

Emergent Gauge Fields and Linear Dichroism in Spin-Textured Altermagnets

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Spin textures are ubiquitous in antiferromagnets, yet their consequences for altermagnets remain largely unexplored. We show that smooth spatial variations of the Néel order act on itinerant electrons as emergent gauge fields that couple to the altermagnetic sublattice-odd channel. For a coplanar spin helix, this channel generates two responses that vanish when it is removed: a momentum-resolved sublattice-odd spectral structure and a net interband linear dichroism whose principal axes are controlled by the helix wave vector, so the optical response can be steered by the magnetic texture. The dichroism exhibits two distinct regimes separated by a crossover: at low frequencies the absorption axis is locked to the crystal axes, while at high frequencies it tracks the helix. Anisotropic dc transport provides a complementary quantitative signature, and all signatures are robust to weak intrinsic spin-orbit coupling. Our results identify polarization-resolved optics and sublattice-sensitive spectroscopy as direct probes of textured altermagnetic states and suggest a route to direction-selective electronic and optical functionality in altermagnets.

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Altermagnets are antiferromagnetically ordered materials with vanishing net magnetization where opposite-spin sublattices are related by a crystal symmetry other than a simple real-space translation or inversion [1–9]. This multipolar order breaks time-reversal and crystal-rotation symmetries separately while preserving their combinations, yielding spin-split bands at finite momentum despite zero net magnetization. Experimental signatures have been reported in materials including RuO_2 , $\alpha\text{-MnTe}$, and Mn_5Si_3 , among others [10–17]. In particular, layered d -wave altermagnets in the $\text{V}_2\text{Se}_2\text{O}$ family, including $\text{KV}_2\text{Se}_2\text{O}$ and $\text{Rb}_{1-\delta}\text{V}_2\text{Te}_2\text{O}$, provide quasi-two-dimensional platforms with spin-resolved evidence of room-temperature spin-momentum locking [18,19].

Domain formation and nonuniform Néel configurations are common in antiferromagnets due to the absence of stray fields [20]. Consistent with this expectation, nanoscale probes already indicate nonuniform Néel textures in candidate altermagnets [21–23]. Yet, the impact of spin textures on the electronic response in altermagnets remains largely unexplored, with only a few recent works addressing texture dynamics and associated transport phenomena [24–29].

While collinear antiferromagnets and altermagnets both have zero net magnetization, their electronic responses to

spin textures can differ sharply. In conventional antiferromagnets, the Hamiltonian is invariant under Néel reversal ($\mathbf{n} \rightarrow -\mathbf{n}$), implying that texture-induced responses are even in $\mathbf{n}$. In altermagnets, by contrast, the electronic structure can distinguish the sign of $\mathbf{n}$, making certain observables potentially sensitive to texture chirality and enabling emergent gauge-field signatures [24,25,30]. These effects are analogous in spirit to texture-induced Berry-phase phenomena in ferromagnets [31].

In this Letter, we develop an effective low-energy theory for itinerant electrons in spin-textured collinear altermagnets using the $\text{SU}(2)$ gauge formalism [32–36] and show that gradients of the Néel order act as emergent gauge fields on the electronic pseudospin. This framework has three generic consequences: it generates a texture-induced pseudospin-orbit coupling, yields an emergent electromagnetic coupling in the presence of texture singularities, and produces a texture-controlled pseudospin splitting. We illustrate these effects for representative d -wave and g -wave altermagnets and, for a coplanar spin helix, show that the helix wave vector controls both the magnitude and the principal axes of the dc conductivity and the optical absorption. The altermagnetic channel is isolated by the momentum-resolved sublattice-odd spectral structure and the interband dichroism, both of which vanish when $K^z = 0$ and neither of which requires intrinsic spin-orbit coupling. We identify two response regimes: at finite but low frequencies the dominant axis is locked to the crystal axes, whereas at high frequencies it enters a tracking regime where it follows the helix orientation. Unlike spin-dependent transport [37] and valley-selective dichroism [38] in *homogeneous* altermagnets, the texture-controlled changes identified here vanish in the uniform crystal.

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Model—A minimal description starts from collinear antiferromagnetic order complemented by hopping processes that do not respect the translation that maps one magnetic sublattice onto the other [39]. In this setting, the antiferromagnetic order enforces the antialignment of the local moments on the two sublattices. In this setting, a spin texture is described by a slowly varying Néel vector $\mathbf{n}(\mathbf{r})$, which defines the local spin quantization axis. We do not model the microscopic stabilization of the magnetic texture. We instead take a smooth Néel texture as a prescribed background and compute the electronic response to it.

The itinerant electronic Hamiltonian in the absence of intrinsic spin-orbit coupling (SOC) is

$$H = H_{\text{kin}}(\mathbf{k})\sigma_0 - J\mathbf{n}(\mathbf{r}) \cdot \boldsymbol{\sigma}\eta_z, \quad (1)$$

where $\mathbf{k} = (k_x, k_y)$ is the crystal momentum in a two-sublattice (A, B) unit cell, $H_{\text{kin}}(\mathbf{k})$ encodes spin-independent hopping processes, η_α (σ_i) are Pauli matrices acting in sublattice (spin) space, and J is the exchange coupling.

