Bilayer paired quantum Hall states and Coulomb drag

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We consider a number of strongly correlated quantum Hall states that are likely to be realized in bilayer quantum Hall systems at total Landau level filling fraction $\nu_f = 1$. One state, the $(3,3,-1)$ state, can occur as an instability of a compressible state in the large $d/l_B$ limit, where $d$ and $l_B$ are the interlayer distance and magnetic length, respectively. This state has a hierarchical descendnet that is interlayer coherent. Another interlayer coherent state, which is expected in the small $d/l_B$ limit is the well-known Halperin $(1,1,1)$ state. Using the concept of composite fermion pairing, we discuss the wave functions that describe these states. We construct a phase diagram using the Chern-Simons Landau-Ginzburg theory and discuss the transitions between the various phases. We propose that the longitudinal and Hall-drag resistivities can be used together with interlayer tunneling to experimentally distinguish these different quantum Hall states. Our work indicates the bilayer $\nu_f = 1$ quantum Hall phase diagram to be considerably richer than that assumed so far in the literature.

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I. INTRODUCTION

Bilayer quantum Hall systems at total Landau level filling factor $\nu_f = 1$, i.e., $\nu = 1/2$ in each layer,1 allow for novel interlayer coherent phases. These phases have attracted a great deal of theoretical and experimental attention over the last 16 years, dating back to a seminal paper by Halperin2 in which the multicomponent generalization of the Laughlin wave function was first considered in a rather general context. In particular, there is strong experimental evidence2,4 and compelling theoretical basis5,6 to believe that a spin-polarized bilayer $\nu = 1/2$ quantum Hall system would have a novel spontaneous interlayer coherent incompressible phase for small values of the interlayer separation $d$, even in the absence of interlayer tunneling. (We consider only the situation without any interlayer tunneling in this paper, our considerations also apply to the physical situation with weak interlayer tunneling. The situation with strong interlayer tunneling is trivial by virtue of the tunneling-induced symmetric-antisymmetric single-particle tunneling gap that leads to the usual $\nu_f = 1$ quantum Hall state in the symmetric band.) In the limit $d \rightarrow \infty$, however, one expects two decoupled layers each with $\nu = 1/2$ and hence no quantized Hall state. The phase diagram for this compressible to incompressible quantum phase transition in $\nu = 1/2$ bilayer systems has been studied extensively in the literature,2 but we still do not have a complete qualitative understanding of the detailed nature of this transition. In particular, one does not know how different kinds of incompressible (and compressible) phases compete as system parameters (e.g., $d$) are tuned, and how to distinguish among possible competing incompressible phases. In this paper, we revisit this issue by arguing that, in principle, there are several interesting and nontrivial quantum Hall phases in the $\nu_f = 1$ bilayer system that could be systematically probed via interlayer drag experiments carried out at various values of the interlayer separation $d$. One goal of our paper is to describe and discuss the rich $\nu_f = 1$ bilayer quantum Hall phase diagram using the Chern-Simons Landau-Ginzburg approach $6^{-10}$.

Among the more interesting quantum Hall states are the so-called paired Hall states, which have been extensively studied theoretically.3,11,12 In these states, the composite fermions form a superconducting paired state in which two composite fermions bind into effective Cooper pairs that condense into a ground state analogous to the BCS state. The well-studied Moore-Read Pfaffian state13 is a spin-polarized version of such a paired Hall state for a single-layer system. Bilayer paired Hall states have been discussed earlier in the literature in the context of $\nu = 1/2$ (and $\nu = 1$) systems, but no definitive idea has emerged regarding their experimental observability or their relation to the more intensively studied (and robust) $(1,1,1)$ state.2,3,5 The main purpose of the current paper is to critically discuss the possible existence of paired $\nu = 1/2$ bilayer Hall states that we, argue, are distinguishable from the better-studied $(1,1,1)$ incompressible state (as well as from compressible states) through interlayer drag experiments. The transitions between these states are described by a variety of Landau-Ginzburg theories. Given the great current interest3,13 in the physics of $\nu = 1/2$ bilayer systems, we believe that the results presented in this paper could shed considerable light on the nature of the possible quantum phase transitions in bilayer systems.

In Secs. II, III, and IV of this paper, we consider the possible bilayer $\nu = 1/2$ quantum Hall phases in the parameter regimes $d \gg l_B$ (Sec. II), $d \sim l_B$ (Sec. III), and $d \ll l_B$ (Sec. IV), where $l_B$ is the magnetic length, which sets the scale for intralayer correlations. We argue that the likely ground states in these three regimes are, respectively, compressible (Fermi-liquid-like) states ($d \gg l_B$), paired Hall states ($d \sim l_B$), and $(1,1,1)$ states ($d \ll l_B$). In Sec. V, we discuss the transition between the $d \gg l_B$ and $d \ll l_B$ limits and introduce yet another state, a hierarchical descendnet of the $(3,3,-1)$...
state with interlayer coherence. We conclude in Sec. VI with a critical discussion of the various drag resistivities that we argue can, in principle, distinguish among these phases and could be used to study bilayer quantum phase transitions experimentally.

II. \( d > l_B \): COMpressible STATE

Let us consider a bilayer system in which each layer has filling factor \( \nu = 1/2 \) in the limit of the layer separation \( d \) being much larger than the typical interparticle spacing, which is of the order of the magnetic length \( l_B \). As a starting point, we will model the system by two almost independent Fermi-liquid-like compressible states, one in each layer. There are two alternative and complementary descriptions of compressible states at even-denominator fractions. We will briefly recapitulate some features of both as they will inform the following discussion of paired states.

