Planar arrays of interacting lines or stripes have become a paradigm for many different physical systems, including domain walls in magnets, layered superconductors, biophysical systems, liquid crystals, and charge density waves (CDW). A unidirectional CDW or the stripe phase was also predicted to form in GaAs two-dimensional (2D) electron systems when the occupation of the third or a higher Landau level is close to \( \frac{3}{2} \). This prediction was attested by a recent discovery of a dramatic magnetotransport anisotropy in the indicated range of filling factors. The easy (low resistance) and hard (high resistance) current directions are thought to be along and across the stripes, respectively. In order to have a better foundation for the magnetotransport studies, one has to first understand the intrinsic dynamics of the stripe phase, without intervening disorder effects. Recent investigations revolved around the instabilities of the stripes against the \( 2\pi \) translational order, which would transform the system into a highly anisotropic Wigner crystal.

According to MacDonald and Fisher, such an instability appears only at temperatures \( T \) below 1 mK. Here we would like to address another interesting regime of relatively high \( T \) where the \( 2k_F \) modulation is absent and moreover, the sole periodic modulation is due to the main CDW harmonic with wave vector \( k_0 \). In classical terms, the electron density at the topmost Landau level is of the form \( n(r,t) = \text{Re} \Psi(r,t) \), where \( \Psi = |\Psi| e^{i\varphi(x-u)} \) is the CDW order parameter. (We chose the \( x \)-direction to be perpendicular to the stripes). Although oversimplified, this classical picture correctly identifies the low-frequency long-wavelength degrees of freedom: the coarse-grained component \( n \) of the total density, and the CDW phase (or the displacement field) \( u \). It is natural to expect that at sufficiently small \( \omega \) and \( q \) the dynamics of \( u \) and \( n \) is governed by a certain hydrodynamic theory. The formulation of such a theory is the subject of this Letter.

We begin by writing down an effective Hamiltonian, from which we derive the hydrodynamic equations of motion, the spectrum of the collective modes, and the structure factor. Next, we study the renormalization of the bare parameters of the theory by thermal fluctuations. Our results can be verified by measuring the transmission of the micro- and surface acoustic waves as well as inelastic light scattering techniques provided the samples are of high enough mobility so that the disorder effects, which we ignore here, are not important.

**Effective Hamiltonian.**— Our first step is to construct an effective Hamiltonian for \( n \) and \( u \), which is (a) rotationally and translationally invariant and (b) reflects the specific properties of the system, such as the long-range Coulomb interaction. Very useful in this process is the similarity to smectic liquid crystals. The resultant form of the effective Hamiltonian is as follows:

\[
H = \frac{1}{2} \int d^2r [Y E^2(u) + E(u) \Delta + K (\nabla^2 u)^2 + \chi^{-1} \delta n^2 + \Phi n + 2CE(u) \delta p + \rho \omega^2],
\]

where \( E(u) = \nabla_x u - \frac{1}{2} (\nabla u)^2 \) in the usual rotationally invariant strain, \( Y \) is the compression and bending elastic moduli, \( \chi \) is the compressibility, \( \Delta \) is an auxiliary parameter (counter-term), needed to guarantee the condition \( \langle \nabla_x u \rangle = 0 \), \( \rho = mn \) is the mass density, \( \delta p = \rho - \rho_0 \) its deviation from the equilibrium value \( \rho_0 = m n_0 \). \( \Phi(r) = \int d^2r'/\delta n(r')U(r-r') \), is the electrostatic (Hartree) potential, with \( U(r) = e^2/kr \) at large \( r \) (see more details in [1]).

The last term in Eq. (1) accounts for the dependence of the CDW periodicity on \( n_0 \), with \( C = (Y/m) \partial \ln n_0 / \partial n_0 \). Finally, the last term is the correction to the kinetic energy, with \( v \) being the velocity of the electron fluid. This term is mainly a bookkeeping device: it vanishes after the projection on a single (topmost) Landau level but enables us to derive the equations of motion via the standard Poisson-bracket method.