We consider a two-dimensional altermagnet and use a continuum parametrization around an inversion-symmetric point,

$$H_{\text{kin}}(\mathbf{k}) = \frac{k^2}{2m} + \left[C^x + \frac{K^x}{2} k^2 \right] \eta_x + \frac{K^z}{n!} g_n(\mathbf{k}) \eta_z, \quad (2)$$

where m sets the isotropic dispersion, C^x and K^x the sublattice hybridization, and K^z is the symmetry-allowed staggered anisotropy associated with the basis function $g_n(\mathbf{k})$, with $g_2(\mathbf{k}) = k_x k_y$ for the d -wave case and $g_4(\mathbf{k}) = k_x k_y (k_x^2 - k_y^2)$ for the g -wave case.

We treat an inhomogeneous Néel texture $\mathbf{n}(\mathbf{r})$ using a local unitary $U(\mathbf{r}) \in \text{SU}(2)$ that aligns the spin quantization axis with the exchange field [32–36] such that $U^\dagger(\mathbf{r})\mathbf{n}(\mathbf{r}) \cdot \boldsymbol{\sigma}U(\mathbf{r}) = \sigma_z$. Spatial gradients generate a connection

$$D_j \equiv \partial_j + U^\dagger \partial_j U \equiv \partial_j - \frac{i}{2} A_j \cdot \boldsymbol{\sigma}, \quad (3)$$

so in the comoving frame the exchange field becomes uniform, while the kinetic energy couples to the texture through $\partial_j \rightarrow D_j$. To handle operator ordering, we define the gauge-coupled kinetic term by Weyl symmetrization as

$$H = \mathcal{W}[H_{\text{kin}}(k_j \rightarrow -iD_j)] - J\sigma_z \eta_z, \quad (4)$$

where $\mathcal{W}[\cdot]$ denotes full symmetrization over noncommuting covariant derivatives.

The unitary $U(\mathbf{r})$ defines a comoving orthonormal triad $(\mathbf{e}_1, \mathbf{e}_2, \mathbf{n})$ and allows one to decompose the connection into components transverse and longitudinal to $\mathbf{n}$,

$$A_j = A_j^\perp + A_j^\parallel \mathbf{n}. \quad (5)$$

The transverse component is uniquely fixed by the texture, $A_j^\perp = \mathbf{n} \times \partial_j \mathbf{n}$, whereas the longitudinal part $A_j^\parallel = -\mathbf{e}_1 \cdot \partial_j \mathbf{e}_2$ is the residual $\text{U}(1)_z$ gauge field (rotations of $\mathbf{e}_{1,2}$ about $\mathbf{n}$). Observables depend only on gauge-invariant combinations such as the texture metric,

$$g_{jk} \equiv A_j^\perp \cdot A_k^\perp = \partial_j \mathbf{n} \cdot \partial_k \mathbf{n}, \quad (6)$$

as well as linear coupling in the transverse connection of the form $A_j^\perp \partial_k$ and, for textures with topological defects, the skyrmion density.

We focus on the d -wave case ($n = 2$), while the general band expansion can be found in Supplemental Material (SM) [40]. Including the texture through covariant derivatives, the resulting Hamiltonian becomes

$$H_d = C^x \eta_x - \frac{1}{2m} \eta_0 D^2 - \frac{K^x}{2} \eta_x D^2 - \frac{K^z}{2} \frac{\{D_x, D_y\}}{2} \eta_z - J\sigma_z \eta_z. \quad (7)$$

The Weyl symmetrized product reduces to an anticommutator and naturally splits into three components:

$$\begin{aligned} \frac{1}{2} \{D_j, D_k\} &= \left(\partial_j - iA_j^\parallel \frac{\sigma_z}{2} \right) \left(\partial_k - iA_k^\parallel \frac{\sigma_z}{2} \right) \\ &\quad - i(A_j^\perp \partial_k + A_k^\perp \partial_j) \cdot \frac{\boldsymbol{\sigma}^\perp}{2} - \frac{g_{jk}}{4} \sigma_0. \end{aligned} \quad (8)$$

The first term represents the *emergent electromagnetic coupling* to the Abelian gauge field $A_j^\parallel$; the second term represents the *emergent pseudospin-orbit coupling* generated by the transverse texture; and the third term is a *scalar potential* proportional to the texture metric.

We work in the exchange-dominated regime, where the doublet separation $2E_0$ exceeds the spin- and sublattice-dependent kinetic scales at the occupied Fermi surface and treat those terms perturbatively. Let $P = \frac{1}{2}(1 + \sigma_z \eta_z)$ denote the projector onto the low-energy subspace spanned by $\{|A\uparrow\rangle, |B\downarrow\rangle\}$. By projecting into the low-energy sector, the resulting effective Hamiltonian reads as

$$\begin{aligned} H_{d,\text{eff}} &= PH_{\text{kin}}P \\ &= \frac{1}{2m} \left[\sum_{\alpha=x,y} \left(-iD_\alpha^\parallel \right)^2 + \frac{g_{\alpha\alpha}}{4} \tau_0 \right] \\ &\quad - \frac{K^z}{2} \tau_3 \left[D_x^\parallel D_y^\parallel - \frac{g_{xy}}{4} \tau_0 \right] \\ &\quad - iK^x (A_x^\perp \partial_x + A_y^\perp \partial_y) \cdot \frac{\boldsymbol{\tau}}{2}, \end{aligned} \quad (9)$$

where $D_\alpha^\parallel = \partial_\alpha - iA_\alpha^\parallel (\tau_3/2)$ and τ_i are the Pauli operators inside the two-component subspace. A nonzero C^x is