One description of a Fermi-liquid-like compressible state at \( \nu = 1/2 \) is based on the lowest Landau-level wave function

\[
\Psi_{\text{LLL}}(\{z_i\}) = \mathcal{P}_{\text{LLL}} \text{Det} M \prod_{i<j} (z_i - z_j)^2,
\]

where \( M \) has the matrix elements \( M_{ij} = e^{i \mathbf{k}_i \cdot \mathbf{r}_j} \). Here, \( \mathbf{r}_i \) is the position of electron \( i \). The \( \mathbf{k}_i \)s are parameters that are chosen so that the total energy of the system is minimized. We will discuss their interpretation below. \( \mathcal{P}_{\text{LLL}} \) projects states into the lowest Landau level; it has the following action:

\[
\mathcal{P}_{\text{LLL}} e^{i \mathbf{k}_i \cdot \mathbf{r}_i} \mathcal{P}_{\text{LLL}} = e^{i \mathbf{k}_i \cdot \mathbf{R}_i},
\]

where \( \mathbf{R}_i \) are the guiding-center coordinates of the electrons.

The corresponding wave function of the double-layer system at \( \nu = 1 \) can be written as

\[
\Psi_{\text{compressible}} = \Psi_{1/2}(\{z_i^\uparrow\}) \Psi_{1/2}(\{z_i^\downarrow\}),
\]

where \( \alpha = \uparrow, \downarrow \) label the two layers.

The lowest Landau-level constraint and Fermi statistics displace the electrons from the correlation holes, i.e., zeros of the wave function or, equivalently, vortices represented by the \( \Pi_{j < i} (z_i - z_j)^2 \) factor in the wave function. The \( \text{Det} M \) factor is necessary to ensure Fermi statistics; it is a displacement operator because \( \mathbf{k}_i \cdot \mathbf{r}_j \) acts as \( \vec{k}_i \cdot \mathbf{z}_j + k_i \partial / \partial z_j \) in the lowest Landau level. Thus the composite fermion made of an electron and two correlation holes has a dipolar structure. Using these ideas, the system can be described as a collection of dipolar “composite fermions” in which each dipolar fermion consists of an electron and the corresponding correlation holes. By Fermi statistics, the \( \mathbf{k}_s \)s must be distinct. The structure of a dipolar composite fermion is given schematically in Fig. 1. The energy of a composite fermion increases with \( \mathbf{k}_s \), so the ground state is a filled Fermi sea in \( \mathbf{k} \) space. In the long-wavelength limit, the total energy of these dipolar com-

![Schematic picture of the dipolar composite fermion. The black and white dots represent the electron and vortex, respectively.](image)

FIG. 1. Schematic picture of the dipolar composite fermion. The black and white dots represent the electron and vortex, respectively. The “wave vector” \( \mathbf{k}_i \) is perpendicular to the dipole moment.

posite fermions can be approximately written as \( \sum j \kappa_j^2 / 2m^* \) with effective mass \( m^* \), which is determined by the interaction potential.\(^{17,19}\)

An alternative formulation arises from the observation that an electron may be represented by a composite fermion together with a Chern-Simons gauge field that attaches two fictitious flux quanta to each fermion.\(^{20-22}\) This representation is mathematically equivalent to the original one but it naturally suggests another approximation. The composite fermions see zero average magnetic field due to the cancelation between the external magnetic field and the average fictitious magnetic field coming from the fictitious flux quanta. Consequently, the system can be described as an almost free (composite-) fermion system in zero effective magnetic field. In this approach, the fictitious flux quanta are introduced to represent the phase winding of the electron wave function around correlation holes associated with the positions of other electrons. The \( \text{Det} [ e^{i \mathbf{k}_i \cdot \mathbf{r}_j} ] \) factor in Eq. (1) can be interpreted as the wave function of the almost free composite-fermion system.\(^{14,22}\) Then \( \mathbf{k}_s \) can be regarded as the “kinetic momenta” of the composite fermions. In the long-wavelength, low-frequency limit, the two formulations are equivalent.

Note that in addition to the Fermi-liquid-like \( \nu = 1/2 \) compressible state in each layer one could have, in principle, other compressible states (e.g., charge-density wave or Wigner crystal in the \( d > l_B \) limit depending on the details of interaction and Landau-level coupling. In addition, strong disorder will lead to localized insulating states. We do not consider these possibilities in this paper.

III. \( d > l_B \): PAIRING

We will now implement the latter formulation in a double-layer system, and show that the compressible state has a strong pairing instability.\(^{12}\) We introduce two composite fermion fields \( \psi_\alpha \) and two Chern-Simons gauge fields \( \mathbf{a}_\alpha \), \( \alpha = \uparrow, \downarrow \). The Hamiltonian is

\[
H = H_0 + H_I,
\]

\[
H_0 = \int d^2r \sum_{\alpha = \uparrow, \downarrow} \frac{1}{2m^*} |\psi_\alpha^\dagger (\mathbf{\nabla} - \mathbf{a}_\alpha)^2 \psi_\alpha|,
\]

\[
H_I = \int d^2r \int d^2r' \delta \rho_\alpha (\mathbf{r}) V_{\alpha \beta} (\mathbf{r} - \mathbf{r}') \delta \rho_\beta (\mathbf{r}')
\]

with the constraints \( \mathbf{\nabla} \times \mathbf{a}_\alpha = 2 \pi \delta \rho_\alpha (\mathbf{r}) \). \( \delta \rho = 2 \) if the filling factor of each layer is \( \nu = 1/2 \). Here \( \delta \rho_\alpha (\mathbf{r}) = \rho_\alpha (\mathbf{r}) - \bar{\rho} \) is the
The interaction potential is given by \( V_{\text{int}} = V_{\text{el}} = \epsilon / r \) and \( V_{\text{int}} = V_{\text{el}} = \epsilon / r \) where \( \epsilon \) is the dielectric constant.

