Photoinduced Floquet topological magnons in Kitaev magnets

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Abstract – We study periodically driven pure Kitaev model and ferromagnetic phase of the Kitaev-Heisenberg model on the honeycomb lattice by off-resonant linearly and circularly polarized lights at zero magnetic field. Using a combination of linear spin wave and Floquet theories, we show that the effective time-independent Hamiltonians in the off-resonant regime map onto the corresponding anisotropic static spin model, plus a tunable photoinduced magnetic field along the [111] direction, which precipitates Floquet topological magnons and chiral magnon edge modes. They are tunable by the light amplitude and polarization. Similarly, we show that the thermal Hall effect induced by the Berry curvature of the Floquet topological magnons can also be tuned by the laser field. Our results pave the way for ultrafast manipulation of topological magnons in irradiated Kitaev magnets, and could play a pivotal role in the investigation of ultrafast magnon spin current generation in Kitaev materials.

Introduction. – Topological band theory of solid-state materials has dominated many aspects of condensed-matter physics over the past decade [1,2]. The original concept of topological band theory is rooted in insulating electronic systems possessing a nontrivial gap in their energy band structure. They are characterized by the appearance of gapless chiral edge electron modes traversing the bulk gap, which are topologically protected by the Chern number or the $\mathbb{Z}_2$ index of the bulk bands [1,2].

Generally, the concept of topological band structure is independent of the statistical nature of the quasiparticle excitations and, therefore, is not restricted to insulating electronic systems. Recently, there has been a tremendous interest in the topological properties of spin excitations in insulating quantum magnets. In fact, bosonic topological spin excitations (magnons and triplons) have been studied in many different insulating quantum magnets [3–14], and the appearance of chiral edge modes and bulk Chern number has been demonstrated [4–6]. Recently, bosonic topological spin excitations mimicking electronic topological insulators have been experimentally observed in the kagome ferromagnet Cu(1,3-bdc) [9], dimerized quantum magnet SrCu$_2$(BO$_3$)$_2$ [10], and honeycomb ferromagnet CrI$_3$ [11].

The Mott-insulating honeycomb Kitaev magnets are currently of great interest [15–33]. Candidate Kitaev materials include Na$_2$IrO$_3$ and α-RuCl$_3$ [26–30]. Recently, topologically protected spin waves have been predicted in the fully polarized phase of the pure Kitaev model [12] and the Kitaev-Heisenberg model [13] at high magnetic field. In the former, the topological magnons and chiral edge states present in linear spin-wave approximation survive magnon-magnon interactions and, therefore, are robust [12]. Indeed, the manipulation of topological magnons and magnon spin currents is essential for their practical applications in ultrafast magnetic data storage, magnetic switching, and magnon spintronics [34].

The tremendous interest in topological quantum phases of matter has led to different alternative ways for inducing them in quantum materials. Recently, irradiated solid-state materials have provided an alternative route to extend the search for topological quantum materials in electronic systems [35–41]. In this formalism, topologically trivial systems can be periodically driven to nontrivial topological systems termed Floquet topological insulators [38,39]. They have an advantage over their static (equilibrium) topological counterpart, in that

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their intrinsic properties can be manipulated and different topological phases can be achieved. In irradiated insulating quantum magnets with charge-neutral spin excitations [42–45], the Floquet physics can emerge from the coupling of the electron spin magnetic dipole moment to the laser electric field through the time-dependent version of the static Aharonov-Casher phase [46,47], which acts as a vector potential or gauge field to the spin current [48]. In this case, the resulting Floquet physics can reshape the underlying Hamiltonian to stabilize magnetic phases and provides a promising avenue for inducing and tuning Floquet topological spin excitations [42–44], with a direct implication of generating and manipulating ultrafast spin current using terahertz (THz) radiation [49]. Lately, THz electric-field amplitude exceeding 100 MV/cm has been recently performed in the candidate material α-RuCl₃ [50]. In this respect, resonant time-domain THz spectroscopy has been recently performed in the candidate Kitaev material α-RuCl₃ [51].

