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High-Temperature Fractional Quantum Hall State in Floquet-Kagome Flat Band

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Abstract: Fractional quantum Hall effect (FQHE) has been predicted in topological flat band (FB) by single-particle band structure combined with phenomenological theory or solution of many-body lattice Hamiltonian with fuzzy parameters. A long-standing roadblock towards realization of FB-FQHE is lacking the many-body solution of specific materials under realistic conditions. Here, we demonstrate a combined study of single-particle Floquet band theory with exact diagonalization (ED) of many-body Hamiltonian. We show that a time-periodic circularly polarized laser inverts the sign of second-nearest-neighbor hopping in a Kagome lattice and enhances spin-orbit coupling in one spin channel, to produce a Floquet FB with a high flatness ratio of bandwidth over band gap, as exemplified in monolayer $Pt_3C_{36}S_{12}H_{12}$. The ED of the resultant Floquet-Kagome lattice Hamiltonian gives a one-third-filling ground state with a laser-dependent excitation gap of FQH state, up to an estimated temperature above 70 K. Our findings pave the way to explore the alluding high-temperature FB-FQHE.

Integer and fractional quantum Hall effect (IQHE and FQHE), arising from intrinsic band topology in crystal lattices, can potentially overcome the low temperature and high magnetic field required for their counterparts in 2D electron gas [1-8]. The IQHE without Landau level (LL) was originally proposed by Haldane in a honeycomb lattice model with a Chern insulating gap [9], and recently predicted [10-12] and observed [13] in Chern insulators. Also, a partially filled Chern flat band (FB) in 2D lattices is predicted to support high-temperature FQHE [14-17]. A widely used numerical technique to identify the existence of FQHE in lattice systems is the exact diagonalization (ED) of a FB Hamiltonian with many-body Hubbard interaction [17-22], where the signatures of Laughlin-like state [20-22], including ground-sate degeneracy, spectral flow, and fractional statistics, can be shown for partially filled Chern FB.

Some 2D materials have been found to host FB from first-principles calculations [23-32]; however, none of the them has a sufficiently large flatness ratio of band gap (Δ) over band width (w), i.e., $\Delta/w >>1$, which is required for realizing high-temperature FQHE [14-17]. It has been shown that to obtain an "isolated" FB with large Δ/w, peculiar negative kinetic hopping integral and strong spin-orbit coupling (SOC) are needed in a Kagome lattice [14]. But such requirements are unlikely to be met with the natural decay of lattice hopping with interatomic spacing in real materials, and indeed, all the FBs found so far have too small a flatness ratio of $\Delta/w \sim 1$ [23-32]. On the other hand, the existing many-body ED calculations are carried out for FB-lattice models using artificial hopping parameters, SOC strength and Hubbard interaction [21, 22], which are disconnected from realistic materials and conditions.

Therefore, there remain at least two challenges towards demonstrating the FB-FQHE in the real world: firstly, developing an approach to tune the lattice hopping parameters, so that they can meet the stringent requirements to attain a large Δ/w; secondly, solving the many-body lattice Hamiltonian corresponding to real materials, to directly show the solution of FQHE. In this Letter, we demonstrate that the Floquet-FB in a Kagome lattice can reach an unprecedented high $\Delta/w >> 1$ through photo-engineered lattice interactions. A time-periodic circularly polarized laser (CPL) strongly alters the steady-state electronic interactions in a Kagome lattice by virtual photon processes (VPPs), in particular to invert the second-nearest-neighbor (2NN) kinetic hopping and enhance the effective SOC in one spin channel locked with CPL helicity.

Consequently, the originally dispersive Chern band is transformed into a FB. In a prototypical example of monolayer $Pt_3C_{36}S_{12}H_{12}$ (HTT-Pt), the so-formed FB reaches a flatness ratio of Δ/w ~ 40. Furthermore, we use ED method to solve the material-specific lattice Hamiltonian of HTT-Pt under laser illumination, based on realistic Floquet hopping, SOC and Hubbard interaction parameters. We show that the original electronic Kagome lattice Hamiltonian *does not* support the FQHE; only the Floquet-Kagome lattice can host the FQH states at the one-third filling of a FB, with a many-body excitation gap tunable by laser intensity, reaching an estimated temperature above 70 K.