It is convenient to change the gauge field variables to \( a_\pm = \frac{1}{2} (a_i \pm a_j) \). Then, \( H_0 \) can be rewritten as

\[
H_0 = \int d^2 r \left[ \frac{1}{2m^*} \psi_i^* \left( \nabla - a_+ - a_- \right)^2 \psi_i + \frac{1}{2m^*} \psi_i^* \right] \times \left( \nabla - a_+ + a_- \right)^2 \psi_i
\]

with \( \nabla \times a_\pm = \pi \bar{\rho}_0 (\mathbf{r}) \pm \bar{\rho}_0 (\mathbf{r}) \). From Eq. (5), we see that \( \psi_i^* \) and \( \psi_i \) have the same gauge "charges" for \( a_- \) but opposite gauge "charges" for \( a_+ \). As a result, there will be an attractive interaction between the composite fermions in different layers via \( a_- \) and a repulsive interaction via \( a_+ \). Composite fermions in the same layer have repulsive interactions. As a result of Coulomb interactions, the attractive interaction mediated by the \( a_- \) gauge field dominates in the low-energy limit and there exists a pairing instability between the composite fermions in different layers. This result can be understood in physical terms as follows. The \( a_- \) and \( a_+ \) fields represent antisymmetric and symmetric density fluctuations. In the presence of Coulomb interactions, symmetric density fluctuations are highly suppressed but antisymmetric density fluctuations can still be large. As a result, the dynamic density fluctuations in the antisymmetric channel become more important in the low-energy limit and lead to a pairing instability.

This pairing instability has a natural explanation in the dipolar composite fermion picture.\(^{16-18}\) Let us take a dipolar composite fermion in layer \( \uparrow \) with wave vector \( \mathbf{k}^\uparrow = \mathbf{k}_F \) and a dipolar composite fermion in layer \( \downarrow \) with wave vector \( \mathbf{k}^\downarrow = -\mathbf{k}_F \). As seen in Fig. 2, this configuration can lower the interlayer Coulomb energy because the electron in layer \( \uparrow \) and the vortex in layer \( \downarrow \) can sit on top of each other and vice versa.

This analysis predicts that, at least in principle, the Fermi-liquid-like compressible state is always unstable to pairing. In practice, the pairing gap will be small in the limit \( d \gg l_B \) and easily destroyed by disorder. As a result, we expect the pairing instability discussed above to be relevant for \( d / l_B \) not too small. When \( d / l_B \) becomes small, on the other hand, the starting point of two decoupled compressible states is no longer sensible and we take a different starting point, as described in later sections.

The wave function of the paired quantum Hall state constructed in this way can be written as

\[
\Psi_\text{pair} = \Psi_\text{pair}^\text{cf} \prod_{i>j}(z_i^\downarrow - z_j^\downarrow)^2 \prod_{k>l}(z_k^\uparrow - z_l^\uparrow)^2,
\]

where

\[
\Psi_\text{pair} = \text{Pf}[f(z_i^\uparrow z_j^\downarrow \uparrow \downarrow)] = \mathcal{A}[f(z_1^\uparrow z_2^\downarrow \uparrow \downarrow, f(z_3^\downarrow z_4^\uparrow \uparrow \downarrow) \cdots],
\]

and \( f(z^\downarrow, z^\uparrow) \) is the pair wave function that depends on the symmetry of the pairing order parameter. \( \text{Pf}[\cdots] \) denotes the Pfaffian, which is defined in the second line, with \( \mathcal{A}[\cdots] \) denoting the antisymmetrized product. Notice that \( \Psi_\text{pair} \) can be regarded as the product of the wave function of the paired composite fermions and that of the (2,2,0) bosonic Laughlin quantum Hall state.

It is not immediately clear what choice of \( f(z_1^\uparrow z_2^\downarrow \uparrow \downarrow) \) is most favorable energetically. In the Chern-Simons theory of Ref. 12 there is a pairing instability in all angular momentum channels.\(^{12}\) In a modified Chern-Simons theory, it was claimed that the leading instability occurs in the \( p \)-wave channel.\(^{23}\) These approximate calculations do not necessarily capture the detailed energetics that determines the pairing symmetry. Hence, we will not enter into a discussion of energetics, but limit ourselves to a discussion of the simplest (and therefore likeliest) possibilities.

The simplest possibility is \( p_x - ip_y \) pairing,

\[
\Psi_\text{pair} = \text{Pf}[\prod_{i>j}(z_i^\uparrow z_j^\downarrow \uparrow \downarrow)] \prod_{i>j}(z_i^\downarrow - z_j^\uparrow)^2 \prod_{k>l}(z_k^\uparrow - z_l^\downarrow)^2.
\]

Using the Cauchy identity,

\[
\prod_{i>j=1}^N (a_i - a_j)(b_i - b_j) = \prod_{i,j=1}^N (a_i - b_j) \text{Det}(a_i - b_j)^{-1},
\]

this can be rewritten as

\[
\Psi_\text{pair} = \Psi_{(3,3,-1)} = \prod_{i>j} (z_i^\downarrow - z_j^\uparrow)^3 (z_i^\downarrow - z_j^\downarrow)^3 \prod_{i,j}(z_i^\downarrow - z_j^\uparrow)^{-1}.
\]

