Collective Excitations in Quantum Hall Liquid Crystals: Single-Mode Approximation Calculations

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A variety of recent experiments probing the low-temperature transport properties of quantum Hall systems have suggested an interpretation in terms of liquid crystalline mesophases dubbed quantum Hall liquid crystals. The single mode approximation (SMA) has been a useful tool for the determination of the excitation spectra of various systems such as phonons in $^3$He and in the fractional quantum Hall effect. In this paper we calculate (via the SMA) the spectrum of collective excitations in a quantum Hall liquid crystal by considering nematic, tetratic, and hexatic generalizations of Laughlin’s trial wave function having two-, four- and six-fold broken rotational symmetry, respectively. In the limit of zero wavevector $\mathbf{q}$ the dispersion of these modes is singular, with a gap that is dependent on the direction along which $\mathbf{q} = 0$ is approached for nematic and tetratic liquid crystalline states, but remains regular in the hexatic state, as permitted by the fourth order wavevector dependence of the (projected) oscillator strength and static structure factor.

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I. INTRODUCTION

For more than two decades two-dimensional electron systems (2DES), and in particular the quantum Hall effect (QHE), have been constant sources of complex and unexpected behavior, perhaps with no equal in the realms of condensed matter physics. The unique combination of extremely high mobilities ($\mu \sim 10^7$ m/Vs) in GaAs/Al$_x$Ga$_{1-x}$As heterostructures, low temperatures ($T < 100$ mk), enhancement of interactions due to the reduced dimensionality, and relative quenching of the kinetic energy in strong magnetic fields due to Landau level (LL) quantization, has allowed the emergence of interactions due to the reduced dimensionality, and relative quenching of the kinetic energy in strong magnetic fields due to Landau level (LL) quantization, has allowed the emergence of complex and striking behavior due to subtle correlation effects. In fact, from the discovery of the integer and fractional QHE’s in the early eighties (leading to the award of two Nobel prizes), to the existence of novel fractionally charged composite particles, and many newer interesting many-body phenomena, the QHE has been a consistently active and exciting area of research.

The physics of 2DES in partially filled LL’s is inherently complex due to the high degeneracy of the “unperturbed” ground state (i.e. without the Coulomb interaction). The first successful theoretical approach to this system was proposed shortly after the discovery of the FQHE by Laughlin, who proposed his famous trial-wavefunction

$$\Psi_{1/m} = \prod_{i<j}^{N} (z_i - z_j)^{m} e^{-\frac{1}{4} \sum_{i=1}^{N} |z_i|^2} , \quad (1)$$

to describe states at filling factor $\nu = 1/m$, with $m$ an odd integer. Here $z_j = x_j + iy_j$ is the position of the $j$th electron in the complex plane and we work in units of the magnetic length $l_B = [\hbar/eB]^{1/2}$. The “goodness” of Laughlin state originates in that the nodal hyper-surfaces on which the many-body wavefunction vanishes coincide with the ones in which the particles are in contact. Whereas the vanishing of the wavefunction when two (spin aligned) electrons are in contact is required by Fermi statistics, Laughlin state has multiple nodes at those points, thus reducing the Coulomb repulsion. Later, Jain proposed the elegant conceptual framework of the composite fermions (CF) which unifies all hierarchies of integer and fractional QHE’s, along with the intermediate states in between QH plateaus: an even number $2p$ of vortices is attached to each electron, also lowering the Coulomb energy, and (in a “mean field” approximation) the “magnetic fluxes” associated with the vortices lead to a reduction of the effective magnetic field: $B^* = B - 2p \Phi_0 n_e$ ($\Phi_0 = h/e$ is the flux quantum and $n_e$ is the electron density), resulting in an effective filling factor $\nu^* = \nu/(1 - 2p\nu)$. It is easy to see that for the strongest FQHE states $\nu^*$ is an integer (e.g. for $\nu = 1/3$ and $p = 1$ $\nu^* = 1$), leading to the interpretation of the FQHE of electrons as a simple IQHE of the CF’s; whereas the intermediate regions (e.g. $\nu = 1/2, 1/4$ correspond to $\nu^* \to \infty$, i.e. a vanishing effective magnetic field: $B^* = 0$.