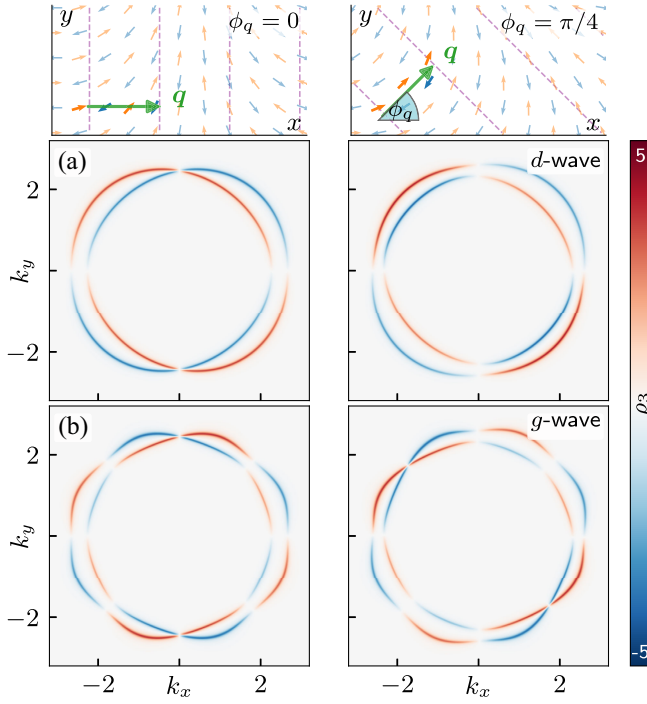


FIG. 1. Real-space spin helices and their sublattice-resolved spectral imprint. Top: schematic planar antiferromagnetic helices with propagation vector $\mathbf{q} = q \cos(\phi_q)\hat{x} + q \sin(\phi_q)\hat{y}$ for $\phi_q = 0$ (left) and $\pi/4$ (right), measured from the crystalline $\hat{x}$ axis. Dashed lines are constant-phase lines perpendicular to $\mathbf{q}$. Rows (a) and (b) show the sublattice-odd spectral weight $\rho_3(\mathbf{k})$ on the Fermi contours in the (k_x, k_y) plane for a d -wave and g -wave altermagnet, respectively, at the same two helix orientations. The helix splits the Fermi surface along $\mathbf{q}$ and reorients the momentum-resolved sublattice-odd weight with ϕ_q .

treated nonperturbatively and renormalizes the coefficients [40]. Concretely, the expansion then requires only the low-energy hierarchy $K^{x,z}k_F^2, v_F q \ll 2E_0$ and a Fermi energy below the gap $2E_0 = 2\sqrt{J^2 + (C^x)^2}$, with k_F and v_F being the Fermi momentum and velocity, respectively.

Equation (9) yields a compact low-energy theory for electrons coupled to an arbitrary smooth Néel texture in a collinear altermagnet. For a conventional collinear antiferromagnet ($K^z = 0$), the sublattice-odd channel is absent and the effective theory contains no τ_3 term, so texture-induced responses remain sublattice even. For an altermagnet ($K^z \neq 0$), the texture couples to the sublattice-odd sector and produces a local τ_3 component set by metric contractions of texture gradients, i.e., an emergent pseudospin splitting. Consequently, textures such as domain walls and spirals can imprint a localized τ_3 signal in observables like the sublattice-resolved spectral function.

Spin helix—A coplanar spin helix, which captures the local rank-one constant-gradient limit of an arbitrary smooth texture, is one of the simplest and most physically relevant spin textures. Its pitch and orientation may be set by Dzyaloshinskii–Moriya interactions, strain gradients,

boundary conditions, or other micromagnetic mechanisms that break inversion symmetry [42], which are not included explicitly in Eqs. (1) and (2). Consider a planar helix with propagation wave vector $\mathbf{q} = (q_x, q_y) = q(\cos \phi_q, \sin \phi_q)$, whose orientation angle ϕ_q is measured from the crystalline $\hat{x}$ axis, parametrized as

$$\mathbf{n}(\mathbf{r}) = \mathbf{u} \cos(\mathbf{q} \cdot \mathbf{r}) + \mathbf{v} \sin(\mathbf{q} \cdot \mathbf{r}), \quad (10)$$

where $\mathbf{u}$ and $\mathbf{v}$ are fixed orthonormal unit vectors spanning the helix plane and $\mathbf{w} \equiv \mathbf{u} \times \mathbf{v}$ is the normal to that plane, such that $\mathbf{n} \times \partial_i \mathbf{n} = q_i \mathbf{w}$ [top panels of Fig. 1]. Choosing the comoving frame with $\mathbf{e}_1 = \mathbf{w}$ fixes $A_j^\parallel = 0$ and $A_j^\perp = q_j \mathbf{e}_1$.