Thus \( \Psi_\text{pair} \) is the \( (3,3,-1) \) state if one takes the \( p_x - ip_y \) pairing. This wave function is well behaved in the long-distance limit but has a short-distance singularity. In the presence of Landau-level mixing, the short-distance part of the wave function can be modified without changing the structure of the wave function in the long-distance limit:

\[
\Psi_{(3,3,-1)} = \text{Pf}[h((z_i^\downarrow - z_j^\downarrow) / \xi) \prod_{i>j}(z_i^\downarrow - z_j^\downarrow)] \prod_{i>j}(z_i^\downarrow - z_j^\uparrow)^2
\]

\[
\times \prod_{k>l}(z_k^\downarrow - z_l^\uparrow)^2.
\]

Here, \( h(0) = 0 \) and \( h(x) \to 1 \) as \( x \to \infty \). In realistic systems, where Landau-level mixing is substantial, \( \Psi_{(3,3,-1)} \) could be a good candidate for the paired quantum Hall state represented by \( \Psi_\text{pair} \). It is natural to assume that \( \Psi_\text{pair} \) does not
have an “interlayer Josephson effect” because there is no gapless neutral mode in the system, in contrast to the case of the (1,1,1) state.\textsuperscript{4,13} We will show this later by direct calculation.

Another possibility for the pair wave function is an exponentially decaying function with a correlation length $\xi$, for example, either of

$$
\Psi_{SP}^{\text{PS}} = \text{Pf}\left[e^{-|z_i - z_j|/\xi} \prod_{i \neq j} (z_i^1 - z_j^1)^2 \prod_{k > l} (z_k^1 - z_l^1)^2\right],
$$

$$
\Psi_{SP}^{\text{PS}} = \text{Pf}\left[e^{-|z_i - z_j|/\xi} \prod_{i \neq j} \left(\frac{1}{z_i^1 - z_j^1} - \frac{1}{z_i^1 - z_j^2}\right)\right] \prod_{i \neq j} (z_i^1 - z_j^1)^2.
$$

(12)

This would correspond to a “strong” pairing (SP) state while the previous choice — the (3,3,−1) state — corresponds to a “weak” pairing state in the terminology of Read and Green.\textsuperscript{7,4} The two different choices of SP wave functions (with $p$- and $s$-wave pairs, respectively) in Eq. (12) can be continuously connected without crossing a phase transition.

IV. $d \leq l_B$ : (1,1,1) STATE

When the layer separation becomes sufficiently small, the interlayer Coulomb interaction can be comparable to or even larger than the intralayer Coulomb interaction. In this case, it should be more advantageous to first form an “interlayer” dipolar object that consists of an electron in one layer and two vortices in the other layer, then form a paired state of these “interlayer” composite fermions, as shown in Fig. 3. In the Chern-Simons formulation, this corresponds to the situation in which the electron in one layer can only see fictitious flux in the other layer. The appropriate Chern-Simons constraint equations are

$$
\nabla \times a_i = 2\pi \Phi \delta \rho_i, \quad \nabla \times a_i = 2\pi \Phi \delta \rho^i.
$$

(13)

As in the previous case, we can form symmetric and antisymmetric combinations of the gauge fields $a_i$ and $a^i$. Again, the antisymmetric combination mediates an attractive interaction. The wave function of the corresponding paired quantum Hall state has the following form:

$$
\Phi_{\text{pair}}^{\text{PS}} = \prod_{i \neq j} (z_i^1 - z_j^1)^2, \quad (14)
$$

where

$$
\Phi_{\text{pair}}^{\text{PS}} = \text{Pf}\left[g(z_i^1, z_j^1, z_i^2, z_j^2)\right]
$$

(15)

and $g(z_i^1, z_j^1, z_i^2, z_j^2)$ is the appropriate pair wave function. Notice that the wave function $\Phi_{\text{pair}}$ can be regarded as the product of the wave function of a paired state of the composite fermions and that of the $(0,0,2)$ quantum Hall state for bosons.

However, this line of thinking appears to conflict with the conventional wisdom that a bilayer quantum Hall system at $\nu_T = 1$ is described by the $(1,1,1)$ state for $d/l_B \sim 1.3$. Fortunately, the $(1,1,1)$ state

$$
\Psi_{(1,1,1)} = \prod_{i \neq j} (z_i^1 - z_j^1)^2 \prod_{i,j} (z_i^2 - z_j^2)^N
$$

(16)

can be rewritten in the form

$$
\Psi_{(1,1,1)} = \text{Pf}\left[g(z_i^1, z_j^1, z_i^2, z_j^2)\right] \prod_{i,j} (z_i^2 - z_j^2)^N
$$

(17)

using the Cauchy identity. In other words,

$$
g(z_i^1, z_j^1; z_i^2, z_j^2) = \frac{1}{z_i^1 - z_j^1}
$$

(18)

is the correct choice for $d/l_B \sim 1$.

Notice that this form of $g(z_i^1, z_j^1; z_i^2, z_j^2)$ corresponds to the (pseudo-)spin triplet $p_x - ip_y$ pairing order parameter for the interlayer composite fermions. From this point of view, it is natural to have the same pairing symmetry (11) for $d > l_B$ but for intralayer rather than interlayer composite fermions.

V. PHASE DIAGRAM AT $\nu_T = 1$ : CHERN-SIMONS LANDAU-GINZBURG DESCRIPTION

In this section, we consider these states within the framework of Chern-Simons effective field theories,\textsuperscript{3,10} the nature of the transitions between them, and also find an additional state that is a hierarchical descendent of the $(3,3,−1)$ state described by a $3 \times 3$ K matrix.