In this paper, we propose a tunable mechanism to induce and manipulate topological magnons in irradiated Kitaev magnets at zero magnetic field. We study the pure Kitaev model [52] and the ferromagnetic phase of the Kitaev-Heisenberg model, which are already present in the zero magnetic-field classical phase diagram of the Kitaev-Heisenberg model on the honeycomb lattice [28]. Using linear spin wave and Floquet theories, we show that when the models are periodically driven by off-resonant linearly and circularly polarized lights, they effectively map onto the corresponding static spin model plus a tunable photoinduced magnetic field along the [111] direction, which is perpendicular to the honeycomb plane. The photoinduced magnetic field precipitates the existence of Floquet topological magnons and chiral edge modes, in a similar fashion to a homogeneous magnetic field in the undriven systems [12,13]. However, the Floquet topological magnons can be tuned by the amplitude and polarization of the laser field. Likewise, we demonstrate that the resulting Floquet thermal Hall conductivity can be tuned by the laser field. The photoinduced magnetic field required to induce magnetic order and Floquet topological magnons in the pure Kitaev model lies in the interval 0 < h(E₀, φ) < 2AS, where E₀, φ are the amplitude and polarization of the laser field, A > 0 is the overall energy scale of the spin exchange interactions and S is the spin value. Therefore, h(E₀, φ) is much smaller than the high magnetic field h > 4AS required to induce topological magnons in the undriven pure Kitaev model [12]. Interestingly, the Floquet topological magnons in the irradiated Kitaev magnets do not require an explicit time-reversal symmetry breaking term from the second-order virtual-photon absorption and emission processes [36], which is strictly required in order to induce Floquet topological states in other irradiated quantum systems [36,37,41–43].

Model. — We study the Kitaev-Heisenberg model on the honeycomb lattice with nearest-neighbour interaction.

The spin Hamiltonian reads [25–32]

\[ \mathcal{H} = 2J_K \sum_{\langle ij \rangle} S_i^\alpha S_j^\beta + J_H \sum_{\langle ij \rangle} \vec{S}_i \cdot \vec{S}_j, \]  

(1)

where the first term corresponds to the bond-dependent Kitaev interaction and the second term to the isotropic Heisenberg interaction. The bond directions are denoted by \( \gamma = \{x, y, z\} \) as shown in fig. 1. We parameterize the interactions as \( J_H = A \cos \vartheta \) and \( J_K = A \sin \vartheta \), where \( \vartheta \in [0, \pi/2] \) and \( A = \sqrt{J_H^2 + J_K^2} > 0 \) is the overall energy scale of the exchange interactions, with \( A \approx 8\) meV in some real materials [30]. The classical phase diagram of eq. (1) has been established in the \( \vartheta \) space [28,29].

The zig-zag phase of eq. (1) is believed to describe the honeycomb magnetic materials Na₂IrO₃ and α-RuCl₃ [28,30]. Recent studies have shown that the fully-polarized phase of the pure Kitaev model (\( \vartheta = \pi/2 \)) [12] and the Kitaev-Heisenberg model (\( \vartheta = 5\pi/4 \)) [13] at high magnetic field possess topological magnon modes. The purpose of this paper is to periodically drive the magnon topologically trivial phases of eq. (1) to Floquet topological magnon insulators for \( \vartheta = \pi/2 \) and \( \vartheta = 5\pi/4 \).

Irradiated Kitaev magnets. — In the presence of an intense laser field with a dominant time-dependent electric-field component \( \vec{E}(\tau) \), the spin magnetic dipole moment of an electron \( \vec{\mu}_S = -g \mu_B \vec{n} \) hopping along the magnetization direction \( \vec{n} \) will accumulate a time-dependent Aharonov-Casher phase [42–45]

\[ \Phi_{ijk}(\tau) = \mu_m \int_{t_i}^{t_j} \vec{E}(\tau) \cdot d\vec{E}, \]  

(2)

where \( \mu_m = g \mu_B / hc \), \( g \) is the spin-g factor, \( \mu_B \) is the Bohr magneton, \( h \) is the reduced Planck’s constant, and \( c \) is the speed of light. Here, \( \vec{E}(\tau) = \vec{E}(\tau) \times \vec{n} \) with \( \vec{E}(\tau) = -\partial_\tau \vec{A}(\tau) \), where \( \vec{A}(\tau) \) is the time-dependent vector potential of the applied laser field.