Recently, irradiating time-periodic light on crystals has been shown to offer an effective way to create and control the experimentally measurable Floquet topological electronic states [33-36], such as the nonequilibrium IQHE [37, 38] and FQHE [39] closely related to the present study, Floquet-Dirac/Weyl semimetals [40-44], and topological FB positions [45, 46]. Here, we illustrate an interesting case of laser created FB with large Δ/w for high-temperature FQHE.

To study coherent interactions between laser field and Kagome lattice, we consider NN (2NN) kinetic hopping integral $\gamma_1{}^0$ ($\gamma_2{}^0$) and SOC strength $\lambda_1{}^0$ ($\lambda_2{}^0$) in two spin channels $s = \pm 1$ (see details in Fig. S1 and Section I of Supplemental Material (SM) [47]), and we adopt a CPL with a time-dependent vector potential $\boldsymbol{A}(t) = A_0(\eta\cos\omega t, \sin\omega t)$, where $\hbar\omega$, A_0, and η are photon energy, laser amplitude, and laser helicity, respectively. Under a time-periodic off-resonant CPL, the Floquet-Bloch Hamiltonian can be expanded as [38, 52]

$$H_F(\boldsymbol{k}) = H_0(\boldsymbol{k}) + \sum_{n\in Z^+} \frac{1}{n\hbar\omega}\left[H_{-n}(\boldsymbol{k}), H_n(\boldsymbol{k})\right] + O\left(\frac{1}{\omega^2}\right) \tag{1}$$

The VPPs, comprising photon absorption ($n' \in Z^+$) and emission ($n' \in Z^-$) terms, $H_{n'}(\boldsymbol{k}) = \frac{1}{T}\int_0^T H(\boldsymbol{k},t)e^{in'\omega t}dt$, will modify ($\gamma_1$, λ_1) and (γ_2, λ_2) [Fig. 1(a)], derived as

$$\gamma_1 = \gamma_1{}^0 J_0(A_1) + s\eta\left\{2J_1(A_1)\left[\sqrt{3}J_1(A_1)B_1 + J_1(A_2)C_1\right] + \sqrt{3}J_2(A_1)\cdot\left[J_2(A_1)B_1 + J_2(A_2)C_1\right]\right\}/(\hbar\omega), \tag{2}$$

$$\lambda_1 = \lambda_1{}^0 J_0(A_1) + s\eta\left\{J_1(A_1)\left[\sqrt{3}J_1(A_1)D_1 - 2J_1(A_2)C_2\right] + \sqrt{3}J_2(A_1)\left[J_2(A_1)D_1/2 - J_2(A_2)C_2\right]\right\}/(\hbar\omega), \tag{3}$$

$$\gamma_2 = \gamma_2{}^0 J_0(A_2) - \sqrt{3}s\eta\left\{2\left[J_1^2(A_1)B_1 + J_1^2(A_2)B_2\right] - \left[J_2^2(A_1)B_1 - J_2^2(A_2)B_2\right]\right\}/(\hbar\omega), \tag{4}$$

$$\lambda_2 = \lambda_2{}^0 J_0(A_2) + \sqrt{3}s\eta\left\{\left[J_1^2(A_1)D_1 - J_1^2(A_2)D_2\right] - \left[J_2^2(A_1)D_1 + J_2^2(A_2)D_2\right]/2\right\}/(\hbar\omega), \tag{5}$$

where $A_1 = aeA_0/2\hbar$ (a is lattice constant), $A_2 = \sqrt{3}A_1$, $B_1 = \gamma_1{}^0\lambda_1{}^0$, $B_2 = \gamma_2{}^0\lambda_2{}^0$, $C_1 = \gamma_1{}^0\lambda_2{}^0 - \gamma_2{}^0\lambda_1{}^0$,

