

Floquet Spin Splitting and Spin Generation in Antiferromagnets

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In antiferromagnetic spintronics, accessing the spin degrees of freedom is essential for generating spin currents and manipulating magnetic order, which generally requires lifting spin degeneracy. This is typically achieved through relativistic spin-orbit coupling or nonrelativistic spin splitting in altermagnets. Here, we propose an alternative approach: a dynamical spin splitting induced by an optical field in antiferromagnets. By coupling the driven system to a thermal bath, we demonstrate the emergence of steady-state pure spin currents as well as linear-response longitudinal and transverse spin currents. Crucially, thermal bath engineering enables a nonrelativistic Edelstein effect—the generation of a net spin accumulation—without relying on spin-orbit coupling. Our results provide a broadly applicable and experimentally tunable route to control spins in antiferromagnets, offering new opportunities for spin generation and manipulation in antiferromagnetic spintronics.

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The frontier of present spintronics research largely focuses on antiferromagnetic systems, due to their ultrafast dynamics and being free of stray magnetic fields [1–5]. However, in many collinear antiferromagnets, the electronic bands remain spin degenerate due to a preserved antiunitary effective time-reversal symmetry, typically realized as the combination of time reversal with spatial inversion or a sublattice translation. This symmetry protection imposes a fundamental constraint on antiferromagnetic spintronics, where direct manipulation of real spin degrees of freedom is essential. To access the spin degrees of freedom, spin-orbit coupling (SOC) is typically required, enabling a variety of spin-related phenomena, such as spin generation, spin Hall effect, and spin-orbit torques [6,7].

Recently, a new class of magnetic materials with nonrelativistic spin splitting, dubbed altermagnets, has gathered extensive attention [8–24]. In altermagnets, the effective time reversal symmetry is absent because the involved lattice symmetry is intrinsically broken by the lattice structure. More importantly, the nonrelativistic spin splitting is typically larger than that arising from SOC, which is favorable for utilizing the spin degrees of freedom.

Another route to lifting spin degeneracy is by explicitly breaking time-reversal symmetry, as in the case of Zeeman splitting induced by a magnetic field. However, the

resulting energy scales are typically negligible compared to the electronic band structure (~ 1 eV). For example, a 1 T magnetic field yields a splitting on the order of $10^{-4} \sim 10^{-5}$ eV. As an alternative, time-reversal symmetry can be broken dynamically using an optical field, offering a more efficient means to lift spin degeneracy. In this Letter, we investigate light-induced spin splitting in antiferromagnets possessing effective time-reversal symmetry. Using Floquet theory [25,26], we demonstrate that an optical field with suitable intensity and frequency can induce substantial spin splitting in the electronic quasenergy bands via nonequilibrium effects. With appropriate thermal bath engineering, the driven system can support SOC-independent steady-state spin and charge currents as well as spin accumulation. Our findings uncover an experimentally tunable mechanism for controlling the spin degree of freedom in spin-degenerate antiferromagnets, opening new avenues for spin generation and manipulation in antiferromagnetic spintronics.

Floquet spin splitting in an antiferromagnet (AFM)—We consider an electronic system with collinear antiferromagnetic order on honeycomb lattice, as realized in materials such as MnPX_3 ($X = \text{S, Se}$) [27,28]. The momentum-space Hamiltonian matrix under basis $(c_{\mathbf{k}A,\uparrow}, c_{\mathbf{k}A,\downarrow}, c_{\mathbf{k}B,\uparrow}, c_{\mathbf{k}B,\downarrow})^T$ is given by

$$H(\mathbf{k}) = \begin{pmatrix} \lambda_{\text{SO}} \xi_{\mathbf{k}} \sigma^z + \lambda \mathbf{n} \cdot \boldsymbol{\sigma} & t \gamma_{\mathbf{k}} \\ t \gamma_{\mathbf{k}}^* & -\lambda_{\text{SO}} \xi_{\mathbf{k}} \sigma^z - \lambda \mathbf{n} \cdot \boldsymbol{\sigma} \end{pmatrix} \quad (1)$$

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TABLE I. Expression of parameters in the Fourier transformed Hamiltonian, where $\mathcal{N}_{\pm} = 1 \pm \frac{\sqrt{3}}{2} \cos \varphi$.

