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Composite Fermi Liquid at Zero Magnetic Field in Twisted $MoTe_2$

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The pursuit of exotic phases of matter outside of the extreme conditions of a quantizing magnetic field is a long-standing quest of solid state physics. Recent experiments have observed spontaneous valley polarization and fractional Chern insulators in zero magnetic field in twisted bilayers of $MoTe_2$, at partial filling of the topological valence band ($\nu = -2/3$ and $-3/5$). We study the topological valence band at *half* filling, using exact diagonalization and density matrix renormalization group calculations. We discover a composite Fermi liquid (CFL) phase even at zero magnetic field that covers a large portion of the phase diagram near twist angle ~3.6°. The CFL is a non-Fermi liquid phase with metallic behavior despite the absence of Landau quasiparticles. We discuss experimental implications including the competition between the CFL and a Fermi liquid, which can be tuned with a displacement field. The topological valence band has excellent quantum geometry over a wide range of twist angles and a small bandwidth that is, remarkably, reduced by interactions. These key properties stabilize the exotic zero field quantum Hall phases. Finally, we present an optical signature involving "extinguished" optical responses that detects Chern bands with ideal quantum geometry.

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Strong interactions can lead to exotic phases of matter such as non-Fermi liquids. A remarkable example is the composite Fermi liquid (CFL) that occurs in a half or quarter filled lowest Landau level (LLL). The CFL is a non-Fermi liquid with an emergent Fermi sea composed of charge neutral "composite fermions" [1–4] and has anomalous responses to a wide variety of experimental probes [5–10]. The gapless CFL state has provided an elegant interpretation for various Abelian [1–4] and non-Abelian gapped topological phases [11].

This work proposes an alternative route to realize CFLs. Our proposal is based on twisted 2D transition metal dichalcogenides (TMD), a family of platforms that have realized a wealth of interesting phenomena [12–27], and generated much theoretical interest for their topological properties [28–41]. A recent experiment [26] provided strong evidence for zero field fractional Chern insulators (FCIs) [42–45] at fillings $\nu = -2/3$ and $-3/5$ in twisted bilayer $MoTe_2$ ($tMoTe_2$). The $\nu = -2/3$ FCI was separately found by Ref. [27]. These experiments were preceded by theoretical models of Chern bands in $tMoTe_2$ [29], as well as numerical works that found FCIs at partial fillings in $MoTe_2$ [46] and in WSe_2 [47,48]. More recently, theoretical studies combining *ab initio* lattice relaxation and exact diagonalization on $tMoTe_2$ [49,50] have also obtained FCIs.

FCIs were previously reported at high magnetic fields [51] by partially filling Hofstadter bands [52] of a substrate-induced moiré potential in graphene. Shortly thereafter, with the discovery of correlated phenomena [53,54] and spontaneous Chern insulators [55–57] in twisted bilayer graphene (TBG), FCIs in zero field were theoretically anticipated in magic-angle TBG [58–60]. Experimental observations of FCIs in this setting soon appeared [61], albeit in a small magnetic field that theory [62] found was needed to improve the bandwidth and quantum geometry. These barriers are strikingly absent in $tMoTe_2$, motivating us to go beyond zero field FCIs to an exotic gapless state—the CFL.

We will focus on the gapless CFL phase, which presents challenges [63–66] relative to the well-understood spectral and entanglement signatures present in gapped FCI phases [67–72]. Combining large scale exact diagonalization (ED) with density matrix renormalization group (DMRG) numerics, we find a broad CFL phase at experimentally realistic parameters of $tMoTe_2$. Furthermore, we present an explicit trial wave function that captures the essential features of the zero field CFL and its low energy spectrum. Finally, we discuss experimental signatures that distinguish the CFL from Fermi liquids, enabling experimental exploration.