**Equations of motion.**— The hydrodynamic fluctuations of our system are governed by the equations

\[
\begin{align*}
\frac{\partial n}{\partial t} &= \mu \nabla \cdot u, \\
\frac{\partial u}{\partial t} &= \frac{\partial}{\partial x} \left[ \frac{1}{Y} \left( \frac{\partial \Phi}{\partial n} \right) \frac{\partial u}{\partial x} \right] - \frac{1}{k_0^2} \delta n - \frac{2 Y E(u)}{k_0^2 \chi}, \\
\frac{\partial \delta n}{\partial t} &= \frac{1}{2} \left( \frac{\partial E(u)}{\partial n} \right) \frac{\partial u}{\partial x} - \frac{\partial}{\partial x} \left[ \frac{1}{Y} \left( \frac{\partial \Phi}{\partial n} \right) \frac{\partial u}{\partial x} \right],
\end{align*}
\]

where \( \mu = m n_0 \) is the mobility, and \( \delta n \) is the deviation of the density from its average value. The above equations are the hydrodynamic equations of motion for a two-dimensional electron liquid in the quantum Hall regime.
\[ F_u = \rho_0 \dot{W} + (\mathbf{v} \nabla)W - \lambda_p |\nabla W|^2, \quad \zeta_u = \zeta, \quad (2) \]

\[ F_j = \frac{\partial (\rho v_j)}{\partial t} + \nabla_j (\rho v_j v_j - \rho \omega_j \varepsilon_{jk} v_k + \nabla_j W \frac{\delta H}{\delta u}) + L_{jk}(i \nabla)v_k + \rho \nabla_k \left( \chi^{-1} \delta \rho + \Phi + CE \right) = \zeta_j, \quad (3) \]

and the continuity equation \( \dot{\rho} + \nabla (\rho \mathbf{v}) = 0 \). The notations here are as follows: \( W = x - u, \ \varepsilon_{jk} \) is the unit antisymmetric tensor, and \( \omega = eB/mc \) is the cyclotron frequency. Functions \( L_{jk}(q) \) describe dissipation. As in 3D case \( [3] \), they can be parametrized by viscosities \( \eta_i \),

\[ L_{jk}(q) = \eta q_j q_k + \delta_{jk} \eta a q^2 + \eta x \delta_{jk}(q^2 + (\eta_i - 4 \eta_3 + \eta_4 - 2 \eta_5) q^2) + \eta x (\eta_i - \eta_4) q x (q_j \delta_{kx} + q_k \delta_{jx}). \quad (4) \]

(There is no analog of \( \eta_2 \) in 2D). One more dissipative coefficient, \( \lambda_p \), [Eq. (3)] describes the permeation, i.e., the mass transport across the stripes. At relatively high \( T \), such that \( |\Psi| \leq n_0, \ \lambda_p \sim (\eta q^2 / \rho_0 a \omega^2)(n_0 / |\Psi|)^2. \) Finally, \( \zeta \)’s in Eqs. \( 3 \) and \( 4 \) are the Langevin noises, which satisfy the fluctuation-dissipation theorem,

\[ \langle \zeta_u(1) \zeta_k(2) \rangle = 0, \quad \langle \zeta_u(1) \zeta_k(2) \rangle = 2k_B T \lambda_p \rho \delta(1 - 2), \quad \langle \zeta_j(1) \zeta_k(2) \rangle = 2k_B T L_{jk}(q) \delta(1 - 2). \quad (5) \]

The main difference of our model from the conventional smectics is the presence of the strong magnetic field. As we will see below, it drastically changes both the linearized and nonlinear dynamics of the system. **Field-theory description.**— Our ultimate goal is to calculate various correlation functions, e.g., the dynamical structure factor,

\[ S(q, \omega) = (N_0 m^2)^{-1} \int dte^{i\omega t} \langle \rho_1(t) \rho_2(0) \rangle, \quad (8) \]

where \( N_0 \) is the total number of electrons at the topmost Landau level, and \( \langle \ldots \rangle \) stands for the thermal averaging, or equivalently, the averaging over the Langevin noise. The latter can be done in a systematic way by means of the Martin-Siggia-Rose (MSR) formalism \([3]\), whose field-theoretic version was previously applied to the conventional smectics by Kats and Lebedev \([2]\).