θ_1	$\theta_{2,3}$	ζ_1	$\zeta_{2,3}$
0	$\mp \text{sign}(\pi - \varphi) \cos^{-1} \left[\frac{1}{2} \mp \frac{\sqrt{3}}{2} \cos \varphi / \sqrt{\mathcal{N}_{\mp}} \right]$	1	$-\sqrt{\mathcal{N}_{\mp}}$
$\tilde{\theta}_1$	$\tilde{\theta}_{2,3}$	$\tilde{\zeta}_1$	$\tilde{\zeta}_{2,3}$
$-\varphi$	$\text{sign}(\pi - \varphi) \cos^{-1} \left[\mp \frac{3}{2} - \frac{\sqrt{3}}{2} \cos \varphi / \sqrt{3\mathcal{N}_{\pm}} \right]$	$\sqrt{3}$	$\sqrt{3\mathcal{N}_{\pm}}$

where $\gamma_{\mathbf{k}} = \sum_{i=1}^3 e^{-i\mathbf{k} \cdot \delta_i}$ with $\delta_1 = a(1, 0)$, $\delta_{2,3} = a(-1/2, \pm\sqrt{3}/2)$ (here, a is lattice constant), and $\xi_{\mathbf{k}} = \sum_{i=1}^3 2 \sin(\mathbf{k} \cdot \boldsymbol{\kappa}_i)$ with $\boldsymbol{\kappa}_1 = a(0, \sqrt{3})$, $\boldsymbol{\kappa}_{2,3} = a(\mp 3/2, -\sqrt{3}/2)$. Here, λ_{SO} is the intrinsic SOC, λ represents the exchange coupling, σ^i is the Pauli matrix acting in spin space, and $\mathbf{n}$ is the Néel vector. We assume $\mathbf{n} \parallel \hat{z}$ to preserve spin conservation. The system preserves combined parity-time ($\mathcal{PT}$) symmetry, ensuring spin degeneracy throughout the band structure.

To break the $\mathcal{PT}$ symmetry, we apply polarized light described by a time-dependent vector potential $\mathcal{A} = A_0[\sin \omega t, \sin(\omega t + \varphi), 0]$, where $A_0 = E_0/\omega$ (set $e = \hbar = 1$) with E_0 being the electric field amplitude. Under inversion $\mathcal{P}$ and time reversal $\mathcal{T}$, the vector potential transforms as $\mathcal{P}: \mathcal{A}(t) \rightarrow -\mathcal{A}(t)$, and $\mathcal{T}: \mathcal{A}(t) \rightarrow -\mathcal{A}(-t)$, which together yield $\mathcal{PT}: \mathcal{A}(t) \rightarrow \mathcal{A}(-t)$. Elliptical or circular polarization ($\varphi \neq 0, \pi$) explicitly breaks the $\mathcal{PT}$ symmetry. In contrast, for linear polarization ($\varphi = 0, \pi$), $\mathcal{PT}$ enforces $\mathcal{A}(-t) = -\mathcal{A}(t)$, which is equivalent to a time translation by π/ω . As a result, all physical observables remain invariant over one driving period, and the $\mathcal{PT}$ symmetry is effectively preserved. The light couples to the system via the Peierls substitution $H(\mathbf{k}) \rightarrow H(t) = H[\mathbf{k} + \mathcal{A}(t)]$. To preserve antiferromagnetic order, we focus on a subgap driving regime $J_{\text{ex}} \ll \hbar\omega \ll \Delta_c$, where J_{ex} is exchange energy between local spins and Δ_c the charge gap associated with their formation [25,29–31]. Typically, $J_{\text{ex}} \sim t_h^2/U$ and $\Delta_c \sim U$, where t_h is the electron hopping energy and U refers to local repulsive Coulomb energy, satisfying $t_h \ll U$. Taking $t_h \lesssim t$, a drive with $\hbar\omega \sim t$ lies within this subgap window. In this regime, magnon absorption is suppressed, and heating via charge excitations occurs only on exponentially long time scales ($\sim e^{U/\hbar\omega}$), which can be further mitigated by coupling to a thermal bath [32–34]. By contrast, antiferromagnetic order may also survive in the ultrafast regime $\hbar\omega \gg U \gg t$, but the resulting spin splitting is perturbatively small.

The resulting periodically driven system is naturally analyzed within the framework of Floquet theory [26,35]. The eigenstate is represented as a Floquet state: $|\psi_n(t)\rangle = e^{-i\varepsilon_n t/\hbar} |\phi_n(t)\rangle$, where ε_n is the quasienergy, and $|\phi_n(t+T)\rangle = |\phi_n(t)\rangle$ with $T = 2\pi/\omega$. The periodic part of the Floquet state respects $(\varepsilon_n + i\partial_t)|\phi_n(t)\rangle = H(t)|\phi_n(t)\rangle$. This equation can be further translated