$C_2 = \gamma_1^0\gamma_2^0 + \lambda_1^0\lambda_2^0$, and $D_1 = (\gamma_1^0)^2 - (\lambda_1^0)^2$, $D_2 = (\gamma_2^0)^2 - (\lambda_2^0)^2$. J_m is the first kind of Bessel function with an order of $m = 0, 1, 2$, which comes from the m-order VPP in Eq. (1). For high-frequency driving, the first-order term ($\sim\omega^{-1}$) dominates the Floquet band structure, so that the higher-order terms can be neglected (Fig. S2 [47]) [35, 40, 52].

By chiral symmetry, the effect of a left-handed CPL ($\eta = 1$) on spin-up ($s = 1$) bands is the same as that of a right-handed CPL ($\eta = -1$) on spin-down ($s = -1$) bands. Thus, below we will focus on the case of $\eta = 1$ and $s = 1$. To illustrate photo-engineered interactions and resulting Floquet-Kagome band structure, we choose a prototypical example having tight-binding hopping integrals of $\gamma_1^0 = 48$ meV, $\lambda_1^0 = 0.180\gamma_1^0$, $\gamma_2^0 = 0.112\gamma_1^0$, and $\lambda_2^0 = 0.038\gamma_1^0$, which are fitted from the first-principles band structure of monolayer HTT-Pt (Fig. S3 and Section II in SM [47]) [30, 31, 53]. Applying a CPL with $\hbar\omega = 8.0\gamma_1^0$, γ_2 decreases with the increasing A_0, as shown in Fig. 1(b). Beyond $A_0 = 476$ V/c (corresponding to 0.039 V/Å or 2.0×10^{10} W/cm^2), γ_2 is inverted from positive to negative, which is essential to flatten the bottom Chern band (see details in Fig. S4 [47]). At the same time, both λ_1 and λ_2 increase with A_0 [Fig. 1(b)], which is critical to increase Δ/w. In sharp contrast to equilibrium Kagome materials, the Floquet-Kagome hopping integrals no longer decay exponentially with the interatomic distance and can have different phases (signs).

The unusual photo-engineered interactions in Fig. 1(b), i.e. inverting the 2NN kinetic hopping and enhancing the SOC in one spin channel locked with laser helicity, lead to an intriguing evolution of Floquet-Kagome bands. The equilibrium spin-up bands have a direct SOC gap $\Delta = 0.887\gamma_1^0$, separating the bottom Chern band ($C = -1$) from the middle Dirac band ($C = 0$) [left panel in Fig. 1(c)]. The bottom Chern band has a sizable width $w = 3.228\gamma_1^0$ and a small flatness ratio $\Delta/w = 0.27$. The CPL significantly reduces its width, such as to $w = 0.022\gamma_1^0$ with $A_0 = 800$ V/c, ~150 times narrower than its equilibrium value, as shown in the right panel of Fig. 1(c). Also, it becomes well separated from the band above by a large SOC gap $\Delta = 0.650\gamma_1^0$. Consequently, an unprecedented high value of $\Delta/w \sim 29$ is achieved. The non-monotonic dependence of Δ/w on A_0, peaking at $A_0 = 800$ V/c for $\hbar\omega = 8\gamma_1^0$, can be better understood by analyzing simplified models with only NN equilibrium hopping integrals, and evaluating the detailed dependence of Δ, w and Δ/w on laser-driven hopping integrals (Section

IV and Figs. S4, S5 in SM [47]). The spin-down Floquet bands evolve differently, e.g., exhibiting bands with a high Chern number (Fig. S6 [47]). With a spontaneous Zeeman splitting of $M = -2\gamma_1{}^0$ induced by partially filled FB [23, 30], the spin-up and -down Floquet bands of monolayer HTT-Pt can be fully separated from each other (Fig. S7 [47]).