It is convenient to introduce orthonormal basis vectors \( (\hat{l}, \hat{m}, \hat{n}) \), where \( \hat{n} \) points along the cubic [111] direction,
perpendicular to the honeycomb plane [53]. We can now write eq. (1) in the new basis. In this new basis, the spin dipole moment of an electron couple to the laser electric field through the Aharonov-Casher phase, in the same way the electron charge couples through the Peierls phase [35,36]. Therefore, the terms that contribute to linear spin-wave approximation can be written as (see supplemental material Supplementarymaterial.pdf (SM))

$$\mathcal{H}(\tau) = \left( J_H + \frac{2 J_K}{3} \right) \sum_{(ij)} \left[ S_i^n S_j^n \right]$$

$$+ \frac{1}{2} \{ S_i^{+} S_j^{-} e^{i\Phi_{ij}(\tau)} + \text{H.c.} \}$$

$$+ \frac{2 J_K}{3} \sum_{(ij)\gamma} \left[ \frac{1}{2} \{ e^{i\phi} S_i^{+} S_j^{+} e^{i\Phi_{ij}(\tau)} + \text{H.c.} \} \right],$$

where $S_j^{\pm} = S_j^{x} \pm i S_j^{y}$ are the usual raising and lowering spin operators, and the angle $\phi$ comes from the rotation of the bond directions (see SM), with $\varphi_{ij} = 2\pi/3, 4\pi/3, 0$ for $x, y, z$ bond directions, respectively. The Aharonov-Casher phase acts as a vector potential or gauge field to the spin current [48]. We consider light propagating along the [111] direction (i.e., perpendicular to the honeycomb plane), given by

$$\Xi(k) = E_0 \left[ \sin(\omega \tau), \sin(\omega \tau + \phi), 0 \right],$$

where $E_0$ is the amplitude of the time-dependent electric field, $\omega$ is the angular frequency of light and $\phi$ is the polarization. Linearly and circularly polarized lights correspond to $\phi = 0$ and $\phi = \pi/2$, respectively. We perform linear spin-wave theory in the polarized phase, which is valid in the large-$S$ limit and for low-energy excitations. This can be done by writing the spin operators in eq. (3) in terms of the linearized Holstein-Primakoff bosons [54]: $S_i^n = S - a_i^\dagger a_i$, $S_j^n \approx \sqrt{2S} a_i$, for $i \in \alpha A$, and $S_n^{\pm} = S - b_j^\dagger b_j$, $S_j^{\pm} \approx \sqrt{2S} b_j$ for $j \in \alpha B$. The resulting linear spin-wave bosonic Hamiltonian is time-periodic $\mathcal{H}_2(\tau + T) = \mathcal{H}_2(\tau)$, where $T$ is the period of the driving field.

We can now implement the machinery of the Floquet theory [55], to study the dynamics of irradiated Kitaev magnets. In the off-resonant limit $\hbar \omega > A$, light simply modifies the band structures [36]. The effect of such off-resonant light is captured in a static effective Hamiltonian $\mathcal{H}_{\text{eff}}$ [35,36], defined through the evolution Floquet operator $U$ of the system after one period $T = 2\pi/\omega$ as

$$\mathcal{H}_{\text{eff}} = i \frac{\pi}{T} \log(U),$$

where $U = T \exp(-i \int_0^T \mathcal{H}_2(\tau) d\tau)$ and $T$ is the time-ordering operator. The effective Hamiltonian can be written as $H_{\text{eff}} = \sum_{i \geq 0} \mathcal{H}_{\text{eff}}^{(i)}(\hbar \omega)^i$. We work in the

off-resonant limit where the photon energy is much larger than the energy scale of the static system, i.e., $\hbar \omega \gg A$. This means that we focus on the zero-photon sector [35], $\mathcal{H}_{\text{eff}}^{(0)} = \mathcal{H}_{\text{eff}}^{(0)}$, where $\mathcal{H}_{\text{eff}}^{(0)} = \frac{1}{T} \int_0^T d\tau e^{-i\omega \tau} \mathcal{H}_2(\tau)$ represents the discrete Fourier components and $n \in \mathbb{Z}$. Next, we Fourier transform $\mathcal{H}_{\text{eff}}^{(0)}$ into momentum space and use the basis vector $|\psi^{(0)}(\vec{k})\rangle = (a_{\vec{k},\alpha A}^{(0)} \dagger, b_{\vec{k},\alpha B}^{(0)} \dagger, a_{-\vec{k},\alpha A}^{(0)}, b_{-\vec{k},\alpha B}^{(0)} \rangle)$. The effective time-independent Hamiltonian is given by

$$\mathcal{H}_{\text{eff}}^{(0)}(\vec{k}) = S \left[ \begin{array}{cc} \mathcal{M}(\vec{k}) & \mathcal{N}(\vec{k}) \\ \mathcal{N}(\vec{k})^\dagger & \mathcal{M}(-\vec{k}) \end{array} \right],$$