Continuum model.—We consider a model [29] for the valence bands of a twisted TMD with gate-screened [73] Coulomb interactions

$$\hat{H} = -\hat{h} + \frac{1}{2A}\sum_{\boldsymbol{q}} V_{\boldsymbol{q}} :\hat{\rho}_{\boldsymbol{q}}\hat{\rho}_{-\boldsymbol{q}}:, \quad V_q = \frac{2\pi \tanh(qd)}{\epsilon_r \epsilon_0 q}, \qquad (1)$$

where $\hat{\rho}_{\boldsymbol{q}}$ is the density operator, A is the sample area, normal ordering is relative to filling $\nu = 0$, d is gate distance, and $\epsilon_r \approx 8$–40 is the dielectric constant [49]. Because of spin-valley locking [29], the low energy holes of the K (K') valley are locked to spin-up (-down). The total kinetic term is $h = h_K + h_{K'}$ with [29]

$$h_K = \begin{bmatrix} h^b(\boldsymbol{r}) + V/2 & T(\boldsymbol{r}) \\ T^\dagger(\boldsymbol{r}) & h^t(\boldsymbol{r}) - V/2 \end{bmatrix}, \qquad (2)$$

where $h^\ell(\boldsymbol{r}) = -(\boldsymbol{p} - \hbar v_F \boldsymbol{K}^\ell)^2/2m^* + \Delta^\ell(\boldsymbol{r})$ and $h_{K'}$ is determined by time reversal. Here the layer-diagonal terms include the quadratic monolayer TMD dispersion centered at rotated monolayer K points $\boldsymbol{K}^{t/b}$, shifted by the displacement field V, and the moiré potentials $\Delta^{b/t}(\boldsymbol{r}) = 2v \sum_{j=1,3,5} \cos(\boldsymbol{b}_j \cdot \boldsymbol{r} \pm \psi)$. The off-diagonal terms are interlayer tunnelings $T(\boldsymbol{r}) = \omega(1 + e^{i\boldsymbol{b}_2\cdot\boldsymbol{r}} + e^{i\boldsymbol{b}_3\cdot\boldsymbol{r}})$, where $\boldsymbol{b}_j$ are the reciprocal vectors obtained by counterclockwise $(j-1)\pi/3$ rotations of $\boldsymbol{b}_1 = (4\pi 3^{-1/2}\theta/a_0, 0)$. We focus on tMoTe$_2$, where recent first-principles calculations [49] (see also [29,50]) found $(a_0, m^*, V, \psi, \omega) = (3.52\ \text{Å}, 0.6m_e, 20.8\ \text{meV}, -107.7°, -23.8\ \text{meV})$. We take $\theta = 3.7°$ throughout.

Flat almost-ideal Chern band.—Figure 1(a) shows the band structure for electrons h_K. The top moiré band has Chern number $C = 1$, due to the skyrmionic character of the layer spinor [29].

Recent experiments [26,27] demonstrate that the many-body ground state is ferromagnetic (valley-polarized) in at least the range $-1.2 \lesssim \nu \lesssim -0.4$. The "parent state" for this regime is the correlated insulating state at $\nu = -1$. Figure 1(b) shows its band structure within self-consistent Hartree-Fock (SCHF), which is strongly renormalized by interactions. Strikingly, the renormalized $C = 1$ band (red) becomes almost exactly flat, with bandwidth 1.6 meV at $\theta = 3.7°$. This *reduction* [74] in bandwidth from interaction effects is highly unusual [75].

The many-body physics of such flat bands is determined by the Bloch wave functions, often through their "quantum geometry." Recent theories [42,58,80–91] emphasize the role of Kähler geometry in FCI stability. We say that a band has "ideal quantum geometry" if the trace inequality $T = \int d^2\boldsymbol{k}[\mathrm{Tr}g_{\mathrm{FS}}(\boldsymbol{k}) - \Omega(\boldsymbol{k})] \geq 0$ is saturated [58,83,87,92]; here g_{FS} is the Fubini-Study metric and Ω is the Berry curvature. Ideal bands are "vortexable" in the sense that $\hat{z}P = P\hat{z}P$, where P is the projector onto the band and $\hat{z} = \hat{x} + i\hat{y}$ [89,93]. Vortexability enables the direct construction of Laughlin-like FQHE trial states that are exact many-body ground states for ideal bands with short-range interactions [89,93,94]. Figure 1(c) shows T, the deviation from ideality, and $\sigma[\Omega]$, the standard deviation of Berry curvature. Both are small in tMoTe$_2$ for $3° \leq \theta \leq 4°$. The top valence band thus has the rare combination of excellent quantum geometry and negligible bandwidth that favors lattice realizations of exotic quantum Hall states at zero magnetic field.