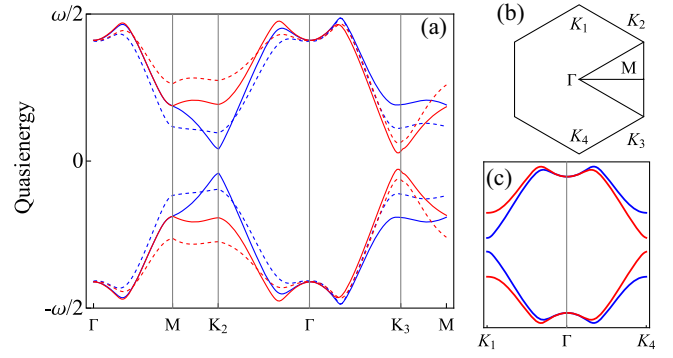


FIG. 1. (a) Quasienergy band structure in honeycomb-lattice AFM with Néel order, where blue (red) color stands for spin up (down), and solid (dashed) line represents the case without (with) SOC (with $\lambda_{\text{SO}} = 0.2$). (b) The Brillouin zone. (c) Quasienergy band structure along $K_1 - \Gamma - K_4$, where SOC is zero. In (a) and (c), $\varphi = \pi/3$, $A_0 a = 1$, $\omega = 4$, and $t = 1$, $\lambda = 0.5$.

to an equation of associated Fourier components: $(\varepsilon_n + m\omega)|\phi_n^{(m)}\rangle = \sum_{m'} H^{(m-m')} |\phi_n^{(m')}\rangle$, where $|\phi_n(t)\rangle = \sum_m e^{-im\omega t} |\phi_n^{(m)}\rangle$ and $H^{(m)} = (1/T) \int_0^T dt e^{im\omega t} H(t)$. The quasienergy ε_n is well defined up to $m\omega$ ($m \in \text{Integers}$). Therefore, it is enough to confine the quasienergy to the first “Floquet-Brillouin zone” (FBZ), $-\omega/2 \leq \varepsilon_n < \omega/2$.

It is straightforward to obtain the Fourier component of Hamiltonian (1):

$$H^{(m)} = \begin{pmatrix} 0 & h^{(m)} \\ (h^{(-m)})^* & 0 \end{pmatrix} \sigma^0 + h_{\text{SO}}^{(m)} \tau^z \sigma^z + \lambda \tau^z \sigma^z \delta_{m,0}, \quad (2)$$

where τ^i is the Pauli matrix in the sublattice space. Here, $h^{(m)} = \sum_{i=1}^3 e^{-i\mathbf{k} \cdot \delta_i} e^{-im\theta_i} J_m(\zeta_i A_0 a)$ and $h_{\text{SO}}^{(m)} = -i\lambda_{\text{SO}} \times \sum_{i=1}^3 [e^{i\mathbf{k} \cdot \boldsymbol{\kappa}_i} (-1)^m - e^{-i\mathbf{k} \cdot \boldsymbol{\kappa}_i}] e^{im\theta_i} J_m(\tilde{\zeta}_i A_0 a)$, where $J_m(\dots)$ is the m th Bessel function; all involved parameters are listed in Table I. The quasienergy band structure can be obtained by applying Eq. (2) to the Fourier-transformed eigenequation. In this system, spin remains a good quantum number, allowing the band structure and corresponding transport properties to be analyzed within spin-resolved subspaces.

The application of light is expected to lift the spin degeneracy of the quasienergy bands. Indeed, the band structure in Fig. 1 is spin nondegenerate, even in the absence of SOC. Importantly, the spin-split energy is surprisingly large, comparable to the scale of the original band structure and exceeding the typical SOC. This offers an efficient way to approach the spin degrees of freedom. It is worth noting that if $\lambda_{\text{SO}} = 0$, the Hamiltonian Eq. (1) has a dual symmetry between the two spin sectors:

$$\tau^x H_{\uparrow}[\mathbf{k} + \mathcal{A}(t)] \tau^x = H_{\downarrow}[-\mathbf{k} - \mathcal{A}(t)]. \quad (3)$$

This symmetry originates from the combined action of spatial inversion ($\mathcal{P}$) and a twofold spin rotation about the

x -axis C_{2x}^{spin} , such that $\mathcal{P}C_{2x}^{\text{spin}}H[\mathbf{k} + \mathcal{A}(t)](\mathcal{P}C_{2x}^{\text{spin}})^{-1} = H[\mathbf{k} + \mathcal{A}(t)]$, where the inversion exchanges the sublattices (represented by τ^x) and transforms $\mathbf{k} + \mathcal{A}$ to $-\mathbf{k} - \mathcal{A}$ and the spin rotation flips spin σ^z to $-\sigma^z$. The dual symmetry leads to a dual relation between quasienergies: $\varepsilon_{u,d}^{\uparrow}(\mathbf{k}) = \varepsilon_{u,d}^{\downarrow}(-\mathbf{k})$ [e.g., see Fig. 1(c)], where u, d refer to the up and down bands in the FBZ. Notice that in Fig. 1, the equivalence among valleys is absent because the threefold rotation symmetry of the original model is broken by the optical field. As a result, the dual relation between quasienergies does not apply to the path $\mathbf{K}_3 - \Gamma - \mathbf{K}_2$. On the other hand, a nonzero intrinsic SOC can explicitly break the dual symmetry in Eq. (3).