The Lagrangian of the $(1,1,1)$ state is

$$
\mathcal{L}^{(1,1,1)} = \Psi_\uparrow \left(i\partial_t + A_\mu^0 - a^0_\mu - a^{0\dag}_\mu\right) \Psi_\uparrow - \frac{1}{2m} \left| \nabla \frac{\Psi_\uparrow}{i} + A_\mu + a_\mu - a^{\dag}_\mu \right|^2 + \frac{1}{4\pi} \epsilon^{\mu\nu\lambda} a^\mu_\nu a^\lambda_\nu \Psi_\uparrow
$$

(19)

Here $\Psi_\uparrow$ and $A_\mu^0$ describe composite bosons and electromagnetic fields in the two layers, and the statistical gauge fields $a^\mu_\nu$ ensure the agreement between this Lagrangian and the $(1,1,1)$ wave function. In the absence of interlayer tunneling, the number of electrons in each layer is conserved, so we can use the dual description\textsuperscript{3} of the $(1,1,1)$ state,
Here $\Phi_{ij}$ describe vortices in the fields $\Psi_I$ with indices $I$ and $J$ labeling the layers $\uparrow$ and $\downarrow$, the dual gauge fields $b^I_\mu$ describe the conserved currents, the Greek indices $\mu$, $\nu$, $\lambda$ include space and time components, and the Gram (or $K$) matrix\(^8\),\(^9\) is

$$K_{IJ} = \begin{pmatrix} 1 & 1 \\ 1 & 1 \end{pmatrix}. \quad (21)$$

It is also convenient to define the charge and spin gauge fields $A_{\mu}^{C,S} = (A_{\mu} + \pm A_{\mu})/2$, $b_{\mu}^{C,S} = b_{\mu}^\uparrow \pm b_{\mu}^\downarrow$, with charge and spin vectors $t_c = (1,1)$ and $t_s = (1, -1)$, and $a = C, S$. Here and henceforth, we use the term “spin” to refer to the charge difference between the two layers, not the physical spin of the electrons, which is assumed to be fully polarized (i.e., spin here refers to an pseudospin associated with the layers).

A generic quasiparticle may now be constructed by taking a composite of $I_1$ vortices of type $\Phi_{\uparrow\uparrow}$ and $I_2$ vortices of type $\Phi_{\downarrow\downarrow}$. $\Phi_{(I_1, I_2)}$ creates such a quasiparticle, which has charge $Q$ and spin $S$.

$$Q = t_C^\uparrow K^{-1} I_1, \quad S = t_S^\uparrow K^{-1} I_1. \quad (22)$$

If the $K$ matrix has a vanishing determinant as it does in Eq. (21), then Eq. (22) will have to be modified. When this occurs, the zero eigenvalue corresponds to the Goldstone modes associated with some broken continuous symmetry. Hence, if we are to use the $K$-matrix formalism to calculate the quasiparticle properties and the degeneracy of the ground states on a torus [usually the degeneracy is det $K$ (Refs. 8 and 9)], some sort of reduced $K$ matrix will be required. We will now describe how this can be done in general. It is helpful to think in terms of the vectors in the condensate lattice;\(^8\) the $K$ matrix is the Gram matrix of the lattice, that is, the matrix of inner products of a basis of vectors in the lattice. The inner product of two vectors $m$, $n$ in the lattice, represented as column vectors of integers [not to be confused with the similar vectors $t$, which lie instead in the dual (excitation) lattice], is then given by $m^T KN$. The vanishing determinant of $K$ implies that we can find a lattice vector $n$ such that $K n = 0$. Then the inner product of $n$ with any other vector, including itself, is zero; we call $n$ a null vector. We choose $n$ to be primitive, that is, not divisible by any integer larger than 1. Two vectors that differ by an integer multiple of $n$ have the same inner product with any other vector because $n$ is null. Hence, we can obtain a reduced lattice in which we identify vectors that differ by integer multiples of $n$ and the inner product remains well defined. The reduced lattice is the quotient of the previous one by $n$. In terms of matrices, the reduced $K$-matrix is obtained by changing basis,\(^8\),\(^9\) taking $n$ as one of the basis vectors. Then in the resulting $K$-matrix, the entries in the row and column corresponding to $n$ are all zero. The reduced $K$-matrix, $K_{\text{red}}$ is obtained by deleting this row and column. The process can be repeated until either a nonzero determinant is obtained or the lattice has dimension zero; in our examples, a single reduction is sufficient.

In the case of the $(1,1)$ state above, the null vector is $n = (1, -1)$ and the reduced $K$ matrix is $K_{\text{red,}\uparrow\downarrow} = (1, 1)$, the same as in the polarized or single-layer $\nu = 1$ state. The physical meaning of the procedure is that the $(1,1)$ state is obtained by condensing composite bosons with pseudospin or by taking a pseudospin-polarized state and tilting the pseudospins into the $XY$ plane. This does not affect the quantum Hall properties of the state, which remain those of the single layer $\nu = 1$ state. The procedure above correctly accounts for disregarding the direction of the pseudospin and implies that there is a single ground state on the torus up to low-lying states associated with the broken symmetry. The use of a reduced $K$ matrix gives the quasiparticle properties; the quasiparticles carry charge $\pm 1$ and are fermions. We see that the merons (vortices in the pseudospin order parameter that carry charge $\pm 1/2$ and ill-defined statistics) are not obtained from $K_{\text{red}}$, but are confined by the logarithmic potential between them, and cannot be separated to infinity with finite energy. Usually, the different degenerate ground states can be obtained from each other by creating a quasiparticle-quasihole pair, transporting one of them around the torus and subsequently annihilating them. The nontrivial statistics (Abelian, in all cases in this paper) of the quasiparticles then require degenerate ground states. Since the merons are confined, they do not contribute to the count of ground states, and indeed dragging one around the torus produces a helical texture in the ground state, increasing the energy by order width/length; we do not regard such a state as a ground state. In general, the ground-state degeneracy is divisible by the denominator of the filling factor $\nu_f$ (equal to 1 here); any ground-state degeneracy beyond that is not exact in a finite-size system, but the energy splitting $\sim \exp(-\epsilon L)$ on a torus of size $L$, where $c$ is a constant. Finally, spin wave states have excitation energies $\sim 1/L$.