$$\mathcal{M}(\vec{k}) = \left( \begin{array}{cc} \rho_0^{(0)}(\vec{k}) & \rho_1^{(0)}(\vec{k}) \\ \rho_1^{(0)}(\vec{k})^* & \rho_0^{(0)}(-\vec{k}) \end{array} \right),$$

$$\mathcal{N}(\vec{k}) = \left( \begin{array}{cc} 0 & \rho_2^{(0)}(\vec{k}) \\ \rho_2^{(0)}(-\vec{k}) & 0 \end{array} \right),$$

where

$$\rho_0^{(0)} = -3 J_H - 2 J_K,$$

$$\rho_1^{(0)}(\vec{k}) = \left( J_H + 2 J_K / 3 \right) \left[ \mathcal{J}_0(\vec{k}) + \mathcal{J}_0(\mathcal{E}_+(\phi)) e^{i \vec{k} \cdot \vec{a}_1} + \mathcal{J}_0(\mathcal{E}_-(\phi)) e^{i \vec{k} \cdot \vec{a}_2} \right],$$

$$\rho_2^{(0)}(\vec{k}) = \left( 2 J_K / 3 \right) \left[ \mathcal{J}_0(\vec{k}) + \mathcal{J}_0(\mathcal{E}_+(\phi)) e^{i \vec{k} \cdot \vec{a}_1 + 2 \pi / 3} + \mathcal{J}_0(\mathcal{E}_-(\phi)) e^{i \vec{k} \cdot \vec{a}_2 - 2 \pi / 3} \right],$$

where $\mathcal{J}_0(x)$ is the Bessel function of order $\ell \in \mathbb{Z}$, and $E_{\pm}(\phi) = \frac{\mathcal{E}_0}{\sqrt{2\pi}} \sqrt{4 \pm 2 \sqrt{3} \cos \phi}$. The dimensionless quantity that characterizes the light intensity is $E_0 = g \mu_B E_0 / h c^2$. The static effective Hamiltonian in eq. (6) can be diagonalized by performing a bosonic Bogoliubov transformation (see SM).

**Fig. 2: Floquet magnon bands for FM Kitaev-Heisenberg model $\vartheta = 5\pi/4$ (top panel) and AFM Kitaev point $\vartheta = \pi/2$ (bottom panel).**

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1See Supplemental Material for detailed information on the Floquet theory, Berry curvature, and the thermal Hall effect.
magnon topology is not well defined at \( \theta \) point (BZ). The flux of the Berry curvature threading the entire Brillouin zone (BZ) defines the Chern number of the Floquet magnon bands as \( = 0 \), which implies a photoinduced magnetic order without the onset of a classical spin liquid [58]. As the laser field is an artifact of an extensive classical degeneracy, and points to the onset of a classical spin liquid [58]. By applying a laser drive, the gap at the K-point does not close, however, the system is now driven to a well-defined topological magnon insulator as we will show below. At the AFM Kitaev point \( \theta = \pi/2 \) [52,56] (bottom panel), the lowest magnon band is a zero energy mode in the undriven system for \( \theta = 0 \) [57]. The presence of zero energy mode in the spin wave excitations of frustrated magnets is an artifact of an extensive classical degeneracy, and points to the onset of a classical spin liquid [58]. As the laser field is applied, the zero energy mode is lifted for \( \phi = \pi/2 \) and \( \phi = 0 \), which implies a photoinduced magnetic order without a high applied magnetic field [12].

To investigate the magnon topology of the system, we define the Chern number of the Floquet magnon bands as the flux of the Berry curvature threading the entire Brillouin zone (BZ): \( C^0_{\text{eff}}(\theta, \phi) = \frac{1}{\pi i} \oint_{\Omega(\tilde{k})} d^2k \Omega^x_0(\tilde{k}) \), where \( \Omega^x_0(\tilde{k}) \) is the Berry curvature of the Floquet magnon bands labeled by \( \alpha = 1, 2 \) (see SM). The Chern number has been computed using the discretized BZ method [59]. In the main panel of fig. 3, we show the evolution of the lowest Floquet Chern number as a function of \( \theta_0 \) for \( \phi = 0 \) with \( \theta = 5\pi/4 \) and \( \theta = \pi/2 \), while the inset shows the Chern number for \( \phi = \pi/2 \). As we mentioned above, the magnon topology of the system is not well defined at equilibrium \( \theta_0 = 0 \), thus we do not consider this case. For \( \phi = 0 \) and \( 0 < \theta_0 \lesssim 1.35 \), where \( h(\theta_0 \sim 1.35, \phi = 0) \sim AS \) for \( \theta = \pi/2 \) (see eq. (12) below), the Chern number of the lowest band is \( C_{\text{eff}} = +1 \) in the FM Kitaev-Heisenberg model \( \theta = 5\pi/4 \) and \( C_{\text{eff}} = -1 \) at the AFM Kitaev point \( \theta = \pi/2 \); but the Chern number is zero for \( \theta_0 > 1.35 \). For \( \phi = \pi/2 \), the Chern number is nonzero provided \( \theta_0 \neq 0 \).