The spin splitting is a cooperative effect of exchange coupling and light. To see this in a clean way, we investigate the special case of off-resonance ($\omega \gg t$) and weak driving field ($A_0 a \ll 1$), for which the original bands are dressed by the light to yield a spin-split term. We focus on the vicinity near the valley $\mathbf{K}_4(\mathbf{K}_1) = [0, \mp(4\pi/3\sqrt{3})]$ point, where $\mathcal{H}_v(t) = v_F[\tau^x \eta^z(q_y + \mathcal{A}_y) + \tau^y(q_x + \mathcal{A}_x)] + s\lambda\tau^z$, and λ_{SO} is set to zero for simplicity. Here, $v_F = 3at/2$, $\mathbf{q} = \mathbf{k} - \mathbf{K}_4$ (or $\mathbf{q} = \mathbf{k} - \mathbf{K}_1$), η^i denotes the Pauli matrix for valley freedoms, and $s = \pm 1$ for up and down spins. The effective Floquet Hamiltonian [36,37] up to $O[(aA_0)^4]$ is

$$H_{\text{eff}}^F \approx \mathcal{H}_v^{(0)} + \frac{[\mathcal{H}_v^{(-1)}, \mathcal{H}_v^{(+1)}]}{\omega} \\ = v_F(\tau^x \eta^z q_y + \tau^y q_x) + s\lambda\tau^z - M\tau^z \eta^z \quad (4)$$

where $M = [(v_F A_0)^2 / \omega] \sin \varphi$, and $\mathcal{H}_v^{(\pm)}$ is the Fourier component of $\mathcal{H}_v(t)$. It is easy to obtain $\varepsilon_{\mathbf{K}_4(\mathbf{K}_1)}^s(\mathbf{q}) = \pm \sqrt{v_F^2 q^2 + (s\lambda - \chi M)^2}$ with $\chi = \pm 1$ being the valley index. It is clear that the quasienergy at each valley becomes spin dependent, driven by the combined effects of exchange coupling and optical driving. When the intrinsic SOC is considered, M is shifted to $M - s3\sqrt{3}\lambda_{\text{SO}}$, thus breaking the dual relation $\varepsilon_{\mathbf{K}_1}^s = \varepsilon_{\mathbf{K}_4}^{-s}$.

In general, a periodically driven isolated system tends toward an infinite-temperature state due to energy absorption from the drive. To avoid this heating problem and reach a steady state, the system must be coupled to a heat bath. Interestingly, we find that the nature of spin transport or accumulation depends sensitively on the type of bath: distinct spin-related behaviors emerge when coupling to bosonic versus fermionic reservoirs.

Steady state and spin currents—We first consider electron coupling to a bosonic phonon bath, described by the electron-phonon Hamiltonian

$$H_{ep} = \sum_{\mathbf{k}, \mathbf{q}} \sum_{\nu, \nu'} \sum_{\lambda} g_{\nu\nu'}^{\lambda}(\mathbf{k} - \mathbf{q}, \mathbf{k}) c_{\mathbf{k}-\mathbf{q}, \nu'}^{\dagger} c_{\mathbf{k}, \nu} (b_{\mathbf{q}\lambda}^{\dagger} + b_{-\mathbf{q}\lambda}). \quad (5)$$

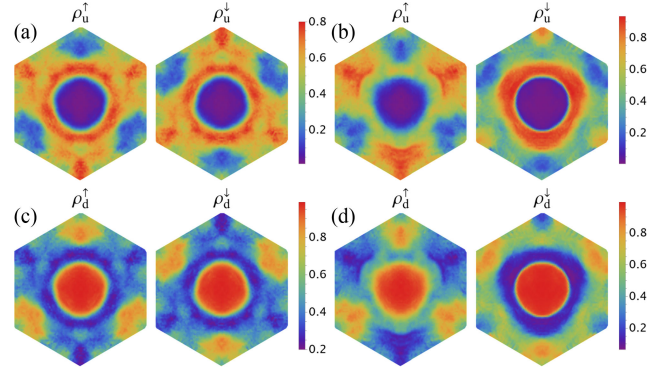


FIG. 2. Steady-state population for the quasienergy band (in the first FBZ) in each spin sector. The first and second rows display the population for upper and lower bands, respectively. In (a) and (c), $\lambda_{\text{SO}} = 0$; in (b) and (d), $\lambda_{\text{SO}} = 0.1$. Other parameters are $\varphi = \pi/2$, $A_0 a = 1$, $\omega = 1$, and $t = 1$.