We have analyzed the laser-driven distribution of Berry curvature (Ω) of the photo-flattened Chern band to identify essential features for realizing FQHE. At equilibrium, Ω of bottom dispersive Chern band is highly localized at K and K' points of the first Brillouin zone (FBZ) [top panel in Fig. 1(d)], which is characterized by a large mean-square deviation of $\langle(\Delta\Omega)^2\rangle = 7.2 \times 10^{-2}$ [54]. Differently, Ω of the ultra-flat and isolated bottom Floquet-Chern band becomes delocalized in the whole FBZ [bottom panel in Fig. 1(d)], characterized with $\langle(\Delta\Omega)^2\rangle = 3.9 \times 10^{-3}$. The uniform distribution of Berry curvature in a FB indicates a very short magnetic length l on the order of lattice constant [14], in analogy to a LL, which favors the FQHE with a large excitation gap and high temperature [8, 55] (see section V of SM [47]).

To identify the optimal laser intensity and photon energy for maximizing band gap and flatness ratio, we have mapped out Δ and Δ/w as functions of $\hbar\omega$ and A_0. Figures 2(a) and 2(b) show their dependence on A_0 for a chosen $\hbar\omega = 8\gamma_1{}^0$, respectively. A maximum $\Delta/w \sim 29$ (also a large $\Delta = 0.650\gamma_1{}^0$) is reached at $A_0 = 800$ V/c. Figures 2(c) and 2(d) show the phase diagrams of Δ and Δ/w in the parameter space of $\hbar\omega$ and A_0. There exists a large experimentally accessible parameter space to achieve a Chern FB with $\Delta > 0.54\gamma_1{}^0$ and $\Delta/w > 10$. In general, as the photon energy increases, the required laser intensity is stronger in order to create the desired band structure, because of an inverse linear scaling between the intensity of VPP and photon energy [Eq. (1)]. With $\hbar\omega = 9.1\gamma_1{}^0$ and $A_0 = 812$V/c, a maximum $\Delta/w \sim 43$ is achieved with a large $\Delta = 0.60\gamma_1{}^0$ (Fig. S8 [47]).

Next, we investigate the possible existence of FQHE and assess its temperature in a specific material under realistic conditions, by performing a series of ED studies of a Kagome lattice with Floquet hopping parameters, corresponding to HTT-Pt under laser illumination. The many-body Hamiltonian is $\hat{H} = \hat{H}_0 + U\sum_{<i,j>}\hat{n}_i\hat{n}_j$, where $\hat{H}_0$ is the HTT-Pt-specific Floquet Kagome Hamiltonian with $\gamma_1{}^0 \sim 0.048$ eV, $\hat{n}_j$ is the on-site particle number operator, and $U \sim$

0.14 eV (3 γ_1^0) is the NN Hubbard repulsion (see details in Section VI of SM [47]). We note that in HTT-Pt, U is larger than the single-particle band gap (Δ). Because the two lower bands have respectively Chern number −1 and 0, any mixing between them will not change the Chern number of the lowest FB with band width w [19]. Hence if $U >> w$, interactions dominate and partial filling of the FB leads to a strongly correlated state, such as a FQH state. So, we exactly diagonalize the many-body Hamiltonian projected to the lowest FB for a finite system with $N_x \times N_y$ unit cells (Total sites = $3 \times N_x \times N_y$). Filling factor equals to $\upsilon = \frac{N_e}{N_x \times N_y}$, where N_e is the number of electrons in the system. Under periodic boundary conditions, we implement translational symmetries and diagonalize the Hamiltonian in each momentum sector $q = (2\pi k_x/L_x\,, 2\pi k_y/L_y)$ with k_x and k_y being the integers.