**Effective spin Hamiltonian in real space.** To understand the origin of the photoinduced topological magnons, we can map the off-resonant effective static Hamiltonian in eq. (6) back to the real-space spin operators keeping in mind the Holstein-Primakoff bosons. In the original cubic coordinate system, the real-space effective static spin Hamiltonian which reproduces eq. (6) is given by

\[
\mathcal{H}_{\text{eff}}^{(0)}(\theta, \phi) = \sum_{(ij) \gamma} J_{ij}(\theta_0, \phi) S^\gamma_i S^\gamma_j + \sum_{(ij)} J_{ij}(\theta_0, \phi) \vec{S}_i \cdot \vec{S}_j + h(\theta_0, \phi) \sum_i (S^x_i + S^y_i + S^z_i),
\]

which is a renormalized Kitaev-Heisenberg model plus a photoinduced magnetic field along the [111] direction. The anisotropic Kitaev interactions are given by \( J_{ij}(\theta_0) = 2J_K J_0(\theta_0) \), \( J_{ij}(\theta_0, \phi) = 2J_K J_0(\theta_0, \phi) \), and \( J_{ij}(\theta_0, \phi) = 2J_K J_0(\theta_0, \phi) \). The Heisenberg interactions are distorted with \( J_{ij}(\theta_0) = J_{ij}(\theta_0, \phi) \) along the vertical \( \delta_1 \) bond, \( J_{ij}(\theta_0, \phi) = J_{ij}(\theta_0, \phi) \) along the diagonal \( \delta_2 \) bond, and \( J_{ij}(\theta_0, \phi) = J_{ij}(\theta_0, \phi) \) along the diagonal \( \delta_2 \) bond (see fig. 1). The photoinduced magnetic field is given by

\[
h(\theta_0, \phi) = (2J_K + 3J_H)S \left[ 1 - \frac{\mathcal{J}(\theta_0, \phi)}{3} \right],
\]

where \( \mathcal{J}(\theta_0, \phi) = J_0(\theta_0) + J_0(\theta_0, \phi) + J_0(\theta_0, \phi) \). Equation (13) stems from the nonrenormalized Kitaev-Heisenberg interaction in eq. (9). Note that eq. (13) vanishes at \( \theta_0 = 0 \), hence eq. (12) reduces to eq. (1). For \( \theta_0 \neq 0 \), however, eq. (13) lies in the interval \( 0 < h(\theta_0, \phi) < (2J_K + 3J_H)S \). Thus, at the AFM Kitaev point \( \theta = \pi/2 \) (\( J_K = 0 \)), the photoinduced magnetic field is \( 0 < h(\theta_0, \phi) < 2AS \), which is much smaller than the high homogeneous magnetic field \( h > 4AS \) required to induce topological magnons in the undriven pure Kitaev model [12]. On the contrary, at the FM Heisenberg point \( \theta = \pi/2 \) (\( J_K = 0 \)), the effective Hamiltonian (12) is simply a distorted fully polarized honeycomb ferromagnet, which does not possess any topological magnon modes (see SM).

One of the hallmarks of 2D topological systems is the existence of gapless chiral edge modes on the boundary of the system [1,2]. In insulating topological magnets, the