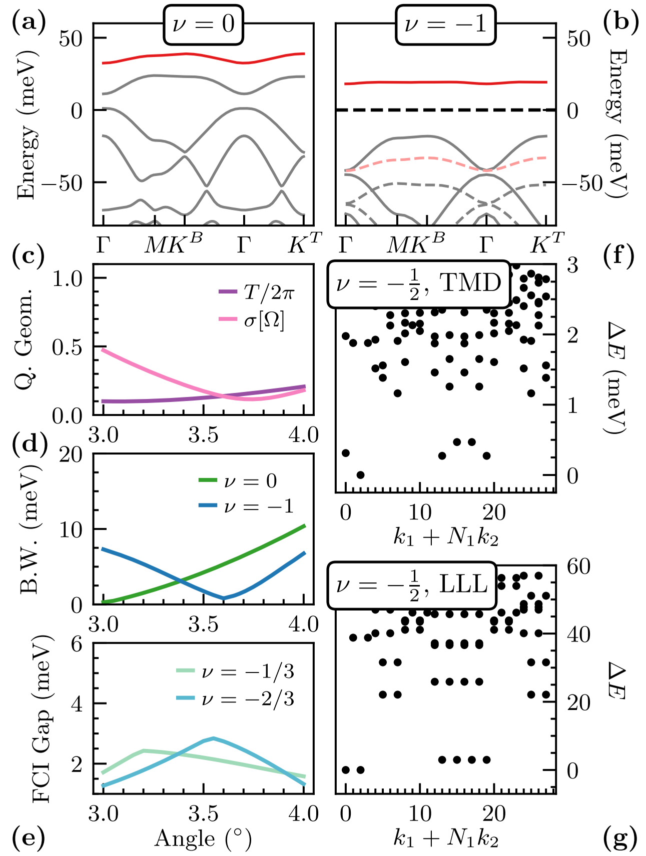

FIG. 1. The top valence band has favorable conditions for fractionalized topological phases. Band structure as seen from (a) charge neutrality and (b) from $\nu = -1$ computed from self-consistent Hartree-Fock. (c) Quantum geometry in terms of trace condition T and Berry curvature deviation $\sigma[\Omega]$. (d) Bare and SCHF bandwidths and (e) the many-body gap of FCIs at $\nu = -1/3$ and $\nu = -2/3$ as a function of twist angle. The FCI gaps are obtained from ED with $N_e = 8$ and 16, respectively. (f) and (g): ED spectrum for 14 particles at half filling for Coulomb interaction in lowest Landau level and screened Coulomb interaction in twisted MoTe$_2$, respectively. Parameters: $(\theta, \epsilon_r, d) = (3.7°, 15, 300\ \text{Å})$ unless otherwise noted.

The interacting physics of the flat band is modeled by projecting Eq. (1) via $-\hat{h} \to \sum_{\boldsymbol{k}} \epsilon(\boldsymbol{k}) \hat{c}^\dagger_{\boldsymbol{k}} \hat{c}_{\boldsymbol{k}}$ and $\hat{\rho}_{\boldsymbol{q}} \to \bar{\rho}_{\boldsymbol{q}} = \sum_{\boldsymbol{k}} \hat{c}^\dagger_{\boldsymbol{k}} \langle u_{\boldsymbol{k}} | u_{\boldsymbol{k}+\boldsymbol{q}} \rangle \hat{c}_{\boldsymbol{k}+\boldsymbol{q}}$, where $\epsilon(\boldsymbol{k})$ and $u_{\boldsymbol{k}}$ are the dispersion and periodic part of Bloch wave function. Figure 1(d) shows the bare ($\nu = 0$) and renormalized ($\nu = -1$) bandwidths versus twist angle, minimized near 3° and 3.6°, respectively. Figure 1(e) confirms that FCIs are stabilized at $\nu = -1/3$, $-2/3$—in accordance with previous results [46,49,50]. The mild angular dependence should make FCIs relatively robust to twist angle disorder. Notably the gap at $\nu = -2/3$ is largest where the bandwidth at $\nu = -1$ is smallest [95]. We therefore expect ~3.6° to be optimal for FQH physics at half filling.

Composite Fermi liquid at $\nu = -1/2$.—We now go beyond gapped FCIs and examine the more exotic gapless

CFL state [1,11]. We focus on $\nu = -1/2$ but our conclusions also apply to $\nu = -3/4$ (data in Supplemental Material [96]).