For comparison, we first calculated the energy spectra for the one-third filling of a 4×6 system for the equilibrium lattice, as shown in Fig. 3(a). There is no clear gap, or the celebrated 3-fold degeneracy of ground state [17-22], or any other identifiable signature of 1/3 FQH state. This indicates that the HTT-Pt by itself does not support FQHE. We then calculated the energy spectra for the Floquet lattices with different laser energies and intensities. One typical example is shown in Fig. 3(b) with $\hbar\omega$ = 8.0γ_1^0; A_0 = 800 V/c. One notices a clear gap separating the three-fold degenerate ground-state manifold from the excited states. Due to the finite size effects, the degeneracy of ground state is slightly lifted, but the energy spread is much smaller than the gap. Also, if one state in the ground-state manifold lies in the momentum sector (k_1, k_2), the next state can be always found at $(k_1 + N_e, k_2 + N_e)$ [modulo (N_x, N_y)]. This correlation implies that the ground state has a non-trivial topology [18, 20]. The key features of finite gap and ground-state degeneracy have been checked for convergence with respect to system size (see Fig. S9 [47]).

To further ascertain the ground state is nontrivial, we also calculated the spectral flow [Fig. 3(c)], under a twisted boundary condition which is equivalent to insertion of magnetic flux. According to Laughlin's gauge argument, if one adiabatically inserts three quantum fluxes into the one-third-filling FQH state, the states should evolve back to their original configuration [2]. This can be clearly seen in Fig. 3(c) for the three states in the ground-state manifold. There is no mixing between them and the excited states throughout the flux insertion. Thus, these states

share a total Chern number of 1, giving a quantized Hall conductance of $\sigma_H = \frac{1}{3}\frac{e^2}{h}$ [56].

The many-body gap is $\Delta_{ED} \sim 6$ meV [Fig. 3(b)], as calculated for the Floquet-Kagome lattice using $\hbar\omega = 8.0\gamma_1{}^0$, $A_0 = 800$ V/c, and $U = 3\gamma_1{}^0$. Taking this Δ_{ED} as a rough estimate of critical temperature (T_c) for FB-FQHE, we obtain $T_c \sim 70\ K$ [14, 47]. In Fig. 3(d), we show that T_c increases monotonically with the increasing laser amplitude. We have calculated Δ_{ED} as a function U at fixed laser intensity $A_0 = 800$ V/c (i.e. fixed lattice hopping), and found $\Delta_{ED} = 0.042\ U$ (Fig. S10 [47]), which is qualitatively same but quantitatively about half of the empirical formula $\Delta_{\mathrm{FQHE}} = 0.09\ U$ [14] for LLs. In addition, we calculated Δ_{ED} as a function of kinetic lattice hopping (γ_1) for constant U, and found $\Delta_{ED} = 0.03\ (\gamma_1{}^0)^2/\gamma_1$ (Fig. S10 [47]). This leads to a combined expression, $\Delta_{ED} = 0.01\ (U/\gamma_1)\ \gamma_1{}^0$, showing the dependence of Δ_{ED} on both U and γ_1. It points to the apparent difference between FB- and LL-FQHE, as the former arises from a lattice system while the latter is lattice-free.

In addition, we have verified the fractional exclusion statistics in the Floquet FB, which is another key characteristics of FQH state in LLs [57]. In Fig. 4(a), we show the calculated quasihole excitation spectrum for the case of Fig. 3(b). This is done by keeping N_e fixed and varying N_x and/or N_y, to introduce a hole. The counting of number of states below the gap in quasihole spectrum must correspond to the counting given by the (1,3)-admissible rule based on the generalized Pauli principle [21],

$$N_{FQH}^{N_e} = N_x N_y \frac{(N_x N_y - 2N_e - 1)!}{N_e!\,(N_x N_y - 3N_e)!} \tag{6}$$

As an example, we performed this calculation for eight electrons in a 5×5 lattice. The number of states below the gap is 25, in exact agreement with the (1,3)-admissible rule.

