systems in which the spin degree of freedom is present, two-dimensional electron systems with two quantum wells, wide single-layer systems in which the electrons spontaneously form a double-layer system due to Coulomb repulsion, or systems in which there may be two valleys for the free quasiparticle spectrum (such as in graphene or SiGe heterostructures), etc. In many of these cases, experimentalists can also tune to some extent the degree of correlation between the two components. For example, in double-layer systems, application of a parallel magnetic field can tune the tunneling and correlation between the layers. There may also be some degree of tunability in the relative densities between the two components in addition to being able to probe FQH states with different numbers of filled Landau levels. With this greatly increased amount of variability and tunability in the effective interactions between electrons in the unfilled Landau levels, it is possible that a non-Abelian state can be realized in a two-component quantum Hall system.

Since the pattern of zeros provides a systematic classification and characterization of a wide variety of quantum Hall states, it provides us with a general sense of how all of the non-Abelian bilayer states are related and which ones are simpler than other ones. Just as we know that the single-layer Pfaffian quantum Hall state is the simplest non-Abelian generalization of the Laughlin states, we can determine the simplest non-Abelian generalization of the Halperin bilayer states and therefore single out some of the possibilities that may be experimentally viable.

In Ref. 27, we have given an overview of some of the simplest non-Abelian bilayer states that we find and that occur at filling fractions at which experiments on two-component FQH systems have already observed incompressible states. Here we briefly summarize that discussion and supplement details of the calculations of various topological properties of the candidate states.

Experiments have so far observed FQH plateaus in two-component systems at $\nu=2/3$, 4/5, 4/7, 4/9, 6/5, 6/7, 1/4, etc. In some cases, these plateaus have been observed in both bilayer and spin-unpolarized single-layer systems while in others, the plateau has only been observed in one of them. At all of these filling fractions, there exists also one (or several) candidate Abelian phase(s); in most cases, it is assumed that these plateaus are described by one of the Abelian phases. However, the pattern-of-zeros construction also yields many simple non-Abelian states at these filling fractions. In some situations, we expect the non-Abelian states to be good candidate states.

There are at least two dimensionless quantities that are important determining factors for which FQH state is realized. The first parameter is $\alpha = V_{\text{inter}}/V_{\text{intra}}$, where $V_{\text{inter}}$ is the potential for interlayer repulsion and $V_{\text{intra}}$ is the potential for intralayer repulsion. The second parameter is $\gamma = \tau V_{\text{intra}}$, where $\tau$ is the interlayer hopping amplitude. In the limit $\alpha \rightarrow 0$ and $\gamma \rightarrow 0$, the system will be a FQH state that consists of two independent single-layer FQH states in each layer. In the limit $\gamma \gg 1$ and $\alpha \sim 0$, a single-layer FQH state may be observed. But if we keep $\gamma \sim 0$ and increase $\alpha$ from $\sim 0$, then the FQH state formed by two independent single-layer FQH states in each layer must undergo a phase transition into either a compressible phase or a new incompressible state. In the latter case, an Abelian hierarchy state (such as a bilayer composite fermion state) may form, which would in most cases be a state described by a $4 \times 4$ or more complicated $K$ matrix and would have four or more edge modes. The other possibility is that a non-Abelian two-component state may form. The pattern-of-zeros construction yields non-Abelian two-component wave functions that have zeros when particles from different layers approach each other, indicating that they can accommodate situations in which $\alpha \sim 1$. Additionally, these states generally have less than four edge modes; if we use the number of edge modes as a measure of the complexity of the state, then the non-Abelian states are simpler and may therefore be realized experimentally.