Returning to the dual Lagrangian in terms of the unreduced $K$ matrix of the $(1,1,1)$ state, in terms of the charge and spin gauge fields and the quasiparticle fields $\Phi_{(m,n)}$, it is

$$L_{d}^{(1,1,1)} = \frac{1}{2} \left[ i \partial_{\mu} - \frac{m + n}{2} \right] b^{C}_\mu \Phi_{(m,n)}^C - \frac{m - n}{2} \left[ b^{S}_\mu \Phi_{(m,n)}^S \frac{1}{2} \right] + 1 \frac{1}{2} \pi b^{C}_\mu \partial_{\nu} b^{S}_\lambda \epsilon^{\mu\nu\rho\lambda} - \frac{1}{2} \pi A^{\mu}_{\nu} \partial_{\nu} b^{a}_{\lambda} \epsilon^{\mu\nu\rho\lambda} + \frac{1}{2} (f^{a}_{\mu\nu})^2. \quad (23)$$

Since there is no Chern-Simons term for $b^{S}_\mu$, it is massless. This gauge field is dual to the Goldstone mode that results when $\Psi_{\uparrow}$, $\Psi_{\downarrow}$ condense, thereby breaking the $U(1)$ pseudospin symmetry. Quantum fluctuations can disorder the pseudospin degree of freedom. This occurs when $\Phi_{(1,-1)}$ (the field for merons) condenses in Eq. (23). The effective theory for this transition is

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We note the existence, in principle, of a state intermediate between the (1,1,1) state and the quantum disordered state. It is gapped, unlike (1,1,1), and has lower energy density. Note that again the $K$ matrix for the SP state and the reduced $K$ matrix for the SP/F state are both just $K_{ij} = (4)$. The transition between the SP and SP/F states is an XY transition at which the spin degeneracy is broken; it can also be viewed as a strong- to weak-pairing transition, similar to Ref. 24, but in the presence of pseudospin order in the pseudospin.

As we have described in the previous section, this state can also be viewed as a paired state. Passing to the dual theory, we have Eq. (20), but with

$$K_{ij} = \begin{pmatrix} 3 & -1 \\ -1 & 3 \end{pmatrix}.$$ (30)

In terms of the charge and spin gauge fields and the quasiparticle fields $\Phi_{(m,n)}$, the dual Lagrangian takes the form

$$\mathcal{L}^{(3,3,-1)} = \mathcal{L}(\Phi_{(m,n)}) + \frac{1}{4\pi} b^a_\mu \partial_\mu b_\nu^a \epsilon^{\mu \nu \rho \kappa} + \frac{2}{4\pi} b^a_\mu \partial_\mu b_\nu^b \epsilon^{\mu \nu \rho \kappa}$$

$$- \frac{1}{2\pi} A^a_\mu \partial_\mu b_\nu^b \epsilon^{\mu \nu \rho \kappa} + \frac{1}{2} (f^a_{\mu \nu})^2.$$ (31)

From Eq. (30) we can see that the charge sector of the (3,3,−1) state is similar to the SP state. It has a quantized charge Hall conductance and supports elementary excitations of charge 1/2. The pseudospin sector is different from that of the charge sector: it is gapped, unlike (1,1,1), and exhibits a pseudospin Hall effect (which is manifested in the Hall drag resistance as we discuss later), unlike (1,1,1) and SP.

The condensation of the neutral semion (1,−1) in Eq. (31) eliminates the pseudospin gauge field $b^a_\mu$ by the Anderson-Higgs effect, thereby leading to the SP state. This is analogous to the situation at $\nu_F = 2/1$, where it was shown in Ref. 24 that the transition between the (3,3,1) and strong-pairing states is a second-order transition at which a Dirac fermion becomes massless. However in the $\nu_F = 1$ case, we are dealing with a semion, rather than a fermion, so we might expect the transition to be analogous to the quantum Hall liquid to insulator transition. Both are described by a single relativistic field coupled to the Chern-Simons gauge field. In the large-$N$ limit this transition was shown to be second
order, but in the relevant $N=1$ limit the gauge-field fluctuations may drive the transition first order. However, similar arguments in the absence of a Chern-Simons term were not conclusive. That transition was argued to be first order in $4-\epsilon$ dimensions, but three-dimensional duality implies that the transition is in the inverted $\chi'$ universality class and, therefore, second order. Therefore, the possibility that the transition in the presence of the Chern-Simons term is second order appears to be still open. In the presence of disorder, at any rate, the transition will be second order. We believe, therefore, that the transition is in the inverted $\chi'$ universality class.

We may, on the other hand, consider the condensation of the boson $(2,-2)$ upon the attachment of two flux quanta,

$$L^{(3,3,-1)} = \frac{1}{2}[(i\partial_\mu - 2b^\mu - \alpha_\mu)\tilde{\Phi}(2,-2)]^2 + V(|\tilde{\Phi}(2,-2)|^2)$$

$$+ \frac{1}{4\pi} \sum_{\mu} A^\mu_\chi \partial_\mu \beta_\chi e^{\mu\nu\chi} + \frac{2}{4\pi} \sum_{\mu} b^\mu \partial_\mu \beta_\chi e^{\mu\nu\chi}$$

$$+ \frac{1}{4\pi} \sum_{\mu} b^\mu \partial_\mu \beta_\chi^C e^{\mu\nu\chi} + \frac{2}{4\pi} \sum_{\mu} b^\mu \partial_\mu \beta_\chi e^{\mu\nu\chi} - \frac{1}{2\pi} A^a_\mu \partial_\mu f^{a\mu\nu} + \frac{1}{2}(f^{a\mu\nu})^2.$$  (32)