In the d -wave case, the long-wavelength Hamiltonian for itinerant electrons reads as

$$H_{d,h} = \frac{1}{2m} \left[k_x^2 + k_y^2 + \frac{q_x^2 + q_y^2}{4} \right] + \frac{K^x \tau_1}{2} [q_x k_x + q_y k_y] + \frac{K^z \tau_3}{2} \left[k_x k_y + \frac{q_x q_y}{4} \right]. \quad (11)$$

The analogous g -wave Hamiltonian, obtained at quartic order in gradients, has the same τ_1 texture coupling and a τ_3 channel that mixes lower harmonics; its explicit form is given in SM [40].

Within the effective theory, a smooth helix generates two texture-induced effects. First, the emergent pseudospin-orbit coupling induces a Fermi surface splitting. Second, it modifies the momentum-resolved sublattice-odd spectral weight $\rho_3(\mathbf{k}, \omega) \equiv (2\pi)^{-1} \text{Im Tr}[\tau_3(G^A - G^R)]$. Its momentum integral defines the net polarization $\mathcal{P}_3(\omega) \equiv \int_{\mathbf{k}} \rho_3(\mathbf{k}, \omega)$. Here $G^{R,A}(\mathbf{k}, \omega) = [\omega + \mu - H(\mathbf{k}) \pm i\Gamma]^{-1}$ are Green's functions, μ is the chemical potential, and Γ is an effective broadening. The key distinction between the d -wave and g -wave models is the momentum structure of this τ_3 contribution: in the g -wave case, the contraction of the intrinsic $\ell = 4$ anisotropy with the texture metric lowers the effective symmetry, making the induced τ_3 polarization alternate in sign around the contour. For helix directions along crystal axes, the continuum model admits symmetry-enforced degeneracy lines that intersect a given Fermi contour at isolated points; away from these special orientations, the helix generically lifts the degeneracy and splits the Fermi surfaces anisotropically [Figs. 1(a) and 1(b)].

The texture-induced τ_3 term acts as an emergent pseudospin Zeeman field controlled jointly by the altermagnetic form factor K^z and the texture metric; it therefore vanishes in either the conventional-antiferromagnetic ($K^z = 0$) or uniform limit. The net polarization $\mathcal{P}_3$ vanishes as $q \rightarrow 0$, although the momentum-resolved $\rho_3(\mathbf{k}, \omega)$ need not. This momentum-resolved structure can be probed by atomically registered Fourier-transform scanning tunneling

spectroscopy: subtracting the quasiparticle-interference maps of the two magnetic sublattices isolates the odd scattering channel [43,44].

For a homogeneous Néel state ($q = 0$), tetragonal symmetry enforces $\sigma_{xx} = \sigma_{yy}$. A spin helix lowers the symmetry and unlocks an anisotropic correction. It is therefore natural to resolve the dc conductivity in the spiral basis $(\hat{q}, \hat{q}_\perp)$, with $\sigma_\parallel = \hat{q} \cdot \boldsymbol{\sigma}^{\text{dc}} \cdot \hat{q}$, $\sigma_\perp = \hat{q}_\perp \cdot \boldsymbol{\sigma}^{\text{dc}} \cdot \hat{q}_\perp$, and $\sigma_\times = \hat{q} \cdot \boldsymbol{\sigma}^{\text{dc}} \cdot \hat{q}_\perp$. The splitting $\sigma_\parallel - \sigma_\perp$ measures the helix-induced anisotropy, while $\sigma_\times$ captures any residual misalignment between the helix and the conductivity principal axes. Altermagnetic helices exhibit $\sigma_\parallel \neq \sigma_\perp$, with a pronounced tetragonal modulation in the d -wave case and a weaker, $\pi/4$ -periodic modulation in the g -wave case [40]. In both cases, $\sigma_\times$ remains subleading, indicating that $\hat{q}$ closely tracks the principal axes of the symmetric conductivity tensor.

This anisotropy is genuine but not, by itself, unique to altermagnetism: a conventional antiferromagnetic helix exhibits it as well. In the matched isotropic conventional-AFM model, $\sigma_\parallel - \sigma_\perp$ is independent of the helix orientation and $\sigma_\times = 0$, so the tensor rigidly follows the helix. The altermagnet is quantitatively distinguished by an anisotropy growing as $\mathcal{O}(q)$ rather than $\mathcal{O}(q^2)$ [40,45]. Reproducing its crystal-locked modulation and finite $\sigma_\times$ in a conventional AFM requires an additional crystal-locked C_4 anisotropy [40]. The dc angular structure is therefore a quantitative rather than absolute discriminator. The qualitative altermagnetic fingerprint instead lies in the interband optical response, to which we now turn.

Turning to optical absorption, we find that the system exhibits *linear dichroism*: the dissipated power depends on the in-plane polarization direction, and two orthogonal linear polarizations are absorbed differently at the same frequency. In linear response this is encoded in the dissipative optical conductivity, $\text{Re}\sigma_{ij}(\omega)$, whose two eigenvalues $\sigma_+(\omega) \neq \sigma_-(\omega)$ correspond to absorption along orthogonal in-plane principal axes. We quantify the dichroism strength by

$$\mathfrak{D}(\omega) \equiv \frac{\sigma_+(\omega) - \sigma_-(\omega)}{\sigma_+(\omega) + \sigma_-(\omega)} \quad (12)$$

and introduce the polarization-averaged response $\bar{\sigma}(\omega) \equiv [\sigma_+(\omega) + \sigma_-(\omega)]/2$, which sets the absorption for unpolarized light. The principal absorption axis $\theta_+(\omega)$ is defined as the eigenvector angle associated with $\sigma_+(\omega)$.