(i) Many body spectrum: Figures 1(f) and 1(g) compare the spectra of twisted $MoTe_2$ and the lowest Landau level (LLL) at half filling with 14 electrons, showing a one-to-one correspondence at low energy. The LLL spectrum uses the same geometry as $tMoTe_2$ with Coulomb interactions. This one-to-one similarity holds at all system sizes $N_e = 8$–14. We thus conclude that the ground state of $\hat{H}$ at $\nu = -1/2$ is the same phase as the half-filled LLL with Coulomb interactions—the CFL. The ground state and low-energy excitations are at precisely the momenta expected for compact composite Fermi sea (CFS) configurations [113]. See Supplemental Material [96] for other system sizes, and detailed matching of degeneracies, momenta, and excitations to CFL expectations.

(ii) Absence of electron Fermi surface: A finite quasiparticle weight $Z > 0$ gives the jump in electron occupations $n(\boldsymbol{k})$ at the Fermi surface in a regular Fermi liquid (FL). As a non-Fermi liquid, composite fermions have vanishing Z, leading to the absence of Fermi surface occupation discontinuities.

To characterize the CFL, we employ large-scale iDMRG [114,115] calculations with the TenPy library [116]. We use an infinite cylinder geometry with circumference $L_y = 5$–10, corresponding to L_y evenly spaced horizontal wires through the Brillouin zone [Fig. 2(c) inset]. We take a computational basis of hybrid Wannier orbitals [117–119], and use "MPO compression" [120,121] to accurately capture gate-screened Coulomb interactions in the flat band. Under weak interactions ($\epsilon_r = 100$), we find the FL expected from band theory at $\nu = -1/2$, with an almost-circular Fermi surface centered at Γ [Fig. 2(a), left] with radius $k_F = (A_{\mathrm{BZ}}/2\pi)^{1/2}$. The Supplemental Material [96] shows electrons, holes, and particle-hole pairs are likely gapless [122], confirming the Fermi liquid.

Under realistic interactions ($\epsilon_r = 15$) with the same parameters, the ground state has quasi-uniform occupations $|n(\boldsymbol{k}) - \frac{1}{2}| < 0.17$ [Fig. 2(a), right]. Because charge $Q_E = 1$ correlations are short ranged, the state is inconsistent with an electronic Fermi liquid. However, the state has high entanglement and significant electrically neutral correlations, consistent with the gapless density fluctuations expected from an emergent CFS. To reveal the "hidden" CFS, we turn to the structure factor.

(iii) Scattering across the composite Fermi sea: Figure 2(b) contrasts the connected structure factor $S(\boldsymbol{q}) = \langle \hat{\rho}_{\boldsymbol{q}} \hat{\rho}_{-\boldsymbol{q}} \rangle - \langle \hat{\rho}_{\boldsymbol{q}} \rangle \langle \hat{\rho}_{-\boldsymbol{q}} \rangle$ between the FL and the CFL. Both nearly vanish when $|\boldsymbol{q}| > 2k_F$, strongly implying that there is a Fermi surface in the CFL phase whose constituent fermions are not electrons. We then match the features of $S(\boldsymbol{q})$ to scattering events with different momentum transfers across the putative CFS in Fig. 2(c), e.g., $\hat{c}^\dagger_{\boldsymbol{k}=G} \hat{c}_{\boldsymbol{k}=B}$ scattering with $q_x \approx 1.94 k_F$. The *tour-de-force* work of