At $\nu=2/3$, experiments on wide single quantum wells have observed a phase transition from a bilayer to single-layer state while experiments on single-layer systems have seen a phase transition from a spin-polarized to a spin-unpolarized state. In the limit $\alpha \sim 0$ and $\gamma \sim 0$, the system should be in the $(3,3,0)$ state. As $\alpha$ is increased while $\gamma \sim 0$, one possibility is the $(1,1,2)$ state. This wave function appears unphysical because it has higher-order zeros as particles from different layers approach each other than particles from within the same layer. Another wave function, which
TABLE I. Proposed explanations for incompressible states at experimentally relevant filling fractions, $\nu=2/3$, 4/5, 4/7, and 1/4, in two-component FQH systems. The bilayer composite fermion state $(n_1, n_2 | m)$ (Ref. 35) refers to the state $\Pi_j (z_j^+ - w_j)^m \Phi_{\nu_1} ((z_j^+) \Phi_{\nu_2} ((w_j))$, where $\Phi_{\nu}$ is a single-layer composite fermion state at filling fraction $\nu$. For $(2/3, 2/3 | m)$, we have taken the single-layer 2/3 state to be the particle-hole conjugate of the Laughlin state. $n_{k1}$ and $n_{l1}$ indicates that there are $n_{k1}$ right-moving edge modes and $n_{l1}$ left-moving edge modes. See Appendix C for details of how to calculate the number of edge modes and the shift $\mathcal{S}$.

| $\nu$ | Proposed states | Edge modes | Shift $\mathcal{S}$ |
|-------|----------------|------------|------------------|
| 2/3   | $(3,3,0)$      | 2          | 3                |
|       | $(1,1,2)$      | 2          | 1                |
|       | $2/3 |_{\text{inter}}$ [see Eq. (94)] | $2 + \frac{1}{2}$ | 2 |
|       | $2/3 |_{\text{intra}}$ [see Eq. (95)] | 3          | 3                |
|       | $\nu=1/3$      | $1 + 1_{L}$ | 0                |
| 4/5   | $su(3,2)_{\nu}(u(1))^2$ [see Eq. (96)] | $2 + 2_{L}$ | 2          |
|       | $(2,3,2/3 | 1)$ | 2          | 0                |
| 4/7   | $su(3,2)_{\nu}(u(1))^2$ [see Eq. (96)] | $2 + \frac{3}{2}$ | 4          |
|       | $(2,5,2/3 | 1)$ | 4          | 4                |
|       | $(2,3,2/3 | 2)$ | $1 + 3_{L}$ | 0                |
| 1/4   | $(7,7,1)$      | 2          | 7                |
|       | Interlayer Pfaffian [see Eq. (98)] | $2 + \frac{1}{2}$ | 7          |
|       | Single-layer Pfaffian | $1 + 1_{L}$ | 5                |

TABLE II. Quasiparticle operators from the CFT for $\nu=2/3$ interlayer Pfaffian states. The scalar bosons $\phi_\sigma$ and $\phi_\rho$ are related to the total and relative density fluctuations of the two layers, respectively. $\sigma$ is the spin field in the Ising CFT, which has scaling dimension 1/16.

| Quasiparticles for interlayer Pfaffian at $\nu=2/3$ | CFT operator | Total charge | Scaling dimension |
|--------------------------------------------------|---------------|--------------|------------------|
| $\Psi_{4/7} = \Phi_{\nu_1}((z_i, w_i)) \Phi_{(2,2,3/2)}((z_i, w_j))$ | $V_{4/7} = \psi_{\nu_1}^{(1/2)} \Phi_{\nu_1}((z_i, w_j))$ | 1           | 3/2             |
| $\Psi_{4/9} = \Phi_{\nu_1}((z_i, w_i)) \Phi_{(4,4,1/2)}((z_i, w_j))$ | $V_{4/7} = \psi_{\nu_1}^{(1/2)} \Phi_{\nu_1}((z_i, w_j))$ | 1           | 3/2             |
| $\sigma^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $\psi^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $1/3$ | 1/3             |
| $\sigma^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $\psi^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $1/3$ | 1/3             |
| $\sigma^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $\psi^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $1/3$ | 1/3             |
| $\psi^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $\psi^{(1/3)} \Phi_{\nu_1}((z_i, w_j))$ | $1/3$ | 1/3             |

The other set of proposed Abelian states are the bilayer composite fermion states $\Phi_{\nu_1}((z_i, w_j))$ is a single-layer FQH state at filling fraction $n_0$. These states have four edge modes, indicating that they may be less stable than the alternative non-Abelian possibilities.

Recently, an incompressible state was found at $\nu=1/4$ and it is unclear what phase this corresponds to and even whether it is a single-layer or double-layer phase. Some possibilities that have recently been considered are the $(5,5,3)$ and $(7,7,1)$ Halperin states and the $\nu=1/4$ single-layer Pfaffian. The pattern-of-zeros construction yields many other alternative possibilities, perhaps the most physical (and simplest) of which is the following interlayer Pfaffian:

$\Psi_{\nu_1}((z_i, w_j)) = \Phi_{\nu_1}((z_i, w_j)) \Phi_{(6,6,2)}((z_i, w_j))$. 