Since 1999, magnetotransport experiments have uncovered a variety of surprising results at low temperatures ($T \lesssim 100$ mK), for example: extreme anisotropies, and apparent competition between different ordered phases in the intermediate regions between quantum Hall plateaus at high LL’s, the melting transition between a Wigner crystal (WC) at $\nu \approx 1/7,2$ and the microwave conductivity evidence of structural phase transitions in partially filled LL’s. A large body of evidence corresponding to seemingly distinct phenomena may be partly understood in terms of a simple conceptual point: that the many-body states involved have an intrinsic crystalline or liquid crystalline order,[9,10,11,12,13,15,16,17] be it smectic,[9,10,12,13] nematic,[10,11,12,13] or tetratic.[12]

Generalizations of Laughlin wavefunction [Eq. (1)] with discrete broken rotational symmetry (BRS) have been proposed in the past[12,13,18] as candidates for nematic or hexatic states[12,13] in order to understand anisotropic transport observed in the intermediate regions[9,10] or the melting of the WC at $\nu = 1/7$.[19] In fact, the motivation for these states arises from the fact that it is generally expected that melting in 2D may occur through a topological Kosterlitz-Thouless-type (KT) transition.[20,21] For 2D melting, the reliable Kosterlitz-Thouless-Halperin-Nelson-Young (KTHNY) theory predicts that, in fact, an intermediate liquid crystalline phase will exist be-
tween a solid and a liquid phase, which will exhibit no trans-
lational order, and only a quasi-long-range order for the orient-
tional order below the KT transition temperature. These arg-
ments were used by Wexler and Dorsey to calculate qual-
itatively correct anisotropic-isotropic transition temperatures
for the quantum Hall liquid crystal in the transitional regions
at high LL's. In this paper we consider the spectrum of collective excita-
tions for a family of liquid crystalline states in a partially filled
LL. These states are generated so as to satisfy the follow-
ing criteria which we consider reasonable for understanding
the dynamics: (i) states must obey Fermi statistics, i.e., they
must be odd under the exchange of any pair of electrons; (ii)
the states must be translationally invariant (far enough from
boundaries); (iii) there must be a broken rotational symme-
try belonging to the proper symmetry group (i.e., $C_2$ for a
nematic, $C_4$ for a tetratic, and $C_6$ for a hexatic; additional
symmetries are possible in principle, e.g., with a $C_{10}$ symme-
try, we have not explored such possibilities); (iv) states and
excitations must reside entirely in the LLL to avoid the large
cyclotron energy cost $\hbar \omega_c$. Note, as we will show later, that
the properties of any excited LL may be readily obtained from
the properties of the LLL.

States that satisfy the abovementioned requirements have
been proposed and studied in detail\cite{12} for filling factors
$\nu = 1/3, 1/5, and 1/7$. These are found by splitting the “ex-
tra” vortices of the Laughlin (or other CF) states around the
electron, while obeying the required symmetries:

$$\Psi_{1/(2p+1)} = \prod_{i<j} N^{2p} \prod_{\mu=1}^{p} (z_i - z_j - \alpha \mu) \times \prod_{i<j} (z_i - z_j) e^{-\frac{1}{4} \sum_{k=1}^{N} |z_k|^2}, \tag{2}$$

where the complex numbers $\alpha \mu$ are distributed in pairs of op-
posite value in the complex plane (to satisfy Fermi statistics).
For the states with the highest discrete symmetry at each fill-
ing factor we may take

$$\alpha \mu = \alpha e^{i \frac{2\pi}{2p+1} \mu}, \quad \mu \in \{1, 2, \ldots, 2p\} , \tag{3}$$

and without loss of generality $\alpha$ can be taken to be real.
The wavefunction in Eq. (2) represents a homogeneous liquid
crystalline state at $\nu = 1/(2p+1)$, is anti-symmetric, lies enti-
rely in the LLL, and is smoothly connected to the isotropic
Laughlin state for $\alpha = 0$. Figure 1 depicts the nodal distri-
bution for various states of Eq. (2). A remarkable feature of
these states is that they posses an underlying charge density
wave (CDW), but these CDWs are melted by fluctuations, and
overall the system is translationally invariant\cite{12}.