Moreover, to rule out other possible ground states, such as Wigner crystal state, we calculated the particle-cut entanglement spectrum (PES). By partitioning N_e into two groups of N_A and N_B, and tracing out the degrees of freedom for N_B, we calculate the reduced density matrix $\rho_A = Tr_B \sum_i |\phi_i><\phi_i|$, where the sum is over the three states in the ground-state manifold and $|\phi_i>$ is their respective many-body wavefunctions. The eigenvalues of this matrix are given by $e^{-\xi}$ [21, 58], and the entanglement energy levels ξ can then be displayed

in groups labeled by the momentum (k_1, k_2) of N_A particles. In Fig. 4(b) we plot the PES for $N_A = 3$ in a 4×6 lattice. A clear entanglement gap can be seen, and the number of states below the gap matches exactly the quasihole counting for three particles in a 4×6 lattice, indicating an FQH state [21, 22, 58]. Besides A_0 = 800 V/c, the 1/3 fractional Hall conductance and key characteristics of the predicted FQHE are also demonstrated for other laser amplitudes (see Figs. S11, 12 [47]).

To facilitate a direct comparison with experimental measurement of Floquet states [34, 36, 59-61], we simulated time- and angle-resolved photoelectron spectroscopy, by adopting the experimental pump-probe scheme (see Section VII and Fig. S13 in SM [47]). The simulation shows that there are ~ 5% electrons photo-excited from the bottom Floquet FB (Fig. S13 [47]) to form Floquet dressed states through VPPs, which would result in a conductivity plateau of "IQHE" deviating from an exact integer (similar to that in Floquet graphene [36]), when the Fermi level lies in the gap above the bottom Floquet FB. Interestingly, however, this may not affect the FQHE measurement, because FQHE corresponds to the many-body ground state that occurs at exact 1/3 FB occupation, usually achieved by gating to tune the chemical potential. The laser intensity and energy for achieving reasonably large Δ/w corresponds to an electric field of ~ 1×10^8 V/m, in the same range as experiments where the materials remain stable [34, 36, 59-61].

In summary, we have carried out a comprehensive study of FQHE in topological FB in specific materials under realistic conditions, by combining single-particle Floquet band theory with exact diagonalization of many-body Hamiltonian. We conclude that the naturally existing electronic materials are unlikely to support the FQHE because of too small a FB flatness ratio. We propose one viable approach to increase the flatness ratio is by Floquet band engineering via photo-inverted lattice hopping coupled with photo-enhanced SOC interaction. Using monolayer HTT-Pt as a prototypical example, we show that a Floquet-Kagome lattice may exhibit a 1/3-filling FB-FQHE above the liquid nitrogen temperature.