The implementation of the MSR method begins with enforcing the equations of motion by \( \delta \)-function-type weight factors, \( \delta (F_a - \zeta_a) \), in the path integral over \( F_a \). These \( \delta \)-functions are then represented by integrals of \( \text{exp}[ip_a(F_a - \zeta_a)] \) over auxiliary dynamical variables \( p_a \), which enables one to average over the Gaussian noise \( \zeta_a \). Finally, one changes the integration variables from \( F_a \) to \( V \equiv \rho v / \rho_0 \) and \( u \). If the denote by \( F^n \) the nonlinear terms in \( F_a \), and introduce vectors \( \phi^n = \{ V^1, P^1 \}, \quad V^1 = \{ u, V_x, V_y \}, \quad P^1 = \{ p_u, p_x, p_y \}, \) then the resultant action becomes

\[ A = -i \int dt \int d^2 p_a \int d^2 F_a^n(u, v) - \ln J + A_0, \quad (9) \]

\[ A_0 = \frac{1}{2} \int \frac{d^2 q}{(2\pi)^2} \int \frac{d^2 \omega}{2\pi^2} \tilde{G}^{-1}(q, \omega) \phi_q \cdot \phi_{q\omega}, \quad (10) \]

where \( J = \text{det} \frac{\partial F_a}{\partial V_b} \) is the Jacobian and \( G_0 \) is the (bare) propagator of the following block-matrix form

\[ G_0 = \begin{bmatrix} G_{VV} & -iG_{VP} \\ -iG_{VP} & 0 \end{bmatrix}. \quad (11) \]

Two particular components of \( G_0 \), \( G_{uu} \) and \( G_{up} \),

\[ G_{uu}(q, \omega) = \frac{-i s^2}{Q(q) s^2 - i \omega \tilde{v}(q) - \omega^2 \omega_x^2}, \quad (12) \]

\[ G_{up}(q, \omega) = \frac{k_B T}{\rho_0 \omega} [G_{up}(q, \omega) - G_{up}(q, -\omega)], \quad (13) \]

will play an important role in the later discussion. Here we introduced the notations \( Q(q) = (q_x^2 + K q_y)^2 / \rho_0 \), \( \omega_x^2(q) = n_0 \tilde{v}_s^2 / U(q) + \chi^{-1} / m \), \( \tilde{v}_i = \eta_i q^2 / \rho_0, \quad \epsilon = q_x / q = \cos \theta, \quad s = q_y / q = \sin \theta, \) and

\[ \tilde{v}(q) = \tilde{v}_3 + c^2 s^2 (\tilde{v}_1 - 4 \tilde{v}_3 + \tilde{v}_4 - 4 \tilde{v}_5). \quad (14) \]

Equations \( (12) \) and \( (13) \) are obtained by neglecting \( \omega, \lambda_p, \tilde{v}_i, \) and \( C \) compared to quantities proportional to “large” frequencies \( \omega_c \) and \( \omega_p \).

**Harmonic theory.**— If only the quadratic part \( A_0 \) of the full action is retained, then the correlators of the velocity fields \( V_a \) are given by the components \( G_{VV} \), \( G_{Vb} \) of \( G_0 \) and can be found after straightforward albeit tedious algebra. The continuity equation enables us to relate these correlators to the structure factor. This way we get

\[ S(q, \omega) = \frac{2}{m \omega} \text{Im} \frac{k_B T q^2 (Q s^2 - i \omega \tilde{v})}{Q_{\omega}^2 (s^2 + \lambda_p \tilde{v}) - i \omega \omega_x^2 - \omega_c^2 \omega^2}, \quad (15) \]

for \( \omega \ll \omega_c \) and, in particular,

\[ S(q, 0) = 2k_B T \frac{q^2 \omega_c \lambda_p + c^2 \tilde{v}_3 + s^2 \tilde{v}_4}{m \omega^2}, \quad (16) \]