II. THE SINGLE MODE APPROXIMATION

To calculate the excitation spectrum we use the single mode
approximation (SMA)\cite{21,22}\ which reliably provides the first
moment (mean) of the energy of the excitations (for a given

wavevector $k$) that are coupled to the ground state by means of
the density operator\cite{21,22,23} The SMA was first used by Feyn-
man in 1953 to accurately calculate the spectrum of phonons
in superfluid $^4$He.\cite{24} The essence of the method originates on
the assumption that the ground state of a system of bosons
has a scarcity of long-wavelength excitations. Under those
circumstances, the variational wavefunction for an excitation
corresponding to a density-wave can be written as

$$\phi_k(r_1, \ldots, r_N) = N^{-1/2} \rho_k \psi_0(r_1, \ldots, r_N), \tag{4}$$

where $\rho_k = \sum_{j=1}^{N} e^{-i k r_j}$ is the density operator, and $\psi_0$ is
the many-body ground state (which is, in fact, unknown for
$^4$He). Note that this trial state automatically builds in the fa-
orable correlations of the ground state. The energy of this
excited state, $\Delta(k) = \langle \phi_k | H - E_0 | \phi_k \rangle / \langle \phi_k | \phi_k \rangle$, can be simply
evaluated:

$$\Delta(k) = N^{-1} \langle \psi_0 | \rho_k^\dagger [H, \rho_k] \psi_0 \rangle \equiv f(k) S(k) , \tag{5}$$

In the last term the numerator is the “oscillator strength” and
takes on the universal value $f(k) = \hbar^2 k^2 / 2m$, and $S(k)$ is
the static structure factor, which is directly measurable by
means of neutron scattering (it is here, in fact, where the He-
He correlations of the ground state are “encoded”). Using the
experimental results for $S(k)$, Feynman could calculate a spec-
trum of remarkable quality, showing the phonon-like
spectrum at small wavevectors, and a roton minimum at wave-
vectors comparable to the inverse interatomic distance\cite{24}.

III. THE SINGLE MODE APPROXIMATION IN THE
QUANTUM HALL EFFECT

The applicability of the SMA to fermion systems is also
well established for two- and three-dimensional systems in the
absence of magnetic fields, giving a good approximation for
the plasmons at long wavelength. For 2DES in presence of a
magnetic field it correctly gives the zero-wavevector mag-
etoplasmon at $\omega_c = eB/m_e$, a result that is guaranteed by
Kohn’s theorem\cite{25} which states that the dipolar excitation is
saturated by the cyclotron mode (this results in the modes of
interest—the intra-LL excitations—having quadrupolar char-
acter, i.e. with an oscillator strength $f \propto q^4$).
For excitations fully contained within a single LL, the cy-
clotron mode is not of primary interest. In 1985 Girvin, Mac-
Donald an Platzman (GMP) proposed an ingenious ansatz for
projected excited states\cite{22}.

FIG. 1: Nodal distribution for $z_{ij} \equiv z_i - z_j$ for a quantum Hall nematic at $\nu = 1/3$, tetratic at $\nu = 1/5$, and hexatic at $\nu = 1/7$.\cite{21,22,23}
The projected oscillator strength comes from the non-projected density operator:

\[
\bar{\rho}_q = \sum_{m,m'} (0,m') e^{-i q \cdot r} |0,m\rangle a_{0,m'}^\dagger a_{0,m} \tag{6}
\]

where \(|0,m\rangle\) correspond to single-particle states in the lowest LL and angular momentum \(m\), and \(a_{0,m}\) is the particle creator operator for such state. As in Feynman’s ansatz [Eq. (4)], Eq. (6) preserves the favorable correlations of the ground state. The exclusion of inter-LL excitations eliminates the problem with the saturation of the dipolar mode. The excited states have a compelling description, in first quantized form:

\[
\bar{\rho}_q \psi(z_1, \ldots, z_N) = \sum_{j=1}^N e^{-|q|^2/2} e^{-i q \cdot z_j}/2 \times \psi(z_1, \ldots, z_{j-1}, z_j - i q, z_{j+1}, \ldots, z_N) \tag{7}
\]

which corresponds to shifting each electron by \(\hat{e}_z \times q\) and superimposing these \(N\) configurations with an amplitude \(e^{-|q|^2/2}\).