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Figures

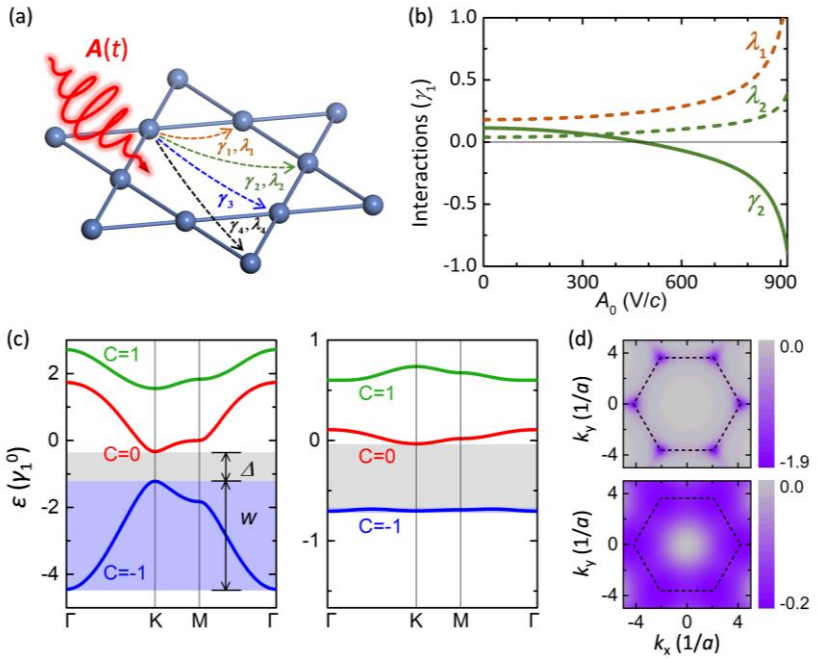

FIG. 1. Photo-engineered Kagome interactions and Floquet bands. (a) Schematic illustration of lattice hopping driven by a time-periodic CPL. γ_1 (λ_1), γ_2 (λ_2), γ_3, and γ_4 (λ_4) represent the NN, 2NN, 3NN, and 4NN kinetic (SOC) hopping, respectively. (b) Evolution of λ_1, γ_2, and λ_2 vs. laser amplitude A_0 for spin-up bands in monolayer HTT-Pt under a CPL with $\hbar\omega = 8{\gamma_1}^0$ and $\eta = 1$. (c) Left panel: Kagome Bands in equilibrium. Right panel: Floquet Kagome bands induced by a CPL with $A_0 = 800$ V/c in (b). (d) Top panel: The Berry curvature Ω for the bottom dispersive Chern band with $\Delta/w \sim 0.27$ in the left panel of (c). The dashed lines mark the FBZ. Bottom panel: The laser-driven Ω of the bottom FB with $\Delta/w \sim 29$ in the right panel of (c).

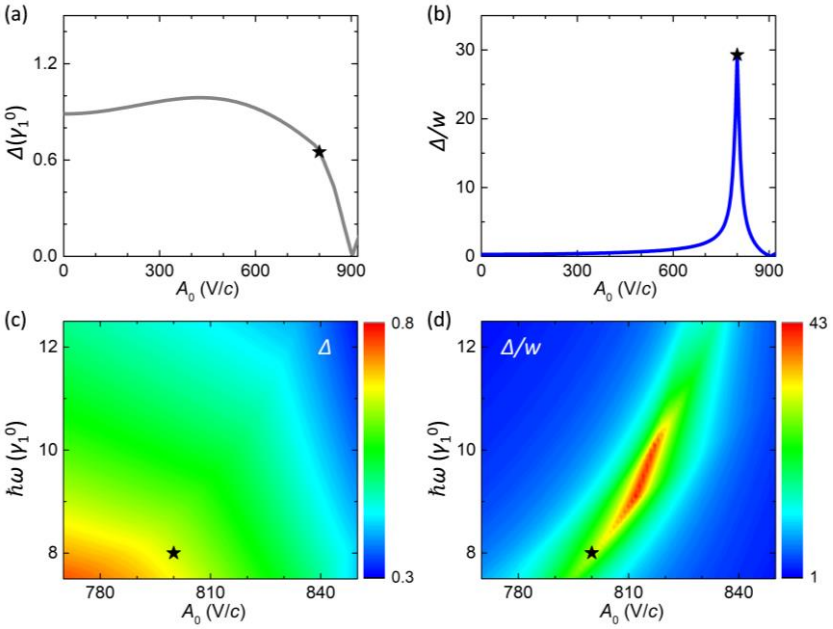

FIG. 2. Evolution of (a) spin-up band gap Δ and (b) flatness ratio Δ/w vs. laser intensity, produced for HTT-Pt under CPL with $\hbar\omega = 8\gamma_1{}^0$. (c, d) Phase diagram of Δ and Δ/w in the parameter space of $\hbar\omega$ and A_0. The star marks the specific case with $\Delta/w \sim 29$ shown in the right panel of 1(c), as used for ED study.