The poles of \( S(q, \omega) \) correspond to the collective modes. Two of them (magnetophonons) are gapless:

\[ \omega_m(q) \approx \sin 2 \theta \left( \frac{V}{\rho_0} \right)^{1/2} \omega_p(q) q - \frac{i}{2} \sigma(q) \quad (17) \]

and its counterpart \( \omega = -\omega_{m}^*(q) \). Their attenuation rate \( \alpha(q) \) is given by

\[ \alpha(q) = \lambda_p Q + \omega_c^{-2} [\omega^2 p(q) \tilde{v} + Q (\tilde{v}_3 s^2 + \tilde{v}_4 s^2)], \quad (18) \]

up to terms proportional to \( C \) and higher powers of \( \tilde{v}_i \) and \( \lambda_p \). The other two collective modes (magnetoplasmsons)
have a large gap \( \omega_c \). Thus, the number of the hydrodynamic (gapless) modes coincides with the number of the hydrodynamic variables \((n \text{ and } \alpha)\) as it should [14].

For Coulomb interaction \( \omega_p(q) \propto q^{1/2} \), hence, within the harmonic theory the magnetophonons have \( \omega_m(q) \propto q^{3/2} \) dispersion, similar to that of the Wigner crystal [15] but with the \( \theta \)-dependent prefactor. Propagating magnetophonon modes exist as long as \( \theta \) is not too small so that \( \text{Re} \, \omega_m(q) \gg \alpha(q) \) in Eq. (17); otherwise, they are replaced by the two overdamped modes:

\[
\omega_{\text{fast}}(q) \approx -i \alpha(q), \quad \omega_{\text{slow}}(q) \approx -i \frac{d}{\alpha(q)} \left( \frac{q_x q_y \omega_p}{q^2 \omega_c} \right)^2 \frac{Y}{p_0}.
\]

Below we will see that at small enough \( q \) the magnetophonon dispersion relation and their damping are significantly modified by the anharmonisms.

**Renormalization.** Our next step is to calculate the propagator \( G = (G_0^{-1} + \Sigma)^{-1} \) of the full theory, Eq. (3), treating the previously ignored nonlinear terms as perturbations. We expect that \( G \) has the same form as \( G_0 \), but with \( Y, K \), and other parameters replaced by \( \omega \) and \( q \)-dependent (renormalized) values \( Y^R(q, \omega), K^R(q, \omega), \) etc. The perturbative corrections to \( Y, K, \) and \( C \) are determined by the energy-component \( \Sigma_{ap}(q, \omega) \). Similarly, the corrections to \( \nu \)'s are determined by \( \Sigma_{pq}(q, \omega) \). To the lowest order in \( T \) they are given by the diagrams shown in Fig. 1.

![Diagram](image)

\[
g(l) = g_0 + \frac{g_0 - g_0}{D}, \quad D \equiv 1 + \frac{g_0}{g_0} [e^{(3-d)l} - 1], \quad \nu^2(l) = e^{2dl} \left[ \nu^2(0) + \frac{1}{2(2d-1)} \frac{\omega^2}{\omega^2_0(\Lambda-l)} \right] D \nu^2, \quad \gamma_Y = 2 \frac{d-1}{d+2}, \quad \gamma_K = 2 \frac{d}{d+2}.
\]

As one can see, for small \( \epsilon = 3-d \), \( g \) flows to a weak coupling fixed point \( g_0 = 16(3-d)(d-1)/(d+2) \), justifying our one-loop RG. The renormalized values of the parameters of the harmonic theory are found by rescaling back to the original coordinates, \( Y^R(q, \omega) = Y(l) e^{-dl}, K^R(q, \omega) = K(l) e^{-(d-2)l}, \nu^2(q, \omega) = \nu_0(l) e^{-dl}, \) etc., where \( l \) is the smallest of the three cutoffs \( l_s, l_l, l_\omega \) to be found from equations \( K(l_s) \Lambda^4 \sim \omega^2_\omega e^{2dl} Y(l_s), \Lambda = q_\perp l_s, \) and \( \omega^2_c \omega = \omega^2_\omega (\Lambda e^{-l_s}) \nu_0(l_\omega) \Lambda^2 e^{-(d+2)l_\omega}. \) These