The optical conductivity is obtained from the standard relation $\sigma_{ij}(\omega) = (i\omega)^{-1}[K_{ij}^R(\omega) - K_{ij}^R(0)]$, where $K_{ij}^R(\omega)$ is the retarded current-current correlator. To analyze anisotropy and symmetry, we focus on the interband paramagnetic contribution,

$$K_{ij}^{\text{para}}(\omega) = -e^2 \int d\mathbf{k} \sum_{n \neq m} \frac{f_{nk} - f_{mk}}{\omega + i\Gamma - \Delta_{mn}} v_i^{nm} v_j^{mn}, \quad (13)$$

where $v_i^{nm}(\mathbf{k})$ are the interband velocity matrix elements encoding the full angular dependence, $\Delta_{mn} = E_{mk} - E_{nk}$, and $f_{nk} = [e^{E_{nk}/T} + 1]^{-1}$ are Fermi-Dirac occupation factors at temperature T .

Since the helix breaks the crystalline C_4 symmetry down to C_2 , the tensor structure of $K_{ij}^{\text{para}}(\omega)$ becomes angle dependent. Specifically, the interband absorption is controlled by the off-diagonal velocity matrix elements that, for the d -wave helix with $q \ll k_F$, give

$$v_i^{+-} \approx -\frac{K^x K^z}{4} \frac{q_i k_i^2}{\sqrt{d_1^2 + d_3^2}}, \quad (14)$$

where $(\bar{x} \equiv y, \bar{y} \equiv x)$; see SM [40]. Equation (14) shows that the optical anisotropy is not trivially locked to $\mathbf{q}$ since x -polarized transitions are controlled by q_y (and vice versa), reflecting interference between the helix-induced mixing through pseudospin-orbit coupling and the crystal form factor.

In conventional collinear antiferromagnets, linear dichroism is usually attributed to intrinsic spin-orbit or crystal-field anisotropies [46,47]. In their absence it is symmetry forbidden: for translation-connected sublattices and a coplanar texture, a π spin rotation about $\mathbf{w}$ composed with the sublattice-connecting translation commutes with

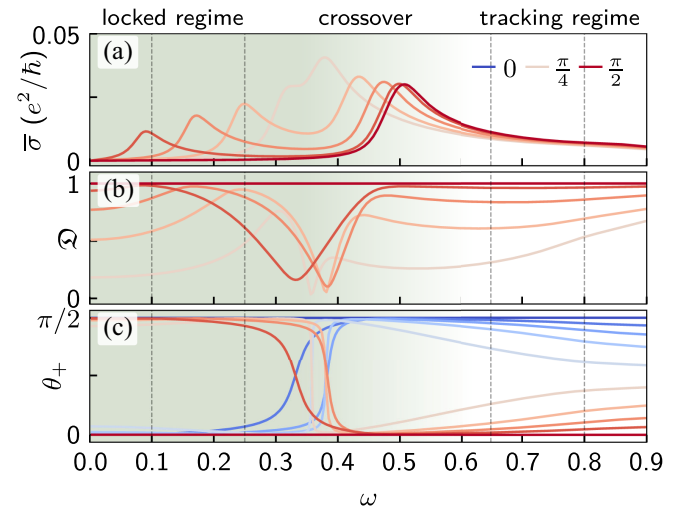


FIG. 2. (a) Total optical absorption $\bar{\sigma}(\omega)$ for ten helix orientations ϕ_q from 0 (blue) to $\pi/2$ (red), color coded as in the legend ($0, \pi/4, \pi/2$ marked). Reflection about the $k_x = k_y$ axis maps $\phi_q \rightarrow \pi/2 - \phi_q$ and leaves $\bar{\sigma}$ and $\mathfrak{D}$ invariant, so the curves at ϕ_q and $\pi/2 - \phi_q$ coincide in (a) and (b); in particular the $\phi_q = 0$ (blue) and $\pi/2$ (red) lines overlap. (b) Corresponding linear dichroism $\mathfrak{D}(\omega)$. (c) Frequency-dependent orientation $\theta_+(\omega)$ of the principal absorption axis. The vertical dashed lines mark the four probe frequencies used in Fig. 3(a).