Similarly to Eq. (5), the excitation spectrum can be readily obtained

\[
\bar{\Delta}_q = \frac{(2N)^{-1} \langle \psi_0 | [\bar{\rho}_q, [\bar{H}, \bar{\rho}_q]] | \psi_0 \rangle}{N^{-1} \langle \psi_0 | \bar{\rho}_q^2 | \psi_0 \rangle} \equiv \frac{\bar{f}(q)}{\bar{S}(q)} \tag{8}
\]

The projected oscillator strength comes from the non-commutation of the projected density operator with terms in the potential energy part of the Hamiltonian also projected onto the LLL:

\[
\bar{H} = \frac{1}{2} \sum_q \nu_q (\bar{\rho}_q^2 - N e^{-|q|^2/2}) \tag{9}
\]

Since \([\hat{\rho}_k, \bar{\rho}_q] = (e^{k \cdot q/2} - e^{k \cdot q^*/2})\bar{\rho}_{k+q}\), we find:

\[
\bar{f}(q) = 2 e^{-|q|^2/2} \sum_k \sin^2\left(\frac{k \cdot q}{2}\right) S(k) \times \delta(r_i - r_{i'}) \tag{10}
\]

For its part, the projected static structure factor \(\bar{S}(q)\) can be calculated from:

\[
\bar{S}(q) = S(q) - (1 - e^{-|q|^2/2}) \tag{11}
\]

where \(S(q)\) is the unprojected static structure factor:

\[
S(q) - 1 = \rho_0 \int d^2 r e^{-i q \cdot r} [g(r) - 1] \tag{12}
\]

the Fourier transform of the pair correlation function

\[
g(r) = n_e^{-2} \left( \sum_{i \neq j} \delta(r_i - r_j) \right) \tag{13}
\]

which is obtainable from the ground state via, e.g., Monte Carlo (MC) simulations. In our case, we considered BRS states [Eq. (3)] corresponding to a \(\nu = 1/3\) nematic, a \(\nu = 1/5\) tetratic, and a \(\nu = 1/7\) hexatic. Figure 2 depicts the pair correlation function for various states. In all cases, correlation functions and SMA excitation spectra were computed for numerous \(\alpha\)'s. The angle dependence significantly increases the burden in the MC simulations since the full angle-dependent \(g(r)\) is needed, rather than the considerably simpler angle-averaged \(g(r)\) of the isotropic systems. The accurate calculation of \(\bar{f}(q)\), with its angular-dependent exponentially large factors required high quality \(\bar{S}(q)\) and hence \(g(r)\). To put things in perspective, to achieve ca. 1% accuracy in \(g(r)\), \(O(10,000)\) counts were accumulated for each 0.01 \times 0.01 (in units of \(l_0\)) r-box in the original histograms (the small boxes were necessary to have precision in the Fourier transform at high wavevectors). This required, for each filling factor \(\nu\) and anisotropy generating \(\alpha\), to run approximately 4–8 \times 10^7 MC steps in systems of \(N_e = 200–400\) electrons taking 100–200 cpu\times days of computation for each \((\nu, \alpha)\) in a 2 GHz Athlon cluster (see Ref. 13 for further details). A relatively large \(N_e\) is required so that the simulations are able to reproduce a system in the thermodynamic limit. The fact that \(g(r) \simeq 1\) over a large area is a guarantee of that achievement (Fig. 2).

Figure 3 shows the projected static structure factors \(\bar{S}(q)\) for a \(\nu = 1/3\) nematic, a \(\nu = 1/5\) tetratic, and a \(\nu = 1/7\) hexatic. From \(\bar{S}(q)\), the oscillator strength \(\bar{f}(q)\) is computed using Eq. (11). Analysis of \(\bar{S}(q)\) and \(\bar{f}(q)\) shows that both are \(O(q^2)\) for small \(q\). This restriction on the the small wavevector behavior originates in Kohn’s theorem as all of the \(O(q^2)\)
FIG. 3: (Color online) Projected static structure factors $\tilde{S}(q)$ for a $\nu = 1/3$ nematic, a $\nu = 1/5$ tetratic, and a $\nu = 1/7$ hexatic. In all cases, $\tilde{S}(q) = O[q^4]$ for small $q$, and $\tilde{S}(\infty) \rightarrow 0$.

pieces in the unprojected parts are saturated by the (uninteresting) inter-LL excitations at $\hbar\omega_c$.