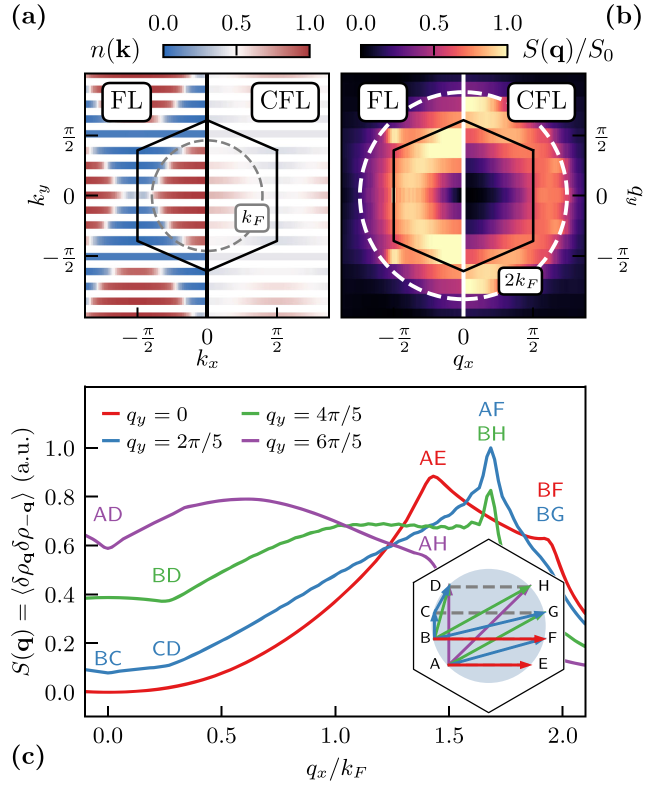

FIG. 2. Numerical identification of the composite Fermi liquid (CFL) from iDMRG. (a) Occupations $n(\boldsymbol{k})$ in the Brillouin zone at $L_y = 8$ for the Fermi liquid (FL, left side) versus the CFL (right side). (b) Connected structure factor $S(\boldsymbol{q}) = \langle \hat{\rho}_{\boldsymbol{q}} \hat{\rho}_{-\boldsymbol{q}} \rangle - \langle \hat{\rho}_{\boldsymbol{q}} \rangle \langle \hat{\rho}_{-\boldsymbol{q}} \rangle$ at $L_y = 8$. Characteristic features of a Fermi surface are visible for both the FL and CFL: near-vanishing weight outside $|\boldsymbol{q}| \approx 2k_F$, and peaks corresponding to momentum transfers inside that radius. (c) Cuts of $S(\boldsymbol{q})$ at constant q_y for $L_y = 5$ for the CFL. Each peak or inflection in $S(\boldsymbol{q})$ quantitatively matches scattering events across the almost-circular composite Fermi surface (Inset). Parameters match Fig. 1 with $\epsilon_r = 15$ (100) for the CFL (FL).

Geraedts *et al.* [63] showed such features are emblematic of the CFS arising from the half-filled LLL. As every feature in $S(\boldsymbol{q})$ corresponds to such a scattering (quantitative matching in the Supplemental Material [96]), we conclude the state has an almost-circular [123] CFS composed of non-Landau quasiparticles. These two independent numerical methods establish a CFL state at $\nu = -1/2$ (see the Supplemental Material [96] for $\nu = -3/4$).

Zero field CFL wave function.—Standard theories of composite fermions apply at $B > 0$, where emergent gauge flux cancels external magnetic flux. These cannot apply directly here at zero magnetic field. We therefore construct an explicit zero-field CFL wave function. To start, we approximate the geometry of the top $tMoTe_2$ band as ideal. Such bands have the general "LLL-like" wave function [58,87],

$$\psi_l(\boldsymbol{r}) = \phi(\boldsymbol{r}) \zeta_l(\boldsymbol{r}) = f(z) e^{-K(\boldsymbol{r})} \zeta_l(\boldsymbol{r}), \tag{3}$$

where $f(z)$ is holomorphic and $\zeta_l(\boldsymbol{r})$ is an orbital-space spinor where $\sum_l |\zeta_l(\boldsymbol{r})|^2 = 1$. Here $\phi(\boldsymbol{r})$ is the wave function of a Dirac particle in an inhomogeneous, periodic, magnetic field $B(\boldsymbol{r}) = \nabla^2 \mathrm{Re} K(\boldsymbol{r})$ with one flux per unit cell [58,124,125]. While ψ is symmetric under ordinary translations, $\phi(\boldsymbol{r})$ and $\zeta_l(\boldsymbol{r})$ are symmetric under *magnetic* translations, with opposite magnetic twists [126], giving a gauge redundancy $\phi(\boldsymbol{r}) \to e^{+i\lambda(\boldsymbol{r})}\phi(\boldsymbol{r})$, $\zeta_l(\boldsymbol{r}) \to e^{-i\lambda(\boldsymbol{r})}\zeta_l(\boldsymbol{r})$. The form Eq. (3) implies that *all* many-body wave functions within the band of interest have the form $\Psi = \Psi_\phi \prod_i \zeta_{l_i}(\boldsymbol{r}_i)$, where Ψ_ϕ is a wave function of flux-feeling particles; in the Supplemental Material [96] we interpret this fractionalization in terms of a new type of Chern band parton theory [127,128]; see also [129–134]. For example, we may use Read and Rezayi's LLL ansatz for the CFL [135] to obtain