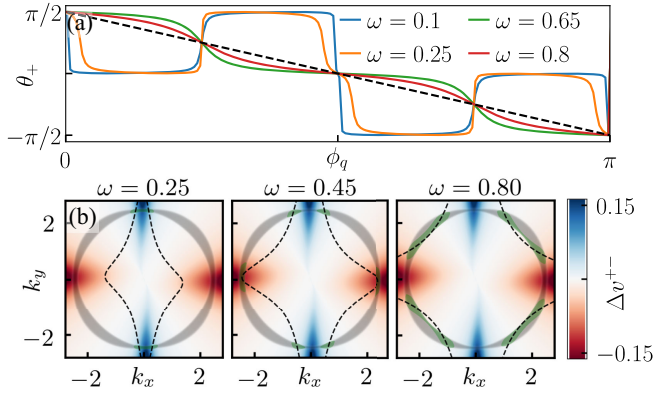


FIG. 3. (a) Principal absorption-axis angle $\theta_+(\omega)$ versus helix orientation ϕ_q . In the *locked* regime θ_+ is pinned to the crystal axes (plateaus with $\pi/2$ jumps), while in the *tracking* regime it varies smoothly and approaches $\theta_+ \simeq \pi/2 - \phi_q$ (dashed). (b) Interband-velocity anisotropy $\Delta v^{+-} \equiv |v_x^{+-}|^2 - |v_y^{+-}|^2$ for the three indicated probe frequencies at fixed helix orientation $\phi_q = \pi/8$. Black dashed: resonance $E_+(\mathbf{k}) - E_-(\mathbf{k}) = \omega$. Thick gray: Fermi surface. Green shading: absorption weight $W_\omega(\mathbf{k})$ (Pauli factor times resonance kernel), concentrated where the resonance contour crosses the Pauli-allowed region near the Fermi surface.

H and with the current, and the two bands are its ± 1 eigenstates, so $v_i^{+-} \equiv 0$ [11,40]. By contrast, in an altermagnet, a distinct *texture-induced interband mechanism* is already present without intrinsic spin-orbit coupling since the textured altermagnet features a nonrelativistic spin-split structure with momentum-dependent eigenvectors that feed directly into interband optical matrix elements. It is also distinct from the valley-selective dichroism of *homogeneous* altermagnets [38], a spin-conserving selection rule that yields no net dichroism unless the fourfold symmetry is broken externally; here the texture supplies both the symmetry breaking and the interband coupling, giving a net dichroism whose axes track the helix.

The overall line shape of $\bar{\sigma}(\omega)$ is controlled by the interband resonance condition, while the relative weight of its features depends on the helix orientation ϕ_q [Figs. 2(a) and 2(b)]. This angular dependence reflects that the relevant interband matrix elements are governed by the texture-induced mixing and its interference with the altermagnetic form factor. Two regimes separated by a crossover frequency $\omega_c(\phi_q) \equiv \text{argmin}_\omega \mathfrak{D}(\omega)$ can be identified. For $\omega \lesssim \omega_c$, the response is strongly polarization selective and $\mathfrak{D}(\omega) \approx 1$ over a broad window, indicating absorption dominated by a single linear polarization [Fig. 2(b)]. In this *locked regime*, the principal axis θ_+ develops broad pinning plateaus and undergoes sharp $\pi/2$ reorientations as ϕ_q is varied [Fig. 3(a)]. Near ω_c , $\mathfrak{D}(\omega)$ shows a pronounced dip accompanied by an abrupt $\pi/2$ rotation of θ_+ , signaling a narrow *crossover window* where the two eigenvalues of $\text{Re}\sigma$ become nearly degenerate. For $\omega \gtrsim \omega_c$ the response crosses over to a *tracking regime* in which $\theta_+(\omega)$ varies smoothly and approaches a helix-tracking

law,

$$\theta_+ \simeq \frac{\pi}{2} - \phi_q. \quad (15)$$

Equation (15) implies that, in the tracking regime, the absorptive eigenbasis is not the spiral frame $(\hat{q}, \hat{q}_\perp)$ but approaches the swapped direction $\hat{e}_{\text{tr}} = (\sin \phi_q, \cos \phi_q)$, i.e., the helix direction reflected by $x \leftrightarrow y$.

This behavior can be understood by expressing the absorptive tensor in terms of interband matrix elements. Schematically, the interband contribution has the form

$$\text{Re}\sigma_{ij}(\omega) \propto \int d\mathbf{k} v_i^{-+}(\mathbf{k}) v_j^{+-}(\mathbf{k}) W_\omega(\mathbf{k}), \quad (16)$$

where $W_\omega(\mathbf{k})$ is a peaked function centered on hot spots placed along the resonance condition $\omega = E_+(\mathbf{k}) - E_-(\mathbf{k})$, with a width set by Γ and T . The interference between the twofold texture gap and the fourfold altermagnetic gap produces two inequivalent minima in the Fermi-surface splitting. Deviations for $\omega \lesssim \omega_c$ arise when $W_\omega(\mathbf{k})$ is dominated by anisotropic hot spots along the resonant contour, so the resonance average is no longer approximately $x \leftrightarrow y$ symmetric. In particular, for $\omega \ll \omega_c$ only the lower valley is accessible, and the weight concentrates near the crystal axes ($k_x \simeq 0$ or $k_y \simeq 0$), causing the switching regime. When the second valley becomes accessible, additional resonant states with a competing polarization enter and drive the dichroism dip, placing ω_c near this threshold. The resulting scaling form is $\omega_c \simeq \frac{1}{2} K^z k_F^2 X(\rho, \phi_q)$, with $\rho = 2K^x q / (K^z k_F)$. For finite weak textures ($\rho \ll 1$), $\omega_c \sim K^x q k_F$ up to an orientation factor [40]. For band parameters representative of layered *d*-wave altermagnets and texture scales already imaged in altermagnets, the locked-to-tracking dichroism spans the terahertz-to-mid-infrared window [40].