In Table I, we summarize some of the filling fractions at which incompressible states have been experimentally observed in two-component FQH systems. For each filling fraction we list some of the proposed wave functions that may characterize the topological order of those phases, the number of edge modes, and their respective shifts on the sphere. We list the quasiparticles, their electric charges, and their scaling dimensions for the interlayer Pfaffian state at $\nu$.
TABLE III. Charge and scaling dimensions of the quasiparticle operators with minimal nonzero total charge in the non-Abelian bilayer states discussed here. In the scaling dimension, the first term comes from the non-Abelian part, the second term comes from the total density fluctuations [the U(1) part], the third term comes from the relative density fluctuations of the two layers [also the U(1) part].

| ν       | Charge q_{min} | Scaling dimension h |
|---------|----------------|---------------------|
| 2/3_{\text{inter}} | 1/3 | 1/16 + 1/12 + 0 |
| 2/3_{\text{intra}}  | 1/6 | 1/16 + 1/48 + 1/16 |
| 4/5      | 1/5 | 1/10 + 1/40 + 1/25 |
| 4/7      | 1/7 | 1/10 + 1/90 + 1/35 |
| 4/9      | 1/9 | 1/10 + 1/90 + 1/35 |
| 1/1      | 1/8 | 1/16 + 1/32 + 0 |

=2/3 in Table II. In Table III, we list the quasiparticles with the minimal electric charge and their scaling dimensions h for the non-Abelian FQH states discussed in this paper [see Eqs. (90), (92), (93), (96), and (98)]. Those minimally charged quasiparticles may dominate interedge tunneling and give rise to the following I-V curve: I ∝ V^{d-1} in the T=0 limit.

In summary, we find many simple non-Abelian bilayer states that occur at experimentally observed filling fractions. For certain effective interactions among the electrons in the unfilled Landau levels, these states may be more favorable than their Abelian counterparts. In these cases, the non-Abelian states have larger interlayer correlations and therefore may be energetically more favorable in situations in which the interlayer repulsion is comparable to the intralayer repulsion.

Conditions on filling the sphere

A useful tool for identifying FQH states in numerical studies of exact diagonalization on finite systems on a sphere is to look at what values of the shift, S=ν⁻¹N − N_{\Phi}, a ground state with zero total angular momentum is found. This then limits the possibilities of which topological phase is realized in the system to those that have that particular value of the shift. Similarly, in such numerical studies of multilayer systems, one can look for the different sets (N_{1},...,N_{\nu};N_{\Phi}^{1},...,N_{\Phi}^{\nu}) that yield a ground state with zero total angular momentum. Each topological phase will have its own list of (N_{1},...,N_{\nu};N_{\Phi}^{1},...,N_{\Phi}^{\nu}) that let it fill the sphere; analyzing this can be a useful way of determining which topological phase is obtained numerically. In Sec. V D, we found conditions that N and N_{\Phi} should satisfy for the FQH state to fill the sphere.

For states that have a cluster structure n=\binom{1}{0}_{2}, we find that the condition (46) becomes trivial as long as \tilde{N}=∑_{\mu}k_{\mu}n_{\mu}, where k_{\mu} is an integer. This means that as long as N_{1} and N_{2} are even and N_{\Phi}^{1}, N_{\Phi}^{2} satisfy Eq. (44), then these states can fill the sphere. In this case, we find that Eq. (44) reduces to the form

\begin{equation}
\begin{pmatrix}
N_{\Phi}^{1} + S \\
N_{\Phi}^{2} + S
\end{pmatrix} = M \begin{pmatrix} N_{1} \\ N_{2} \end{pmatrix},
\end{equation}

where M is a 2×2 matrix and S is the shift, which can be calculated using Eq. (C11). The states that we have been considering are of the form \Phi=\Phi_{\text{sc}} × \Phi_{\text{(ξ,β,y)}}, for which M = \binom{\beta \beta}{\alpha \alpha}. Tables IV and V lists some examples.