When $\tilde{\Phi}(2,-2)$ condenses, the resulting Meissner effect enforces the condition $2b^\mu = -\alpha_\mu$ (up to gauge transformations). Hence, the following quantum Hall state results:

$$L = \frac{4}{4\pi} b^\mu_\mu \partial_\mu \beta_\chi e^{\mu\nu\chi} + \frac{2}{4\pi} \beta_\mu \partial_\mu \beta_\chi e^{\mu\nu\chi} + \frac{1}{4\pi} b^\mu \partial_\mu \beta_\chi^C e^{\mu\nu\chi}$$

$$+ \frac{2}{4\pi} b^\mu \partial_\mu \beta_\chi e^{\mu\nu\chi} - \frac{1}{2\pi} A^a_\mu \partial_\mu f^{a\mu\nu} + \frac{1}{2}(f^{a\mu\nu})^2.$$  (33)

or

$$L = \frac{1}{4\pi} K_{IJ} b^I_\mu \partial_\mu b^J_\nu e^{\mu\nu\chi} + \frac{1}{2}(f^{a\mu\nu})^2 - \frac{1}{2\pi} A^a_\mu t^I_\alpha \partial_\mu b^J_\alpha e^{\mu\nu\chi},$$  (34)

where $I,J=1,2,3$, $b^3 = -\beta_\mu$, $t^I_\chi = (1,1,0)$, $t^I_S = (1,-1,0)$, and

$$K_{IJ} = \begin{pmatrix} 3 & -1 & 2 \\ -1 & 3 & -2 \\ 2 & -2 & 2 \end{pmatrix}.  \quad (35)$$

This state is a hierarchical descendent of the $(3,3,-1)$ state (though the construction differs from the usual hierarchy by condensing neutral quasiparticles, which does not change the filling factor). A wave function for it can be constructed along the lines of Ref. 8. When the number of flux quanta on the sphere is $N_\phi = N-3$ [as in the $(3,3,-1)$ state], it contains two merons and two antimerons, which are bound in pairs to form two charge $+1$ excitations so the natural ground state has $N_\phi = N-1$. Also, it can occur for $N$ odd as well as for $N$ even, like the $(1,1,1)$ state and unlike the $(3,3,-1)$ state. The state is distinct from the $(1,1,1)$ state, despite the fact that it breaks pseudospin symmetry (since $\Phi(2,-2)$ carries pseudospin, the same value as the electron) and has a gapless Goldstone mode. The $3\times3$ $K$ matrix has determinant zero and hence a reduced $K$ matrix is required. This can be obtained most easily by first making the basis change for the condensate lattice to basis vectors $(1,0,-1)$, $(0,1,1)$, and $(0,0,1)$ (relative to the previous basis). The resulting $K$ matrix is

$$K_{IJ}' = \begin{pmatrix} 1 & 1 & 0 \\ 1 & 1 & 0 \\ 0 & 0 & 2 \end{pmatrix}.  \quad (36)$$

Since this contains the $(1,1,1)$ state as a block, it is clear that the reduced $K$ matrix is

$$K_{\text{red}}' = \begin{pmatrix} 1 & 0 \\ 0 & 2 \end{pmatrix}.  \quad (37)$$

Hence, the ground-state degeneracy on the torus is 2. Incidentally, the block containing only 2 represents a “hidden SU(2)” in this state; the corresponding edge theory is an SU(2) current algebra at level 1, even though this will presumably not be a symmetry of the Hamiltonian.

We can complete the circle and return to our starting point, the $(1,1,1)$ state, if the quasiparticle $\Phi(1,1,1)$ (where the vortices are relative to the original basis) condenses in the state (34), thereby eliminating one of the neutral gauge fields by the Anderson-Higgs effect. The proliferation of these vortices leaves intact only those condensates (composite boson fields) that do not wind on going around these vortices. These condensates lie on a sublattice (of the unreduced lattice), which is the same as that of the $(1,1,1)$ state; in fact in the basis used for $K'$ in Eq. (36) or for $K_{\text{red}}'$ in Eq. (37), the transition has destroyed the condensate described by the $1\times1$ block at the lower right, leaving the $(1,1,1)$ state. This transition could be first order or second order in the absence of disorder according to the conflicting conventional wisdom discussed above.

A more useful form for the critical theory for the transition between $(3,3,-1)$ and its interlayer coherent descendent, along the lines of Eq. (25), may be derived from Eq. (32) by making the change of variables $\alpha_\mu \to -\alpha_\mu - 2b^\mu$ and integrating out $\beta_\mu$. The Lagrangian takes the form

$$L = \frac{1}{2}[(i\partial_\mu - \alpha_\mu)\tilde{\Phi}(2,-2)]^2 + V - \frac{1}{8\pi} \alpha_\mu \partial_\mu \alpha_\chi e^{\mu\nu\chi}$$

$$+ \frac{1}{2\pi} \alpha_\mu \partial_\mu b^\chi e^{\mu\nu\chi} + \frac{1}{4\pi} b^\mu \partial_\mu b^C_\chi e^{\mu\nu\chi} + \frac{1}{2}(f^{a\mu\nu})^2 - \frac{1}{2\pi} A^a_\mu \partial_\mu f^{a\mu\nu}.$$  (38)

The gauge field $b^\mu_\chi$ only appears linearly in the Lagrangian, so we may integrate it out thereby resulting in the constraint that $\alpha_\mu = A^\chi_\mu$ up to a gauge transformation. The final Lagrangian is then
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\[ \rho_{ij}^{a\beta} = \rho_{ij}^{cf} a_{i} b_{j} + \epsilon_{ij} \rho_{ij}^{cs} a_{i\beta}, \]

where \( i, j = x, y \), \( \alpha, \beta = \uparrow, \downarrow \), \( \epsilon_{ij} \) is the antisymmetric tensor, and \( \rho_{ij}^{cs} a_{i\beta} = 2 \delta_{\alpha\beta} \) in the compressible and \((3,3,1)\) states (intralayer composite fermions) while \( \rho_{ij}^{cs} a_{i\beta} = 2 \sigma_{a\beta} \) in the \((1,1,1)\) state (interlayer composite fermions).