Here, ν, ν' label the sublattices, $b_{\mathbf{q}\lambda}^{\dagger}(b_{-\mathbf{q}\lambda})$ is the creation (annihilation) operator for λ phonon modes, and $g_{\nu\nu'}^{\lambda}(\mathbf{k} - \mathbf{q}, \mathbf{k})$ is the electron-phonon coupling matrix [38,39]. The phonon modes are described by $H_{ph} = \sum_{\lambda, \mathbf{k}} \hbar\omega_{\lambda\mathbf{k}} b_{\mathbf{q}\lambda}^{\dagger} b_{\mathbf{q}\lambda}$. The steady-state occupation can be obtained by numerically solving the kinetic equation

$$\partial_t \rho_{\mathbf{k}\alpha} = \sum_{\mathbf{k}'\alpha'} W_{\mathbf{k}\alpha, \mathbf{k}'\alpha'} (1 - \rho_{\mathbf{k}, \alpha}) \rho_{\mathbf{k}', \alpha'} - W_{\mathbf{k}'\alpha', \mathbf{k}\alpha} \\ \times (1 - \rho_{\mathbf{k}'\alpha'}) \rho_{\mathbf{k}\alpha}, \quad (6)$$

where $\rho_{\mathbf{k}, \alpha}$ denotes the occupation number for α quasiband at momentum $\mathbf{k}$, and $W_{\mathbf{k}\alpha, \mathbf{k}'\alpha'}$ is the scattering rate determined by the electron-phonon coupling and the Floquet eigenstates; see details in Ref. [40,41].

The steady-state distribution obtained from Eq. (6) deviates markedly from the conventional Fermi-Dirac form. In particular, the initially fully occupied lower Floquet bands become partially depleted, while the empty upper bands acquire finite occupation, as shown in Fig. 2. This redistribution leads to an effectively metallic steady state despite the presence of an initial band gap, as evident in the finite conductivity shown in Fig. 3(c). Notably, the dual symmetry in Eq. (3) imposes a corresponding symmetry in the momentum-resolved occupations: $\rho_{u,d}^{\uparrow}(\mathbf{k}) = \rho_{u,d}^{\downarrow}(-\mathbf{k})$; see Figs. 2(a) and 2(c). This duality in population plays a key role in enabling spin transport, as we discuss below.

In the nonequilibrium steady state, a nonvanishing current is allowed to exist. The averaged spin current density over one period is given by $\bar{\mathbf{j}}^s = \bar{\mathbf{j}}^{\uparrow} - \bar{\mathbf{j}}^{\downarrow}$, with

$$\bar{\mathbf{j}}^{\sigma} = \sum_n \int \frac{d^2 \mathbf{k}}{(2\pi)^2} \rho_n^{\sigma}(\mathbf{k}) \partial_{\mathbf{k}} \varepsilon_{n\mathbf{k}}^{\sigma}, \quad (7)$$

where $\sigma = \uparrow, \downarrow$ [41]. The inversion symmetry in each spin sector is broken; see Figs. 1(c) and 2. In the absence of

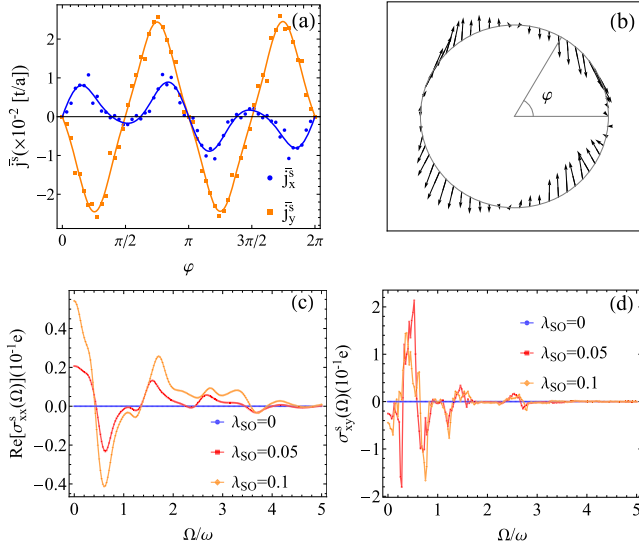


FIG. 3. (a) Steady-state spin current with $\lambda_{SO} = 0$, where the unit $t/a \sim 10^9$ eV/m. (b) The diagram for the current direction and magnitude. (c),(d) Longitudinal and transverse optical conductivity in honeycomb AFM with $\varphi = \pi/2$. In the plots, other parameters are $A_0 a = 1$, $\omega = 1$, $t = 1$, $\lambda = 0.5$, and $T_{ph} = 0.01t$.