XI. SUMMARY AND OUTLOOK

In this paper, we generalized the pattern-of-zeros characterization and classification of FQH states to multicomponent cases. We found that the topological orders in a multicomponent FQH state can be characterized by the following data: a matrix n that describes the cluster structure, a matrix m and a sequence \{S_{\alpha}\} that describes the pattern of zeros.

Our pattern-of-zeros characterization gives us a general quantitative view on a large class of Abelian and non-Abelian bilayer FQH states, which allow us to determine which states are simpler than other states. We find some of the simplest non-Abelian generalizations of the Laughlin-Halperin Abelian bilayer states. Those simple non-Abelian states may describe some of the bilayer/spin-unpolarized FQH states observed in experiments and numerical calculations.

The discussion here has been restricted to idealized ground-state wave functions for multicomponent FQH states. In order to describe the corresponding quasihole wave functions, we must generalize the pattern-of-zeros framework for
TABLE V. Values of \((N_1^1, N_2^0)\) that yield rotationally invariant states on the sphere for various choices of \((N_1, N_2)\) for the \(\nu=4/5\) \(su(3)/u(1)^2\) parafermion state.

| \(N_2\) | 2   | 4   | 6   | 8   | 10  |
|-------|-----|-----|-----|-----|-----|
| 6     | (2,2) | (6,3) | (10, 4) | (14,5) | (18,6) |
| 8     | (3,6) | (7,7) | (11,8) | (15,9) | (19,10) |
| 10    | (4,10) | (8,11) | (12,12) | (6,13) | (20, 14) |
| 12    | (5,14) | (9,15) | (13,6) | (17,17) | (21,18) |
| 14    | (6,18) | (10,19) | (14,20) | (18,21) | (22,22) |

quasiparticles, which was developed for single-layer states in Refs. 11 and 12. We expect that generalizations of the constraints found in Refs. 11, 12, and 38 may provide interesting insights into the nature of topological order in multicomponent \(FQH\) states.

Finally, we comment on the extent to which this multicomponent pattern-of-zeros framework may be a full classification of possible topological orders in \(FQH\) states. It is known that all Abelian \(FQH\) states can be described within the \(K\)-matrix formalism. Consider, for example, the \(\nu=2/5\) hierarchy state for a single layer of electrons, which is described by a \(2 \times 2\) \(K\) matrix. As mentioned in Sec. IX A, while we have model wave functions for this state—in terms of Jain’s composite fermion wave functions or in terms of the Halperin-Haldane construction—there is no single-layer ideal wave function that has the topological properties of this state. However, we can write down a double-layer ideal wave function that does have the same topological properties as this state. In this sense, ideal wave functions can classify all possible topological orders that may appear in Abelian \(FQH\) states. The multicomponent pattern-of-zeros framework described here can be viewed as a generalization of the \(K\)-matrix classification: given any possible topological phase that can appear in a \(FQH\) system, we expect that there is an ideal wave function—possibly for multiple flavors of electrons and described by the multilayer pattern-of-zeros framework—which has the same topological properties. In this sense the multilayer pattern-of-zeros framework may provide the foundation for a full classification of topological orders in \(FQH\) systems.

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APPENDIX A: OCCUPATION NUMBER CHARACTERIZATION

In the single-component pattern-of-zeros description, there is an occupation number characterization that is a useful way to understand both the ground states and the quasiparticles in \(FQH\) states. The generalization to multilayer states does not appear to be quite as simple or useful but for the sake of completeness we will analyze it below.

A convenient set of single-particle basis states for particles in the lowest Landau level are monomials of the form \(z^m\), for integer \(m\). Thus, a basis for symmetrized wave functions of \(N_f\) particles of type \(I\) is given by

\[
\Phi_{\{\eta\}} = \sum_{\{P\}} \prod_i P_i \prod_{\eta_i} (z_{P_i(\eta_i)})^{\eta_i},
\]

where \(P_i\) is a permutation of the particles of type \(I\), \(l_i^I\) is an integer, and \(\eta_i\) is a vector whose component \(\eta_i^I\) is the number of particles of type \(I\) that occupy the \(l_i\)th orbit. The polynomial \(\Phi\) that we are interested in can be expanded in terms of these basis states as

\[
\Phi = \sum_{\{\eta\}} C_{\{\eta\}} \Phi_{\{\eta\}},
\]