\[3\text{Note that eq. (12) is valid in linear spin wave approximation for the magnetically ordered state considered in this paper. Conversely, the effective static spin Hamiltonian that is manifested directly from eq. (3) will be different, because no specific magnetically ordered state is assumed in eq. (9).} \]
The Heisenberg model is stable for topological magnons in magnetically ordered systems. However, it has been shown that the high magnetic-field–induced virtual-photon absorption and emission processes, which is a consequence of the Berry curvature (Chern number) of the lowest magnon band. Recent results indicate that the existence of Floquet topological magnon insulators in periodically driven pure Kitaev model and ferromagnetic phase of the Kitaev-Heisenberg model at zero magnetic field. The main result of our study can be summarized as follows. In the off-resonant limit, the Floquet physics stabilizes magnetic order and the effective time-independent Hamiltonians map onto the corresponding anisotropic static spin model, plus a tunable photoinduced magnetic field along the [111] direction, which facilitates the existence of Floquet topological magnon modes in a similar fashion to a homogenous magnetic field in the undriven systems [12,13]. One of the advantages of the current results is that the photoinduced topological magnons and the chiral edge modes can be tuned by varying the amplitude and polarization of the laser field. Another interesting feature of irradiated Kitaev magnets is that the existence of the Floquet topological magnon insulator does not require the explicit time-reversal symmetry breaking term from the second-order virtual-photon absorption and emission processes, which is mandatory for the existence of Floquet topological states in irradiated graphene [36,37] and irradiated honeycomb ferromagnets [42–44]. We also showed that irradiated Kitaev magnets exhibit a tunable photoinduced thermal Hall effect. A direct experimental implication of the current proposal is that ultrafast magnon spin currents can be generated in irradiated Kitaev materials using different experimental techniques such as the inverse Faraday effect [49] and THz spectroscopy [51]. This could pave the way for topological opto-magnonics and opto-spintronics [34] using Kitaev materials.

In future work, we plan to address the effect of magnon-magnon interactions and see how they modify eq. (12). However, it has been shown that the high magnetic-field–induced undriven topological magnons and chiral edge modes present in the linear spin-wave approximation of magnons is close to thermal equilibrium. In this limit, the thermal Hall effect mimics that of equilibrium systems where a longitudinal temperature gradient $-\partial_y T$ induces a transverse heat current $J_y^{\parallel} = -\kappa_{xy} \partial_y T$, where $\kappa_{xy}$ is the thermal Hall conductivity, derived in ref. [63] (see SM (see footnote 1)). In fig. 5, we show the $\varepsilon_0$-dependence of $\kappa_{xy}$ for $\phi = 0$ and $T/A = 0.3, 0.35, 0.4$, in the FM Kitaev-Heisenberg model $\theta = 5\pi/4$ and at the AFM Kitaev point $\theta = \pi/2$ (inset). We note that $\kappa_{xy}$ is ill-defined for $\varepsilon_0 = 0$ at low temperatures (not shown). The thermal Hall conductivity is dominated by the Berry curvature of the lowest magnon band at low temperatures and its sign is consistent with the sign of the Berry curvature (Chern number) of the lowest magnon band. At low temperature $T/A \ll 1$ and for $\varepsilon_0 > 1.35$, $\kappa_{xy}$ is very small and approaches zero consistently with the vanishing of the Chern number and the absence of traversing chiral edge modes for $\phi = 0$ as shown above. The low-temperature dependence of $\kappa_{xy}$ for $\phi = 0$ is shown in the SM.

**Conclusion and outlook.** – We have proposed the existence of Floquet topological magnon insulators in periodically driven pure Kitaev model and ferromagnetic phase of the Kitaev-Heisenberg model at zero magnetic field. The main result of our study can be summarized as follows. In the off-resonant limit, the Floquet physics stabilizes magnetic order and the effective time-independent Hamiltonians map onto the corresponding anisotropic static spin model, plus a tunable photoinduced magnetic field along the [111] direction, which facilitates the existence of Floquet topological magnon modes in a similar fashion to a homogeneous magnetic field in the undriven systems [12,13]. One of the advantages of the current results is that the photoinduced topological magnons and the chiral edge modes can be tuned by varying the amplitude and polarization of the laser field. Another interesting feature of irradiated Kitaev magnets is that the existence of the Floquet topological magnon insulator does not require the explicit time-reversal symmetry breaking term from the second-order virtual-photon absorption and emission processes, which is mandatory for the existence of Floquet topological states in irradiated graphene [36,37] and irradiated honeycomb ferromagnets [42–44]. We also showed that irradiated Kitaev magnets exhibit a tunable photoinduced thermal Hall effect. A direct experimental implication of the current proposal is that ultrafast magnon spin currents can be generated in irradiated Kitaev materials using different experimental techniques such as the inverse Faraday effect [49] and THz spectroscopy [51]. This could pave the way for topological opto-magnonics and opto-spintronics [34] using Kitaev materials.

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remain intact in the presence of magnon-magnon interactions \cite{12}. We also plan to study the nonequilibrium distribution function \cite{66,67} of magnons in this system. Moreover, it would also be interesting to investigate whether tunable topological magnons can be photoinduced in the zigzag phase of the Kitaev-Heisenberg model.

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