At high ω , the resonant weight samples the contour almost uniformly, so the resonance average becomes approximately $x \leftrightarrow y$ symmetric [Fig. 3(b)]. In this regime, using Eq. (14), we can rewrite Eq. (16) as

$$\text{Re}\sigma_{ij}(\omega) \propto q_i q_j \left\langle \frac{k_i^2 k_j^2}{d_1^2 + d_3^2} \right\rangle_{w_\omega}, \quad (17)$$

and $x \leftrightarrow y$ symmetric sampling implies the factorization $\text{Re}\sigma \propto (q_y, q_x)(q_y, q_x)^T$, resulting in the tracking law in Eq. (15).

Discussion—Néel textures provide a direct way to distinguish altermagnets from conventional collinear antiferromagnets in electronic response. The texture induces three generic ingredients: an emergent U(1) gauge coupling tied to topological defects, a metric-controlled pseudospin splitting, and a transverse texture-induced pseudospin-orbit coupling.

For a coplanar spin helix, the altermagnetic contribution is isolated by two fingerprints that vanish when $K^z = 0$: the

momentum-resolved sublattice-odd spectral structure and the SOC-free interband linear dichroism, a more readily accessible optical probe of the same channel. Interband absorption becomes polarization selective and undergoes a characteristic frequency-dependent reorientation, producing linear dichroism that tracks the helix orientation. Together, sublattice-sensitive spectroscopy and polarization-resolved optics provide direct probes of textured altermagnetic states and a practical discriminator from collinear antiferromagnets. As shown in the End Matter, these signatures survive weak intrinsic spin-orbit coupling, which does not spoil the SOC-free mechanism but instead renders the helix-plane normal, and hence the texture chirality, optically observable.

Additionally, the dc conductivity becomes anisotropic, but this is a quantitative rather than qualitative discriminator as related antiferromagnetic helices can also exhibit anisotropic transport.

The framework is independent of microscopic details and extends naturally to other altermagnetic symmetries via the appropriate form factors. It can incorporate higher-gradient corrections and time-dependent textures, and it provides a route to texture-engineered phenomena, including defect-pinned bound states [34,35] and spatially programmable anisotropies relevant for spintronics [48]. Combined with superconducting proximity [49,50], it enables texture-controlled Andreev spectra [51–56] as well as magnetoelectric [57] and optoelectronic [58–60] responses.

Note added—Recently, we have become aware of an independent preprint developing a closely related low-energy description of electrons in altermagnetic textures [61]. The results are complementary to this work.

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Data availability—The data that support the findings of this article are openly available [62].

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End Matter

The Hamiltonian (1) omits intrinsic (relativistic) spin-orbit coupling. Since a spin helix is frequently stabilized by the Dzyaloshinskii–Moriya interaction, which itself originates in SOC, it is important to establish that weak intrinsic SOC does not spoil the texture-generated response and to identify how it modifies it. A full account is given in SM; here we summarize the result.

To linear order in $\mathbf{k}$, and once inversion is broken, invariance under time reversal, the altermagnetic antiunitary $[TC_{4z}\eta_x]$, and the mirror-plus-sublattice-swap $M_x\eta_x$ leave three SOC channels [40,63],

$$H_{\text{SOC}} = \lambda_A \eta_0 (\mathbf{k} \times \boldsymbol{\sigma})_z + \lambda_B \eta_x (\mathbf{k} \times \boldsymbol{\sigma})_z + \lambda_C \eta_y [\mathbf{k} \times (\mathbf{n} \times \boldsymbol{\sigma})]_z, \quad (\text{A1})$$

two Néel-even Rashba terms ($\lambda_{A,B}$) and one Néel-odd, texture-locked term ($\lambda_C \propto J$). Projecting onto the low-energy doublet intrinsic SOC does not open any new operator at first order: it enters *only* as an additive shift of the same emergent gauge fields generated by the texture. For $C^x = 0$, Channels B and C fuse into a single transverse shift $\propto \bar{\alpha} = \lambda_B + \lambda_C$, while the uniform Rashba λ_A shifts the Abelian field. Explicitly, projecting Eq. (A1) onto the doublet shifts the two emergent connections of the texture,

$$\delta \mathbf{A}_j^\perp = -\frac{2\bar{\alpha}}{K^x} (\hat{\mathbf{z}} \times \hat{\mathbf{x}}_j)_{\perp \mathbf{n}}, \quad \delta A_j^\parallel = -2m\lambda_A (\mathbf{n} \times \hat{\mathbf{z}})_j, \quad (\text{A2})$$

where $\{\hat{\mathbf{x}}, \hat{\mathbf{y}}, \hat{\mathbf{z}}\}$ are the lab-frame axes, $(\cdot)_{\perp \mathbf{n}}$ denotes the component transverse to $\mathbf{n}$, and the Abelian shift carries

the emergent magnetic field $\delta b = 2m\lambda_A \nabla \cdot \mathbf{n}_{\parallel}$ [40]. Unlike the texture connections, whose direction is set by the texture and whose magnitude is $\propto \partial \mathbf{n}$, the SOC shift is crystal pinned and depends on $\mathbf{n}$ itself.