Now we may ask what kind of boson occupations \(\{\eta_i\}\) will be present in the sum (A2) for a polynomial \(\Phi\) with a given pattern-of-zeros \(\{S_{\eta}\}\). To answer this question, let us set \(z_1^I = 0\) in \(\Phi(\{z_2^I\})\). Since \(\Phi\) is nonzero when \(z_1^I = 0\) due to translation invariance, there must be a boson occupation \(\{\eta_i\}\) in the above sum that contains at least one boson occupying the \((z_1^I)^{m_1}\) orbital. That is, there is a term in the above sum with \(n_0^I > 0\). Now, suppose that we bring a second particle of the same type, \(z_2^I\) to 0. The minimal power of \(z_2^I\) in \(\Phi(0,z_2^I,z_3^I,\ldots)\) is

\[
\Phi(0,z_2^I,z_3^I,\ldots) \simeq (z_2^I)^{D_{z_2^I z_1^I}} P_2(z_3^I,z_4^I,\ldots) + O((z_2^I)^{D_{z_2^I z_1^I}+1}).
\]

Thus, among those \(\{\eta_i\}\) that have at least one boson of type 1 occupying the \((z_1^I)^{m_1}\) orbital, there must also be an \(\{\eta_i\}\) that contains a second boson of type 1 occupying the \((z_1^I)^{m_2}\) orbital where \(l_{z_2^I} = D_{z_2^I z_1^I} = S_{z_2^I} - S_{z_1^I}\). Next, assume that two bosons occupy the \((z_1^I)^{m_1}\) and \((z_1^I)^{m_2}\) orbital, and bring a third particle of type 1 to 0; the minimal power of \(z_3^I\) is \(D_{z_3^I z_1^I}\):

\[
P_3(z_3^I,z_4^I,\ldots) \simeq (z_3^I)^{D_{z_3^I z_1^I}} P_3(z_4^I,z_5^I,\ldots) + O((z_3^I)^{D_{z_3^I z_1^I}+1}).
\]
this way, we see that there must be a type 1 boson occupying the orbitals \( l_{\alpha l}^1 = S_{\alpha l} - S_{(a-1)e_l} \) for \( a=1, \ldots, N_l \). After taking all the type 1 particles to zero, we may begin to take the type 2 particles to zero, one by one, thus obtaining that there must be a type 2 boson occupying the orbitals \( l_{\alpha l}^2 = S_{\alpha l}^2 - S_{(a-1)e_l} \). Continuing this argument for bosons of every type, we find that there must be a term in the sum (A2) with occupation number described by the above sequence of \( l_{\alpha l}^i \).

However, in the above argument, we chose a particular sequence in which to take various particles to zero. We first took all of the type 1 particles to zero one-by-one, and then all of the type 2 ones, and so on. But we could just as well have made the argument with any sequence. Suppose that after taking \( i \) particles to the origin, there is a \( \tilde{a}_i \) cluster at the origin. Thus \( \{\tilde{a}_i\} \) is a sequence that describes the order in which we take particles to the origin until all particles are at the origin. For every such sequence, we may make the above argument and argue that if \( \tilde{a}_{a_l} = \tilde{a}_i + e_f \), then there must be a type \( l \) boson occupying the orbital \( l_{\alpha l}^i = S_{\alpha l} - S_{\tilde{a}_l} \). If we enumerate all the different sequences \( \{\tilde{a}_i\} \) by an integer \( \alpha \), then by considering each \( \alpha \), we see that there must be a term in the sum (A2) with occupation number \( \tilde{n}_i^\alpha \). \( n_i^\alpha \) would be the number of \( i \), along the sequence \( \{\tilde{a}_i\} \), for which \( l_{\alpha l}^i = l \). Notice that \( l_{\tilde{a}_l}^i \) must be nonzero. Thus we have the following important condition on \( S_{\tilde{a}_l}^i \):

\[
l_{\tilde{a}_l}^i = S_{\tilde{a}_l} - S_{\tilde{a}_l}^i \geq 0. \tag{A5}
\]