![FIG. 1. Self-energy diagrams.](image)

Doing the power counting with the help of Eqs. (12), (13), and the relation \( k_x \sim k^2(K/Y)^{1/2} \), where \( k \) is the loop momentum, we quickly discover that \( \Sigma_{apb} \) and \( \Sigma_{apb} \) are infrared-divergent. To resolve this problem, we utilize the renormalization group (RG) procedure formulated in \( d = 3 - \epsilon \) spatial dimensions \( \text{one } x\text{-dimension plus } (d-1) \text{ “} \perp \text{”-dimensions}. \) Each step of our RG transformation consists of integrating out the cylindrical shell \( \Delta e^{-l} < k < \Lambda \) in the loop diagrams of Fig. 1 followed by rescaling of the momenta \( k = k/e^{-l} \) to restore the ultraviolet cutoff \( \Lambda \sim g_0 \). It is convenient to rescale the \( u \)-field as well \( u(r) = u'(r') e^{l} \) to preserve the structure of \( E(u) \). We did not find it necessary to rescale the frequencies or impose a cutoff on \( k_x \) [10].

The one-loop RG equations are given by

\[
\begin{align*}
\frac{d}{dt} Y &= dY - \frac{g}{8} Y, \quad \frac{d}{dt} C = dC - \frac{g}{8} C, \\
\frac{d}{dt} K &= (d-2)K + \frac{g}{8(d-1)} K, \\
\frac{d}{dt} g &= (3-d)g - \frac{2d}{16(d-1)} g^2, \\
\frac{d}{dt} \nu^2 &= 2d \nu^2 + \frac{g}{2(d-1)} \times \left[ \frac{\omega^2}{\omega^2_0(\Lambda-l)} \right] K(l) e^{(d+2)l} + \frac{\nu^2}{4}.
\end{align*}
\]

where \( \nu \equiv k_B T Y^{1/2} K^{-3/2} \Lambda^{d-3/S_{d-1}} \) is the dimensionless coupling constant, with \( S_{d-1} = (4\pi/d) \Gamma(d-1)/2 \). Equations for other viscosities are similar to (22) and are not shown.

The solutions of the RG equations [19, 23] are as follows:

\[
\begin{align*}
g(l) &= g_0 + \frac{g_0 - g_0}{D}, \quad D \equiv 1 + \frac{g_0}{g_0} \left[ e^{(3-d)l} - 1 \right], \\
Y(l) &= Y_0 e^{dl} D^{-\gamma_Y}, \quad C(l) = C_0 e^{dl} D^{-\gamma_Y}, \quad \gamma_Y = 2 \frac{d-1}{d+2}, \\
K(l) &= K_0 e^{(d-2)l} D^{\gamma_K}, \quad \gamma_K = 2 \frac{d}{d+2}.
\end{align*}
\]

Equations give rise to the characteristic crossover length and frequency scales \( \xi_y = (K_0/\nu)^{1/2}/(k_B T)^2, \xi_y = (K_0/\nu)^{1/2}/k_B T, \) and \( \omega_0 = (K_0/p_0)^{1/2} \omega_p (\xi_y^{-1}/\omega_0 \xi_y)^{1/2} \).

Taking Eq. (24) for the face value, we obtain \( \gamma_K = \gamma_K = \gamma = 1/2 \) in 2D \( (\epsilon = 1) \) and discover the following three types of asymptotic behavior:

\[
\begin{align*}
\nu^R \sim Y_0 (\xi_y q_0)^{1/2}, \quad K^R \sim K_0 (\xi_y q_0)^{-1/2}, \\
\nu^K \sim \frac{K_0}{\rho_0 \xi_y q_0^{3/4} (\xi_y q_0)^{3/4}}, \quad \nu^{R_k} \sim \frac{\alpha Y_0}{\rho_0 \xi_y q_0^{3/4} (\xi_y q_0)^{3/4}}.
\end{align*}
\]
for \( q_z \ll \xi_x^{-1}(q_y\xi_y)^{3/2} \), \( q_y \ll \xi_y^{-1} \), \( \omega \ll \omega_y(q_y\xi_y)^{9/4} \),