$$\Psi(\{\boldsymbol{r}_i\}) = \mathcal{P}\det_{ij}\psi^{\mathrm{CF}}_{\boldsymbol{k}_i}(\boldsymbol{r}_j)\prod_{i<j}(z_i - z_j)^2 \prod_i e^{-K(\boldsymbol{r}_i)}\zeta_{l_i}(\boldsymbol{r}_i). \quad (4)$$

Here $\mathcal{P} = \prod_i P_i$ is the many-body projector to the top band, and $\psi^{\mathrm{CF}}_{\boldsymbol{k}_i}$ fill a Fermi sea [136].

Experimental signatures.—We conclude with experimental implications of the quantum geometry and CFL phase. Figures 3(a) and 3(b) show phase diagrams of tMoTe_2. At $\nu = -1$, SCHF finds the $|C| = 1$ phase transitions to a valley and layer polarized phase at large V. At $\nu = -1/2$, we find a broad CFL phase centered around 3.8° that competes with layer polarized phases and $C = 1$ Fermi liquids at larger V. The layer polarized region is estimated from SCHF at $\nu = -1$, where an interaction-driven layer-polarized state is more favorable. The phase diagram at $\nu = -3/4$ is similar (see Supplemental Material [96]), except the CFL is more sensitive to the displacement field.

The almost ideal quantum geometry manifests optically. If a band with projector P is vortexable, then $\hat{z}P = P\hat{z}P$ implies the velocity operator $\hat{v}^{\pm} = -i[\hat{x} \pm i\hat{y}, \hat{H}]$ must obey $(I - P)\hat{v}^{+}P = 0$, i.e., left-circularly polarized transitions are "extinguished." This gives perfect circular dichroism:

$$\frac{\sigma^+ - \sigma^-}{\sigma^+ + \sigma^-} = 1; \qquad \sigma^{\pm}(\omega) = \frac{ie^2}{\hbar}\sum_{\boldsymbol{k},a\neq b=0}\frac{f_{ab}}{\epsilon_{ab}}\frac{|\langle\psi_{\boldsymbol{k}a}|\hat{v}^{\pm}|\psi_{\boldsymbol{k}b}\rangle|^2}{\omega - \epsilon_{ab}}. \quad (5)$$

Here $\epsilon_{ab} = \epsilon_a - \epsilon_b$ are energy differences and $f_{ab} = f(\epsilon_a) - f(\epsilon_b)$ are Fermi factors. Figure 3(c) shows $\sigma^{\pm}$ for tMoTe_2 at $\nu = -1$. As the $C = +1$ band is nearly vortexable, transitions from the second and third valence bands to the empty top valence band nearly vanish, giving nearly perfect circular dichroism > 0.9 at resonance. The inset shows a control experiment: the Haldane model has Chern bands $C = \pm 1$ but not ideal geometry; σ^- is not extinguished there.