First, consider the longitudinal drag resistivity. In the compressible state, if we neglect gauge-field fluctuations, \( \rho_{xy}^{cf \parallel \parallel} = 0 \), so \( \rho_{xy}^{11} = 0 \). Including these fluctuations, it vanishes as \( T^{d_{c}} \) at low temperatures.\(^{30}\) In \( \Psi_{(1,1,1)} \) and \( \Psi_{(3,3,1)} \) (as well as its interlayer coherent descendant), \( \rho_{xy}^{cf \parallel} \) and, hence \( \rho_{xy}^{11} \) vanish at zero temperature and are activated at low temperatures.\(^{31}\)

Now, let us consider the Hall drag resistivity. In the compressible state, both terms on the right-hand side of Eq. (40) vanish so the Hall drag resistivity vanishes. In the \((3,3,1)\) state, \( \rho_{xy}^{cf} \) is that of a \( p_{xy} \) + \( ip_{xy} \) superconductor, which has vanishing charge resistivity (since it is a superconductor) but quantized spin Hall resistivity.\(^{24,33}\) In other words, \( \rho_{xy}^{cf \cdots} = 0 \), \( \rho_{xy}^{cf \cdots} = 0 \), \( \rho_{xy}^{00} = 1 \). Consequently, \( \rho_{xy}^{11} = \rho_{xy}^{11} = 3 \) and \( \rho_{xy}^{11} = -1 \).

In the \((1,1,1)\) state, \( \rho_{xy}^{cf} \) is identical but \( \rho_{xy}^{00} \) is different so \( \rho_{xy}^{11} = \rho_{xy}^{11} = 1 \) and \( \rho_{xy}^{11} = 1 \), in agreement with Ref. 32. The same result can be deduced physically by noting that interlayer coherence requires that the voltage be the same in both layers. If we run a current in one layer alone, then this condition can only be satisfied if \( \rho_{xy}^{11} = \rho_{xy}^{11} \). On the other hand, the Hall resistance of the system is \( \rho_{xy}^{11} + \rho_{xy}^{11} / 2 \). Since this must equal 1, we obtain the previously stated result. Note that the same logic applies to the interlayer coherent descendant of the \((3,3,1)\) state that must, therefore, have \( \rho_{xy}^{11} = \rho_{xy}^{11} = 1 \). In other words, the full resistivity tensor of the \((1,1,1)\) state is identical to that of the interlayer coherent descendant of the \((3,3,1)\) state.

Note that the interlayer coherent descendant of the \((3,3,1)\) state discussed in the previous section is a pseudospin Hall superconductor. From Eq. (39), we see that the pseudospin conductivity tensor is of the form

\[ \sigma_{ij}^{ps} = \begin{pmatrix} \kappa / i \omega & i \pi \\ -i \pi & \kappa / i \omega \end{pmatrix}, \]

where \( \kappa \) is a constant. Hence, there is nonvanishing spin Hall conductivity. However, upon inverting this tensor, we see that the spin Hall resistivity vanishes, as it must in order to satisfy \( \rho_{xy}^{11} = \rho_{xy}^{11} = 1 \). The \((1,1,1)\) state, on the other hand, has vanishing spin Hall conductivity. The distinction between a pseudospin Hall superconductor and an “ordinary” pseudospin superconductor is reminiscent of the difference between a Hall insulator and an ordinary insulator (but inverted).

Thus far, we have focused on the situation in which the layers are perfectly balanced. If they are unbalanced due to the presence of an external bias field, for example, this is analogous to introducing a pseudospin Zeeman field along

\[ \right]
the $z$ direction. This will have a pair-breaking effect on the paired states and will be expected to weaken the quantum Hall effect. This should, as a consequence, increase the longitudinal drag (along with the total longitudinal resistance). The presence of external bias can thus be used to distinguish the paired state from other incompressible states.

We conclude by summarizing our results. We have shown that the $\nu = \frac{1}{3}$ ($\nu_T = 1$) bilayer quantum Hall system (in the absence of interlayer tunneling) is likely to have as its ground state a novel paired Hall state (possibly of $p$-wave symmetry) for intermediate layer separations $d > l_B$, which gives way to the usual $(1,1,1)$ state for smaller layer separations ($d \leq l_B$), and to compressible Fermi-liquid-type states (two decoupled Halperin-Lee-Read $\nu = \frac{1}{3}$ layers) for large layer separations ($d \gg l_B$). We argue that the quantum phase transitions separating the paired states from other incompressible states, the $(3,3,-1)$ and $(1,1,1)$ states may occur via an intermediate state that is either the SP state or the interlayer coherent hierarchical descendent of the $(3,3,-1)$ state and in either case one of the two transitions will be in the $XY$ universality class.

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1 Throughout this paper, we use $\nu_T$ to denote the total filling factor and $\nu$ to denote the filling factor in each layer, $\nu_T = 2\nu$.
2 See, for example, the articles by J. P. Eisenstein, S. M. Girvin, and A. H. MacDonald, in Perspectives in Quantum Hall Effects, edited by S. Das Sarma and A. Pinczuk (Wiley, New York, 1997), and references therein.