intrinsic SOC, the dual symmetry [Eq. (3)] ensures that the spin-resolved currents are equal in magnitude and opposite in direction, resulting in a pure spin current with vanishing net charge current. When intrinsic SOC is introduced, this dual symmetry is lifted, allowing both spin and charge currents to coexist. The magnitude and direction of the steady-state spin current in the SOC-free case are shown in Figs. 3(a) and 3(b) as functions of the light polarization angle, revealing a high degree of tunability via the optical field.

In Fig. 3(a), the oscillating spin current exhibits two characteristic symmetry properties. (i) The current is antisymmetric about $\varphi = \pi$, i.e., $\bar{\mathbf{j}}^s(2\pi - \varphi) = -\bar{\mathbf{j}}^s(\varphi)$. (ii) In the interval $[0, \pi]$, the components $\bar{j}_x^s$ and $\bar{j}_y^s$ are, respectively, symmetric and antisymmetric about $\pi/2$, satisfying $\bar{j}_x^s[(\pi/2) - \varphi] = \bar{j}_x^s[(\pi/2) + \varphi]$ and $\bar{j}_y^s[(\pi/2) - \varphi] = -\bar{j}_y^s[(\pi/2) + \varphi]$. Property (i) follows from the relation $H_\sigma[\mathbf{k} + \mathcal{A}(2\pi - \varphi, t)] = T H_\sigma^*[-\mathbf{k} + \mathcal{A}(\varphi, t)] T^{-1}$, which implies $\varepsilon_{n\mathbf{k}}^\sigma(2\pi - \varphi) = \varepsilon_{n,-\mathbf{k}}^\sigma(\varphi)$ and $\rho_n^\sigma(\mathbf{k}, 2\pi - \varphi) = \rho_n^{\sigma*}(-\mathbf{k}, \varphi)$; inserting these into Eq. (7) directly yields the antisymmetry. Property (ii) originates from the $\mathcal{M}_x T$ symmetry of the lattice, where $\mathcal{M}_x$ denotes mirror reflection about the x -axis. Under $\mathcal{M}_x T$, the gauge field transforms as $\mathcal{A}[(\pi/2) + \varphi, t] \rightarrow \mathcal{A}[(\pi/2) - \varphi, t]$, while the spin current transforms as $(j_x^s, j_y^s) \rightarrow (j_x^s, -j_y^s)$, enforcing the symmetry relations in (ii).

Beyond the steady-state spin current, linear-response spin transport induced by a weak external electric field provides an additional degree of control over spin manipulation [35,44,45], particularly in systems where inversion

symmetry prohibits a steady-state current. Unlike the steady-state response, the linear optical spin conductivity can be strongly constrained by the dual symmetry in Eq. (3), which enforces cancellation between contributions from the two spin sectors. Consequently, a finite response requires the inclusion of intrinsic spin-orbit coupling. Figures 3(c) and 3(d) show the computed longitudinal and transverse spin conductivities, respectively, where the oscillatory behavior originates from resonances between the external ac field frequency Ω and the Floquet quasi-energy gaps. Notably, both components are finite in the dc limit, despite the fact that the undriven, half-filled system is a band insulator. This indicates that the interplay between periodic driving and phonon-mediated relaxation leads to an effectively metallic steady state, enabling finite spin transport under small electric fields.

Nonrelativistic Edelstein effect—The spin-conserving nature of electron-phonon coupling prohibits net spin accumulation, as particle numbers in each spin sector are conserved. This constraint can be lifted by coupling the system to a fermionic reservoir through tunneling terms of the form $c_i^\dagger d_j$, which enables exchange of particles between the system ($c_i^\dagger$) and the leads (d_j). In the presence of a spin-split Floquet band structure, effective spin nonconservation then emerges as an imbalance between spin-up and spin-down particle inflow into the driven system, giving rise to a finite spin accumulation [46]. Importantly, spin remains locally conserved, and spin current is well defined in this mechanism, in sharp contrast to scenarios where spin-orbit coupling explicitly breaks local spin conservation [47].