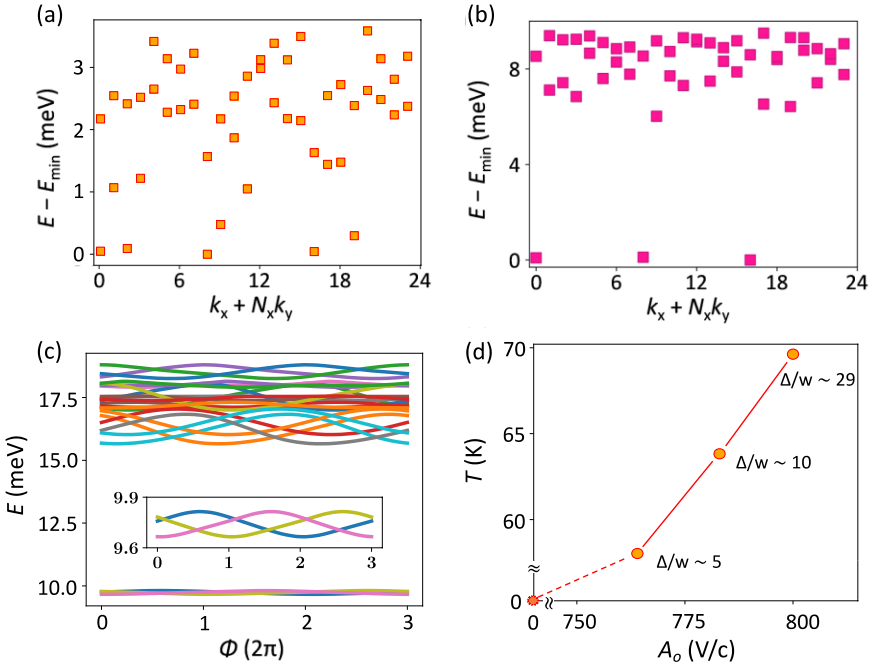

FIG. 3. ED solution of the Floquet FB-FHQE in HTT-Pt at $U = 3\ \gamma_1{}^0$. (a) Low-energy spectrum of $N_e = 8$ electrons in the equilibrium lattice with 4 × 6 Kagome unit cells. (b) Low-energy spectrum of $N_e = 8$ electrons in the same lattice as in (a) driven by CPL with $A_0 = 800$ V/c, $\hbar\omega = 8\gamma_1{}^0$. (c) Spectral flow spectrum: evolution of (b) under the flux insertion along y direction. (d) Temperature estimated by the many-body gap of 1/3 FQH state vs. A_0.

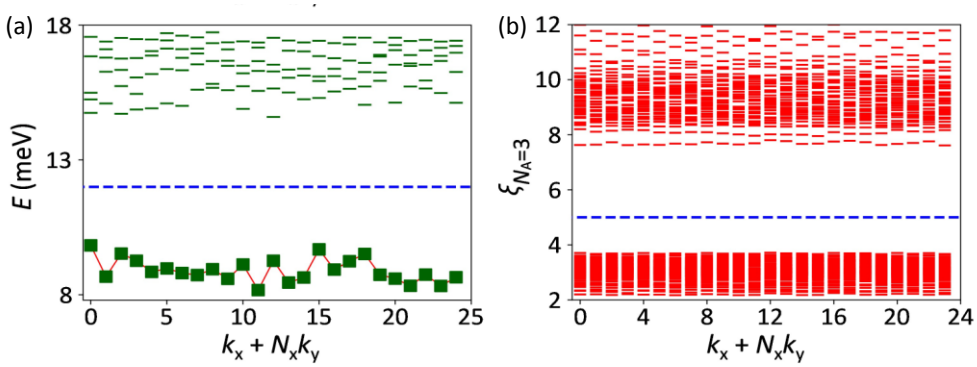

FIG. 4. Fractional statistics for the 1/3 FQH state as in Fig. 3. (a) Low-energy spectrum of 8 electrons in the cluster with 5 × 5 Kagome cells (equivalent to quasihole excitation of 8 electrons in 4 × 6 Kagome cells). (b) PES probing a group of $N_A = 3$ electrons in Fig. 3(b).