The analysis above can be thought of in the following way. Consider an \( N_f \)-dimensional lattice \( \mathbb{Z}^{N_f} \), where \( N_f \) is the number of layers. At every site \( \tilde{a} \) of this lattice (\( a_i \equiv 0 \)) we can associate the nonnegative integer \( S_{\tilde{a}} \). On each link (\( \tilde{a}, \tilde{a} - e_f \)), of the lattice we may also associate an integer \( l_{\tilde{a} - e_f} = S_{\tilde{a}} - S_{\tilde{a} - e_f} \). Now consider any directed path from the origin to \( N \) (\( N \) is the number of particles of type \( i \), in which the sum of the coordinates of every point on the path is one larger than the sum of the coordinates of the point preceding it, and enumerate the set of these paths by \( \alpha \). To each such path we associate an occupation number sequence \( \tilde{n}_i^\alpha \), where \( n_i^\alpha \) is the number of links along the path \( \alpha \) whose \( l_{\tilde{a} - e_f} = l \). If \( \Phi \) has a pattern-of-zeros \( \{S_{\tilde{a}}\} \), then its basis expansion (A2) must contain a term with occupation number \( \tilde{n}_i^\alpha \). Thus we may rewrite Eq. (A2) as

\[
\Phi = \sum_\alpha C_{\alpha} \Phi_{n_i^\alpha} + \sum_{\{\tilde{a}_i\}} D_{\{\tilde{a}_i\}} \Phi(\{\tilde{a}_i\})(\{l_{\alpha l}^i\}). \tag{A6}
\]

The two sequences \( \{S_{\tilde{a}}\} \) and \( \{\tilde{n}_i^\alpha\} \) contain the same information and are one-to-one labelings of each other. However, \( \{\tilde{n}_i^\alpha\} \) is redundant in the sense that it does not need to be specified for every \( \alpha \) in order to reconstruct \( l_{\tilde{a}_l}^i \). The \( \tilde{n}_i \) that appear in the second sum characterize the subleading terms that appear when coordinates are brought together; thus those \( \{\tilde{n}_i\} \) correspond to sequences \( \{S_{\tilde{a}}\} \) where \( S_{\tilde{a}} \equiv S_{\tilde{a} - e_f} \).

In the single-layer case, \( \{l\} \) naturally define an occupation number sequence \( \{n_l\} \), which also described the FQH state in the thin-cylinder limit. In the multicomponent generalization, we have \( \{l_{\alpha l}\} \), which seems to admit no simple generalization of the above occupation number sequence. Instead, one has such occupation number distributions for a large number of sequences which we enumerated above by \( \alpha \). We have not analyzed on general grounds which particular sequences contribute the most weight to the wave function in the thin cylinder limit.

**APPENDIX B: \( su(3)_2/u(1)^2 \) PARAFOERM CONFORMAL FIELD THEORY**

Some of the simplest non-Abelian bilayer states are closely related to \( su(3)_2/u(1)^2 \) paraffin CFT. This CFT has central charge \( c=6/5 \) and has three simple currents, \( \psi, \psi\beta, \psi\alpha \), all of which square to the identity and have scaling dimension 1/2.

There are four other primary fields, which are associated with the fundamental representation of \( su(3) \). Their scaling dimensions are listed in Table VI. The fusion rules for these fields all follow from the following fusion rule:

\[
\sigma \times \sigma = 1 + \sigma\beta. \tag{B1}
\]

**APPENDIX C: CALCULATIONS FOR CANDIDATE STATES**

1. Number of edge modes

The total number of edge modes is equal to the central charge of the corresponding CFT for the states that are described by the pattern of zeros. For the hierarchy states, the number of edge modes is given by the rank of the \( K \) matrix. Furthermore, in the latter case, the number of right- (left)-moving edge modes is given by the number positive (negative) eigenvalues of the \( K \) matrix.

The interlayer Pfaffian states are described by a CFT that consists of the Ising CFT, with \( c=1/2 \), and two scalar boson CFTs, each with \( c=1 \). Thus the number of edge modes for the interlayer Pfaffian is \( 2+2 \).

The intralayer Pfaffian states have two Ising CFTs in addition to the two scalar boson CFTs, so the total number of edge modes is 3.

The central charge of the \( su(3)_2/u(1)^2 \) paraffin CFT is \( c=6/5 \); the two-component FQH states based on this are described by the \( su(3)_2/u(1)^2 \) theory and two scalar bosons, for a total of \( 2^2\delta \) edge modes.