To model the bath engineering effect, we consider a system coupled to two fermionic electrodes on its left and right sides, characterized by chemical potentials μ_L and μ_R , respectively. The coupling to each lead is described by the parameters Γ_L and Γ_R . The resulting spin-related phenomena are analyzed using Floquet-Keldysh theory [36,41,48–50].

When analyzing spin accumulation, the dual symmetry expressed in Eq. (3) remains crucial. If the two leads are symmetric, i.e., $\mu_L = \mu_R$, the dual symmetry enforces vanishing spin accumulation even through the spin degeneracy and particle conservation are broken, necessitating finite SOC to induce spin accumulation [Fig. 4(a)], where the oscillatory pattern is dictated by the same symmetry that governs the steady-state spin current discussed before [41]. Remarkably, the dual symmetry can also be broken by introducing asymmetry between the leads. For example, in Fig. 4(b), we consider leads with chemical potentials $\mu_{L/R} = \mu_0 \pm V/2$. This asymmetry lifts the dual symmetry and leads to a sizable spin accumulation, which emerges as a linear response to the applied voltage V . The sign reversal of the spin signal serves as a hallmark of bath engineering. This SOC-independent effect is known in the literature as the nonrelativistic Edelstein effect [51–54].

Experimental feasibility—To assess experimental feasibility, we consider the following parameters: $\omega = 1$ eV

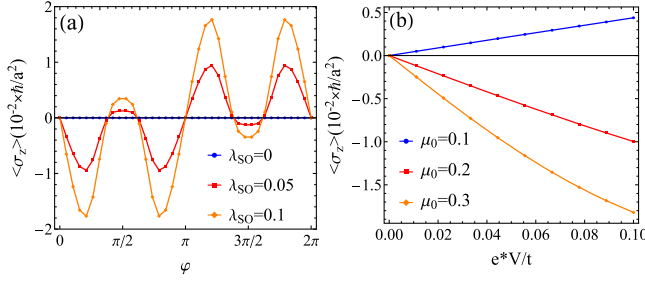


FIG. 4. Spin accumulation by contacting the system to electrodes with chemical potential μ_L, μ_R on its left and right. (a) Two leads are symmetric with $\mu_L = \mu_R = 0$. (b) Voltage induced spin accumulation, where $\mu_{L/R} = \mu_0 \pm V/2$ and $\lambda_{SO} = 0$, $\varphi = \pi/2$. Other parameters used in the calculations are $t = 1$, $A_0 a = 1$, $\lambda = 0.5$, $\Gamma_L = \Gamma_R = 0.1t$, and $\omega = 1$; the system contains 10 unit cells along the longitudinal direction.

(set $t = 1$ eV), $A_0 a = 1$, and $a = 5$ Å. These correspond to an electric field strength $E_0 = A_0 \omega / e = 2 \times 10$ MV/cm, resulting in a laser intensity $I \simeq 5 \times 10^{11}$ W/cm², which lies well within experimentally accessible regimes [36,55]. The input power delivered to the system can be estimated as $P_{in} = (1/T) \int_0^T dt \mathbf{E}(t) \cdot \mathbf{j}^e(t) \sim E_0 |\mathbf{j}^e| \sim 10^8 - 10^9$ W/cm², where $\mathbf{E}(t) = -\partial_t \mathcal{A}$ is the electric field, and $\mathbf{j}^e$ is the charge current. The current can be estimated as $\mathbf{j}^e \sim \frac{e}{\hbar} \mathbf{j}^{\uparrow/\downarrow} \approx 10^{-5} \sim 10^{-6} (e/\hbar) [\text{eV/m}]$ [see Fig. 3(a)] in each spin sector. In the steady state, this input power must be balanced by dissipation. Taking phonons as a representative bath, the scattering rate is estimated as $R \sim 2\pi g_0^2 D_{ph} / \hbar \sim 10^{13} \text{ s}^{-1}$, with electron-phonon coupling strength $g_0 \approx 0.05$ eV and phonon density of state $D_{ph} \sim 1 \text{ eV}^{-1}$ [41]. The corresponding dissipated power is roughly $P_{diss} \sim R[(D_{ph} \Delta E) n_{ph} / S_{u.c.}] \varepsilon_{ph} \sim 10^7 - 10^9$ W/cm², where $\Delta E \approx 0.1 \sim 1 \hbar \omega$ is the phonon energy window, $n_{ph} \sim 1$ is the number of relevant phonon modes, $S_{u.c.} \sim a^2$ is the area of unit cell, and $\varepsilon_{ph} \approx 0.1 \sim 1 \hbar \omega$ is the phonon energy. From these estimates, the power balance condition $P_{in} = P_{diss}$ appears feasible. In practice, additional dissipation channels—such as external reservoirs, substrates, or magnons—can further help mitigate heating in driven systems.