\[
Y^{R} \sim \frac{Y_{0}(\xi_x q_x)^{1/3}}{\rho_{0} \omega \xi_y q_y^{1/2}}, \quad K^{R} \sim \frac{K_{0}(\xi_x q_x)^{-1/3}}{\rho_{0} \omega \xi_y^{5/3} q_y^{3}}, \quad \nu_{3}^{R} \sim \frac{a_{1} Y_{0}(\xi_x q_x)^{1/3}}{\rho_{0} \omega \xi_y^{5/3} q_y^{3}} \quad (26)
\]

for \( q_z \ll \xi_x^{-1} \), \( q_y \ll \xi_y^{-1} q_x^{2}/\xi_x^{3/2} \), \( \omega \ll \omega_y(q_x\xi_x)^{3/2} \), and

\[
\nu_{3}^{R} \sim \frac{K_{0}}{\rho_{0} \omega y^{2/3} \xi_y^{-1/3}} \quad , \quad \nu_{4}^{R} \sim \frac{a_{1} Y_{0}}{\rho_{0} \omega y^{2/3} \xi_y^{-1/3}} \quad (27)
\]

for \( \omega \gg \omega_y(q_y\xi_y)^{9/4} \), \( \omega \gg \omega_y(q_x\xi_x)^{3/2} \).

Here \( a_{1} = (1 - r)^{2}, \quad a_{4} = r^{2}, \quad a_{5} = (1 - r)r, \quad \) and \( r = \rho_{0} C_{0}/Y_{0} \). Thus, \( a_{1} a_{4} = a_{5}^{2} \), which entails the relation \([\Omega] \Delta \eta_{1} \Delta \eta_{4} = (\Delta \eta_{3})^{2} \) derived earlier for the conventional smectics.

**Discussion.** Our one-loop-level results for \( Y^{R} \) and \( K^{R} \) turn out to be exact \([17]\) for the static limit \([2] \) and \([20] \). It is possible that this success transcends to the dynamics, in which case the scaling exponents for viscosities in Eqs. \((23) \) \((27) \) would be exact as well.

Within the domain of anomalous hydrodynamics summarized by Eqs. \((23) \) \((27) \), the magnetophonon dispersion relation becomes

\[
\text{Re} \omega_{m}(q) \sim s c^{7/6}(\xi_{x} q)^{5/3} \frac{\omega_{p}^{2}}{\omega_{x} \xi} \frac{Y_{0}}{\rho_{0}} \quad (28)
\]

and

\[
\text{Im} \omega_{m}(q) \sim s c^{5/6}(\xi_{x} q)^{11/6}(1 - 2r)^{2} \frac{\omega_{y}^{2}}{\xi_{y}} \frac{K_{0}}{Y_{0}} \quad (29)
\]

for a finite fixed \( r \) and \( q \ll \xi_{x}^{-1} \). In contrast to the case of conventional smectics \([11] \) \([3] \), the scaling of \( \text{Im} \omega_{m}(q) \) is cut off by \( t_{c} \) \([24] \) not \( t_{\omega} \) \([27] \).

Strictly speaking, the smectic behavior in 2D can persist only up to the length scale smaller than the average distance \( \xi_{z} \sim \exp(E_{d}/2k_{B}T) \) between dislocations, \( E_{d} \sim K_{0} \) being the dislocation energy \([18] \). The anomalous hydrodynamics can be observed if \( \xi_{x} < \xi_{d} \), which is satisfied at low \( T \) where \( \xi_{d} \) is exponentially large \([18] \).

In conclusion, we formulated a novel long-wavelength low-frequency effective theory of the stripe phase arising in a two-dimensional electron liquid near half-integral fillings of high Landau levels. Our theory applies at relatively high temperatures and in the clean limit. We demonstrated that the collective mode properties of the system exhibit nontrivial power-law scaling, which can be verified by microwave or surface acoustic wave measurements at a finite wave vector or by inelastic light scattering experiments. Our future plans include applying the proposed hydrodynamic approach to the problem of the magnetotransport.

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