For the \( (m,m,l) \) states, the \( K \) matrix is \( K=(m-1)^l \). These states have two edge modes; if \( m=1 \), all edge modes move in
the same direction; if \( m < l \), then there is one right-moving and one left-moving edge mode.

For the states \( (\nu_0, \nu_l|m) \), the \( K \) matrix is
\[
K = \begin{pmatrix}
K_{11}^0 & K_{12}^0 & m & 0 \\
K_{21}^0 & K_{22}^0 & 0 & 0 \\
m & 0 & K_{31}^0 & K_{32}^0 \\
0 & 0 & K_{41}^0 & K_{42}^0
\end{pmatrix},
\]
(C1)
where \( K^0 \) is the \( K \) matrix in the hierarchical basis of the state \( \Phi_{\nu_0} \). For the \( \nu=2/5 \) state, \( K^0 \) in the hierarchical basis is
\[
K^0 = \begin{pmatrix}
3 & -1 & 0 \\
-1 & 2 & 0 \\
0 & 0 & 0
\end{pmatrix}.
\]
(C2)
For the \( \nu=2/7 \) state, \( K^0 \) in the hierarchical basis is
\[
K^0 = \begin{pmatrix}
3 & 1 & 0 \\
1 & 1 & 0 \\
0 & 0 & 0
\end{pmatrix}.
\]
(C3)
For the \( \nu=2/3 \) P-H conjugate of the 1/3 Laughlin state, \( K^0 \) in the hierarchical basis is
\[
K^0 = \begin{pmatrix}
1 & 1 & 0 \\
1 & 1 & 0 \\
0 & 0 & 0
\end{pmatrix}.
\]
(C4)
Using this, we find that the P-H conjugate of the 1/3 Laughlin state has edge modes \( 1_R + 1_L, (2/3, 2/3) \) has edge modes \( 2_R + 2_L \), and \( (2/3, 2/3) \) has \( 1_R + 3_L \).

2. Shifts on sphere

In the hierarchy basis, the formula for the shift is given by
\[
S = \frac{1}{\nu} \sum_J (K^{-1})_{IJ}K_{JJ}.
\]
(C5)
Using this formula, we find \( S=0 \) for the particle-hole conjugate of the \( \nu=1/3 \) Laughlin state.

Now consider the bilayer composite fermion state \( (\nu_0, \nu_l|m) \),
\[
\Phi_{(\nu_0, \nu_l|m)} = \prod_{i,j} (z_i - w_j)^m \Phi_{\nu_0}(\{z_i\}) \Phi_{\nu_l}(\{w_i\}).
\]
(C6)
Let \( N_\Phi^0 \) be the maximum power of \( z_i \) in \( \Phi_{\nu_0}(\{z_i\}) \). It satisfies
\[
N_\Phi^0 = \nu_0^{-1} N_1 - S_0,
\]
where \( S_0 \) is the shift of the state \( \Phi_{\nu_0} \). The factor \( \prod_{i,j}(z_i - w_j)^m \) increases the power of \( z_i \) by \( mn_2 \). Thus the maximum power of \( z_i \) in \( \Phi_{(\nu_0, \nu_l|m)} \) is \( N_\Phi^0 \).

3. Electron and quasiparticle operators for \( su(3)_2/u(1)^2 \) states

The electron operators for the \( su(3)_2/u(1)^2 \) FQH states that we discuss are of the form
\[
V_{e1} = \psi e^{i(1/\nu)\phi_+ + i\phi_-},
\]
\[
V_{e2} = \psi e^{i(3/\nu)\phi_+ + i\phi_-},
\]
where \( \nu = 3/2, 1/2, \) and \( \sqrt{3}/2 \) for the \( \nu=4/5, 4/7, \) and \( 4/9 \) states, respectively. The quasiparticle operators with minimal total charge are of the form
\[
V_{qp} = \sigma_0 e^{iQ(1/\nu)\phi_+ + i\phi_-},
\]
(C13)
and have scaling dimension \( h_{qp} = 1 + \frac{Q^2}{2\nu} + \frac{(\nu_0)^2}{2} \). The total charge of the quasiparticle is \( Q=1/5, 1/7, \) and \( 1/9 \) for the \( \nu=4/5, 4/7, \) and \( 4/9 \) states, respectively. \( s_{qp}=1/3, 1, \) and \( 1/7 \), respectively, for these states.

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