Conclusion—In recent years, there has been a great interest in magnetic materials with compensated magnetic order but without spin degeneracy. We propose a new, nonrelativistic way to lift the Kramers spin degeneracy in antiferromagnetic systems by using light. The optically induced spin splitting is useful for generating both non-perturbative (steady-state) and perturbative (linear-response) spin currents as well as spin accumulation. Our proposal uncovers the great potential of the optical method for spin generation in antiferromagnetic spintronics. Experimental confirmation and application of our prediction are highly feasible, given that the proposal is generically valid for various antiferromagnetic materials.

In the future, optical engineering for magnetic torque will be an intriguing direction. Moreover, controlling spin-related transport via tuning the thermal bath is also a very attractive idea.

Note added—Recently, we became aware of two relevant studies by S. Huang et al. [56] and T. Zhu et al. [57], which report findings partially overlapping with ours.

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Data availability—The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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End Matter

Other examples—We have illustrated the optically induced spin splitting and associated spin generation phenomena using a prototypical honeycomb antiferromagnet, but the underlying mechanism applies broadly to antiferromagnetic systems with appropriate symmetries, as demonstrated below.

First, we consider an AFM model on square lattice with nonsymmorphic symmetry (possible materials include SrMnPb , SrMnSn , etc. [58]), which is discussed in the context of antiferromagnetic Dirac semimetal [58–60]. The Hamiltonian reads $H(\mathbf{k}) = H_0(\mathbf{k}) + H'(\mathbf{k})$, where

$$H_0(\mathbf{k}) = -2t \cos \frac{k_x}{2} \cos \frac{k_y}{2} \tau^x - t'(\cos k_x + \cos k_y) + \lambda \tau^z \boldsymbol{\sigma} \cdot \mathbf{n}, \quad (\text{A1})$$

$$H'(\mathbf{k}) = w \sin \frac{k_y}{2} \cos \frac{k_x}{2} \tau^y. \quad (\text{A2})$$

Here, H_0 describes the square-lattice model with AFM order, and H' reflects the anisotropic nearest hopping strength [58]. As shown in Fig. 5(a), the light-induced quasienergy bands are spin-split even without involving

SOC. It is worth noting that H' breaks the inversion symmetry, which is important for the spin splitting [41].

Another example with PT symmetry is the minimal model for tetragonal CuMnAs [61,62],

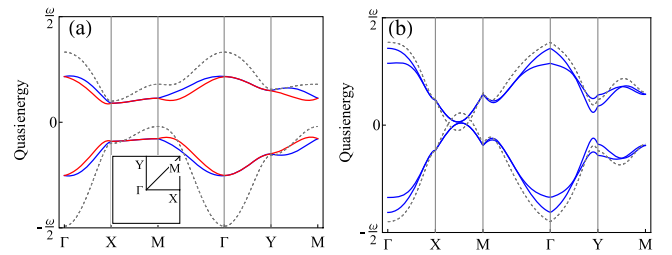


FIG. 5. Quasiband structure in the first Floquet Brillouin zone. (a) The quasienergy band of a nonsymmorphic AFM model, where blue (red) color denotes spin-up (-down) bands, and parameters are $\varphi = \pi/2$, $A_0 a = 2$, $\omega = 5$, $t = 1$, $t' = 0.2$, $w = 0.6$, and $\lambda = 0.5$. (b) Quasienergy bands of the minimal model of tetragonal CuMnAs (each band is not spin-resolved), where $\varphi = \pi/2$, $A_0 a = 1$, $\omega = 5$, $t = 1$, $t' = 0.08$, $\lambda = 0.6$, $\alpha_R = 0.8$, and $\mathbf{n} = (1, 0, 0)$. In both (a) and (b), the dashed gray lines represent the original degenerate bands, and they share the same Brillouin zone; see the inset of (a).

$$H(\mathbf{k}) = H_0(\mathbf{k}) + \alpha_R \tau^z (\sigma^y \sin k_x - \sigma^x \sin k_y), \quad (\text{A3})$$

where H_0 is given by Eq. (A1), and α_R is the Rashba spin-orbit coupling parameter. In contrast to the above examples, even with Kramers degeneracy, spin is not conserved here.

Therefore, optical driving not only removes the spin degeneracy [see Fig. 5(b)] but also affects the spin texture in each quasienergy band; see Refs. [41,63]. This offers an opportunity for engineering Néel torque, which is an interesting direction for future study.