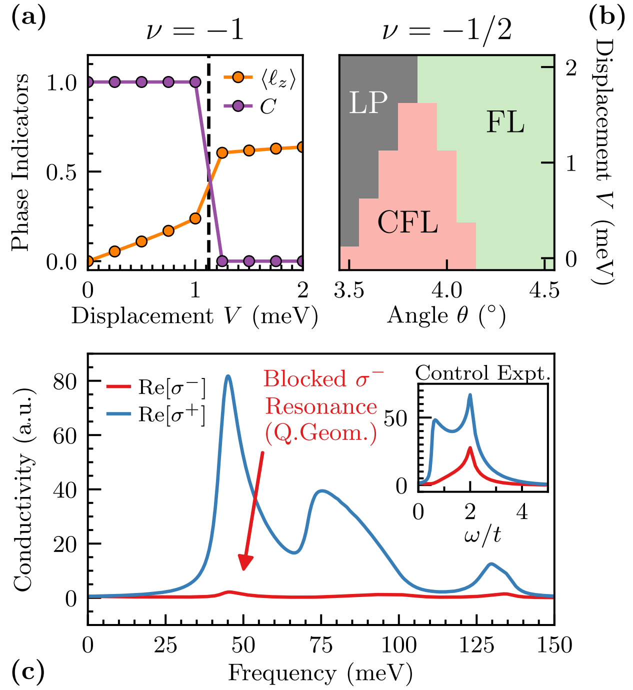

FIG. 3. Many-body phase diagrams and optical responses. (a) Phase diagram at $\nu = -1$ with $\theta = 3.7°$ showing a transition from $C = 1$ layer-unpolarized state to a $C = 0$ layer polarized state. (b) Phase diagram of the topological regime at $\nu = -1/2$: The CFL phase is shown in red, whereas the green region corresponds to the FL phase. Here LP indicates a layer polarization instability determined from $\nu = -1$ SCHF. (c) Direct optical probe of almost-ideal quantum geometry via an "extinguished" valence-valence optical responses in σ^-, Inset: the Haldane model at $(t, t_2) = (1, 0.05)$ has nonideal geometry. Parameters match Fig. 1.

Finally, we discuss direct experimental probes of the zero-field CFL. While the CFL and the FL are both compressible and metallic, they differ in that the CFL's excitations have vanishing overlap with the electron in the limit of low energies, and CFs themselves are best thought of as (doubled) *vortices* in the electronic fluid [3,4,137–139]. This observation leads to a number of striking physical responses that differ strongly from Fermi liquids. These include (i) a "pseudogap" in the tunneling density of states $A(\omega) \propto e^{-\omega_0/\omega}$ [140] as a function of bias ω, which has been observed between two CFLs with a tunnel barrier [5]; (ii) distinct dc conductivity in the clean limit: $\sigma_{xx} \to 0$ in a CFL in the absence of disorder $k_F l \to \infty$, whereas in the FL, even in a Chern band, σ_{xx} diverges [141]; (iii) strong violation of the Wiedemann-Franz law [138,139] which compares heat and charge transport; (iv) quantum oscillations with doping, that CFs feel a magnetic field $\propto (\nu - 1/2)$ and can fill Landau levels, leading to Jain-like [1] FCIs when fully developed, which can further be probed using geometric resonance with a one-dimensional periodic grating [3,6,7]; (v) vanishing thermoelectric conductance $\alpha_{xx} = j_x/(-\partial_x T)$ due to approximate emergent particle-hole symmetry [142,143]; (vi) surface acoustic wave attenuation, a contactless probe that measures $\sigma_{xx}(\boldsymbol{q}) \propto |\boldsymbol{q}|$ in the CFL [3], as opposed to $\sigma_{xx} \propto |\boldsymbol{q}|^{-1}$ in a clean FL [8].

Finally, we highlight properties of zero field CFLs that transcend LLL physics. First, the Chern bands of $MoTe_2$ have one effective magnetic flux quantum per moiré unit cell, translating to 160 T at 3.7°. This vastly exceeds laboratory magnetic fields, leading to enhanced energy scales. The lack of *real* quantizing magnetic fields, however, opens up the possibility of employing zero field experimental probes such as high resolution angle-resolved photoemission spectroscopy (ARPES). Furthermore, the exponentially suppressed tunneling density of states of the CFL could be probed through tunneling from a proximate Fermi liquid state, or via spatial variation of the twist angle, which can be used to create a CFL-FL interface within the same sample. Our work does not rule out the possibility of a continuous quantum phase transition, driven by displacement field, between the CFL and FL [133], which could be studied experimentally. Since the effective magnetic field of the TMD originates from spontaneous breaking of time reversal symmetry through valley polarization, rather than external magnetic field, domains between opposite valley polarizations and hence between time-reversal-related CFLs are expected. Transport properties across such a domain wall would interrogate composite fermions in an entirely new regime, and potentially shed light on their proposed Dirac character [4,137,139]. Finally, we note that moiré phonons occur on the same scale as the effective magnetic length in this system; their interplay with CFL physics is unclear at present and worthy of future study.

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Note added.—Recently, we were alerted by Taige Wang to the importance of layer polarization after our Letter was announced on arXiv. Also, [144,145] appeared recently, which overlap with parts of this work, and Ref. [146] overlaps with the optical responses discussed here. Subsequent to our work, transport experiments [147] on twisted $MoTe_2$ were performed and the results are consistent with our findings.

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