Figure 4 presents some of our results for the excitation spectra in the lowest LL. The results are consistent with those obtained by GMP for the isotropic $\nu = 1/3$ and 1/5 FQHE cases, which were qualitatively confirmed experimentally. They show that the collective excitation spectrum remains gapped, albeit with a deep magnetoroton, for modes coupled to the ground state via the density operator. Not surprisingly, BRS states have significant anisotropy in their spectra. However, it is noteworthy that for the nematic and tetratic cases the spectrum is singular, with an angle dependence on the excitation energy $\Delta(q)$ as $q \rightarrow 0$. By contrast, the hexatic liquid crystal has a regular spectrum in the long wavelength limit. The apparent disparity have, of course, to do with the different rotational symmetries of the different states: as $\Delta = f/S$, and both numerator and denominator are $O[q^4]$ for small $q$, there is no possibility of generating a $C_6$ symmetric form with terms that can only depend on $q_0^2$, $q_0^4$, and $q_0^2q_0^4$. This can also be understood from the point of view of an effective elasticity theory for 2DES (valid in the long wavelength limit, $q \rightarrow 0$)\textsuperscript{27}, the elasticity tensor being of the fourth rank is not compatible with a six-fold rotational symmetry. This elastic interpretation will be published elsewhere.\textsuperscript{17}

The presence of this singular spectrum suggests that microwave conductivity experiments\textsuperscript{28} may be able to discern such structures as a signature of, e.g., the quantum Hall nematic suggested by magnetotransport experiments.\textsuperscript{28}

Another interesting feature (Fig. 5) is that the magnetoroton minima and the gap at the origin appears to collapse at some finite anisotropy factor $\alpha$. The appearance of elastic modes that go soft may be a precursor of the appearance of charge density waves in the system.\textsuperscript{28} Spectra in higher LL’s can be obtained by multiplying the projected density operator $\tilde{\rho}_L$ by $L_L(q_0^2/2)$, where $L_L(x)$ is the Laguerre polynomial and $L$ corresponds to the desired LL.\textsuperscript{28} The rightmost panel of Fig. 4 shows the modification of the SMA spectrum for the first excited LL. Numerical error due to the large $L_L(q_0^2/2)$ factors makes it difficult to get reliable results for higher LL’s at this point.

Theoretical predictions of which type of order is lowest in free energy (required to decide which state is favorable at finite temperatures, as shown in Refs. \textsuperscript{6,7}) are still incomplete, however, as the entropy is likely to be dominated by gapless modes originating, e.g., from the Goldstone modes associated with the spontaneous breaking of the continuous rotational symmetry of the isotropic states. Generalization of the SMA to gapless rotational modes (see e.g. \textsuperscript{15}) will require, however, 3-body operators which demanding considerably higher computing capabilities.\textsuperscript{28}

IV. SUMMARY

Summarizing: in connection with the recent variety of experimental evidence supporting liquid crystalline phases in quantum Hall systems\textsuperscript{5,6,7,8}, we have calculated the collective excitation spectrum in the SMA approximation for a quantum Hall liquid crystal. We found that the spectrum of excitations coupled to the ground state by the density operator remains gapped, but develops a significant anisotropy, which in the case of the nematic and tetratic liquid crystals has a singular gap in the long wavelength limit.

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dramatic angular dependence of the spectra, and the appearance of a singular gap as $q \to 0$ for the nematic and tetratic. The lower right panel shows also the spectrum for the $\nu = 1/3$ nematic in the first excited LL ($L = 1$).

FIG. 4: (Color online) Single mode approximation spectra for $\nu = 1/3$ nematic, $\nu = 1/5$ tetratic, and $\nu = 1/7$ hexatic. Note the dramatic angular dependence of the spectra, and the appearance of a singular gap as $q \to 0$ for the nematic and tetratic. The lower right panel shows also the spectrum for the $\nu = 1/3$ nematic in the first excited LL ($L = 1$).

FIG. 5: Possible collapse of the magneto-roton minimum as function of the anisotropy parameter $\alpha$. Dotted lines are guides to the eye.

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