For a helix with plane normal $\mathbf{w} = \mathbf{u} \times \mathbf{v}$, averaging Eq. (A2) over the helix phase yields a rigid renormalization of the helix wave vector in the τ_1 channel,

$$\mathbf{q}_{\text{eff}} = \mathbf{q} + \underbrace{\frac{2\bar{\alpha}}{K^x}(\hat{\mathbf{z}} \times \mathbf{w})}_{\delta \mathbf{q}}, \quad \zeta \equiv \frac{2\bar{\alpha}}{K^x q} |\mathbf{w} \times \hat{\mathbf{z}}|, \quad (\text{A3})$$

with ζ being the dimensionless ratio of the SOC shift to the texture gauge field. The effective theory is thus *form-invariant*: the entire SOC-free machinery applies with $\mathbf{q} \rightarrow \mathbf{q}_{\text{eff}}$, and the locked-to-tracking crossover survives.

Intrinsic SOC therefore makes the helix-plane normal $\mathbf{w}$ observable. In the SOC-free theory this orientation can be changed by a global spin rotation, but Eq. (A3) couples $\mathbf{w}$ to the crystal axis $\hat{\mathbf{z}}$ through $\delta \mathbf{q}$. This gives three simple limits. For an in-plane spiral ($\mathbf{w} \parallel \hat{\mathbf{z}}$), $\delta \mathbf{q} = 0$, so the first-order SOC correction vanishes. For a proper screw with chirality $s = \pm 1$ ($\mathbf{w} = s\hat{\mathbf{q}}$), $\delta \mathbf{q} \perp \mathbf{q}$ and $\mathbf{q}_{\text{eff}}/q = \hat{\mathbf{q}} + s\zeta(\hat{\mathbf{z}} \times \hat{\mathbf{q}})$, so its azimuth is $\phi_{\text{eff}} = \phi_q + \arctan(s\zeta)$. Since the tracking law becomes $\theta_+ \simeq \pi/2 - \phi_{\text{eff}}$, the absorption axis shifts by the chirality-odd angle $\delta\theta_+ = -\arctan(s\zeta) = \mp \arctan \zeta$. For a cycloid ($\mathbf{w} \perp \hat{\mathbf{q}}$), $\delta \mathbf{q} \parallel \mathbf{q}$ and $\mathbf{q}_{\text{eff}} = (1 \mp \zeta)\mathbf{q}$, so $\omega_c \rightarrow (1 \mp \zeta)\omega_c$ [Fig. 4(b)], as derived in SM [40]. Since $\delta \mathbf{q}$ changes sign under helix-chirality reversal $\mathbf{w} \rightarrow -\mathbf{w}$, intrinsic SOC makes the optical response sensitive to the texture handedness selected by the stabilizing interaction, a sensitivity absent in the SOC-free theory.

These first-order conclusions require $\zeta \ll 1$; the projection additionally requires $|\lambda_i|k_F \ll 2E_0$ for $i = A, B, C$. An exact four-band calculation retaining all winding terms confirms

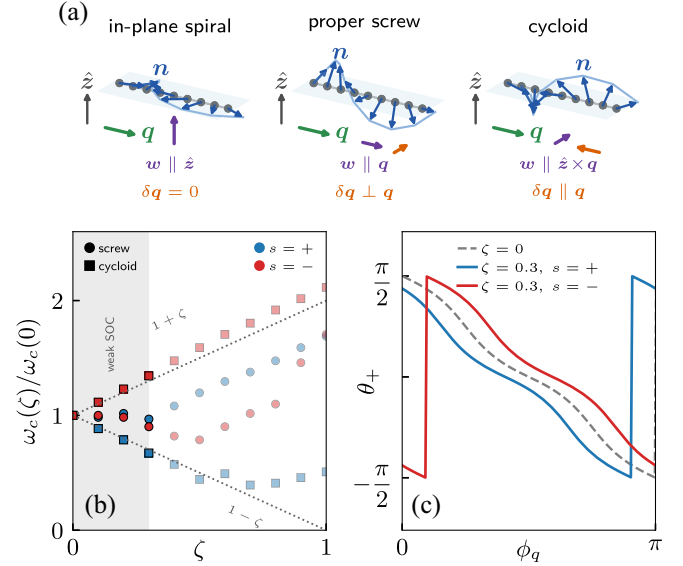


FIG. 4. Effect of intrinsic SOC on the optical crossover, parametrized by ζ [Eq. (A3)]; blue and red denote the two helix chiralities $s = \pm$ ($\mathbf{w} \rightarrow -\mathbf{w}$). (a) Representative in-plane spiral, proper screw, and cycloid in a quasi-two-dimensional interface. Blue arrows show the Néel vector $\mathbf{n}$; green, purple, and orange arrows show $\mathbf{q}$, $\mathbf{w}$, and $\delta \mathbf{q}$, respectively. (b) Numerically exact crossover frequency $\omega_c(\zeta)/\omega_c(0)$, identified by the dichroism minimum that carries the locked-to-tracking principal-axis switch, from a superlattice Bloch-ladder calculation of the four-band helix model with all winding terms retained, for the screw (circles) and cycloid (squares). Dotted lines are the cycloid guides $1 \pm \zeta$. (c) Principal absorption axis $\theta_+(\phi_q)$ in the tracking regime from the projected model, for the screw at $\zeta = 0$ (dashed) and $\zeta = 0.3$ (solid); SOC shifts the whole tracking law rigidly by the chirality-odd angle $\mp \arctan \zeta$.

the locked-to-tracking switch and its chirality-dependent displacement in this regime [Fig. 4(b); see SM [40]].