

Dispersion-mediated modal reconfiguration in a binary photonic lattice

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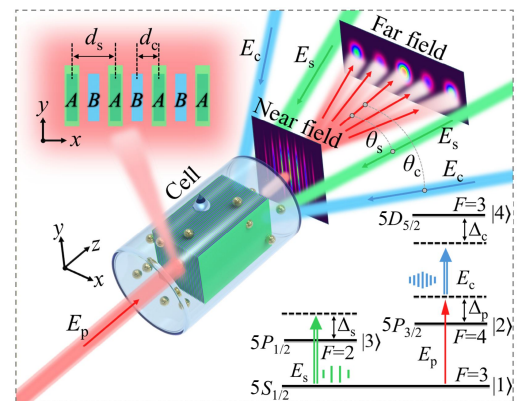
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ABSTRACT

Dynamic control of on-site potentials is key to controlling band structures and light transport in lattice systems. Here, we demonstrate a dynamically tunable binary photonic lattice in a four-level atomic system, formed by superimposing two standing-wave fields with a 1:2 period ratio, under far-detuned electromagnetically induced transparency. In this regime, absorption is strongly suppressed, and the refractive-index response is dominated by its dispersive component. By tuning the laser frequency detuning, we continuously reshape the refractive-index contrast between two sublattices, enabling all-optical control of the on-site potentials. This control drives a clear reconfiguration of the optical modes, whereby the probe field shifts its concentration in real space, accompanied by an alternating distribution of reciprocal-space energy between the first- and second-order diffraction channels. Our results demonstrate a low-loss route to on-demand mode reconfiguration, opening new possibilities for programmable light steering and reconfigurable photonic functions in atomic media.

Keywords binary photonic lattice, electromagnetically induced transparency, all-optical control, low-loss transmission



1 Introduction

Photonic lattices, enabled by spatial refractive-index modulations, provide a versatile platform for exploring light propagation [1], topological properties [2], and non-Hermitian physics [3–5]. Among the parameters that govern their behavior, the local optical potentials at different lattice sites are particularly important, as they set the relative index depths, dictate localization and diffraction, and allow photonic systems to emulate

condensed-matter Hamiltonians [6–8]. Within this context, binary photonic lattices, consisting of two periodically alternating sublattices, represent the simplest model in which on-site potentials can be independently controlled [9]. Their two-site unit-cell structure maps naturally onto Dirac-type Hamiltonians, enabling photonic analogues of relativistic phenomena such as Zitterbewegung [10, 11], Bloch–Zener oscillations [12, 13], Dirac solitons [14], Klein tunneling [15], and electron-positron pair creation and annihilation [16].



Dynamically controlling the relative potentials of lattice sites is key to realizing programmable photonic materials. Most existing lattices rely on fixed structural designs that cannot be adjusted once fabricated [1, 17]. These constraints motivate the search for alternative mechanisms that enable *in situ* and continuous tuning of lattice-site potentials without relying on rigid photonic architectures. Atomic vapors provide a compelling solution, as their refractive index can be continuously controlled through coherent light-matter interactions in real time [18–26]. This makes atom-based photonic lattices a natural platform for reconfigurable optical control [27–30] and for exploring quantum photonic devices [31–33]. A particularly powerful strategy is to use two independent periodic optical fields within the vapor that create spatial regions with distinct electromagnetically induced transparency (EIT) conditions, enabling local modification of lattice-site potentials or selective addressing of individual sublattices [34–37]. Previous demonstrations of non-Hermitian photonic graphene used this strategy to introduce loss difference between sublattices, leading to energy imbalance and localization at the less lossy sites [35]. More recently, an electromagnetically induced binary photonic lattice with a tunable flat band was realized in a three-level atomic system through Rabi-frequency control of two standing-wave fields [37]. An open question is whether similar control over intra-lattice energy redistribution can be realized through purely dispersive, low-loss modulation.

In this work, we exploit this opportunity by constructing a reconfigurable binary photonic lattice in a four-level ^{85}Rb atomic vapor under far-detuned EIT condition. The refractive-index contrast between the two sublattices can be continuously tuned and even inverted by adjusting the frequency detuning of a single laser field. This enables nearly loss-free reshaping of both real-space concentration and reciprocal-space diffraction distributions. These results reveal a near-lossless, dispersion-driven route to controlling binary photonic lattices, establishing atomic vapors as a versatile platform for programmable and reconfigurable photonic band engineering.

2 Experimental setup

Figure 1(a) shows the experimental configuration for constructing a tunable binary photonic lattice in an ^{85}Rb vapor. The vapor cell has a length of 10 cm and is maintained at 110°C , corresponding to a rubidium atomic density of approximately $1.12 \times 10^{13} \text{ cm}^{-3}$. The probe beam waist is approximately $880 \mu\text{m}$, defined as the $1/e^2$ intensity radius of the Gaussian beam. Three external-cavity diode lasers provide the probe (E_p , wavelength $\lambda_p \approx 780 \text{ nm}$), coupling (E_c , $\lambda_c \approx 776 \text{ nm}$), and signal (E_s , $\lambda_s \approx 795 \text{ nm}$) fields. These fields drive the transitions

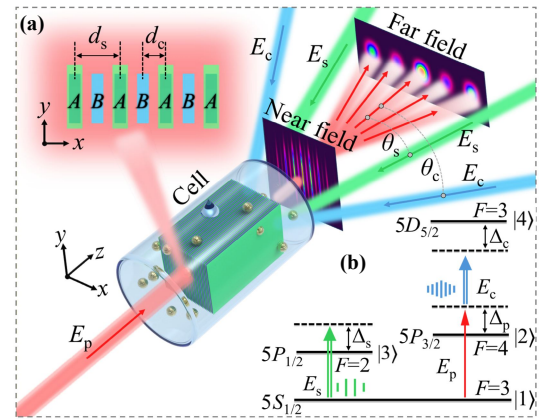


Fig. 1 Construction of a binary photonic lattice. (a) Experimental sketch. (b) Four-level V+ Ξ type configuration in ^{85}Rb atomic vapors.

of a V+ Ξ -type atomic system, and their corresponding detunings are denoted as Δ_p , Δ_c , and Δ_s , as illustrated in Fig. 1(b). To sculpt the spatial modulation, the coupling beam E_c is split into two paths that intersect inside the vapor cell at a small angle $2\theta_c \approx 0.4^\circ$, forming a standing-wave field with period $d_c \approx 110 \mu\text{m}$ ($d = \lambda/(2 \sin \theta)$). The signal beam E_s is treated similarly but with a smaller intersection angle of $2\theta_s \approx 0.2^\circ$, yielding a second standing-wave field with twice the period, $d_s \approx 220 \mu\text{m}$. The two angles shown in Fig. 1(a), are intentionally exaggerated for visual clarity, so that the spatial distribution of the beams can be displayed more clearly. By precisely aligning their transverse positions, the two periodic fields overlap at their antinodes, forming a binary coupling field inside the atomic vapor cell. The probe beam E_p counter-propagates with the binary coupling field inside the vapor cell, forming an effective one-dimensional binary photonic lattice (i.e., a binary grating) under EIT conditions [38]. As illustrated in the inset of Fig. 1(a), sublattice A is defined by the positions where the antinodes of E_s overlap with those of E_c , while sublattice B is defined by the positions where only the antinodes of E_c are present, resulting in different transmission responses of E_p . At the same time, the near-field and far-field images of E_p are recorded simultaneously in real time to monitor its transmission behavior. For completeness, additional details regarding the experimental setup are provided in Appendix A.

Figure 2 shows the four-level EIT spectra for different detuning configurations of the probe and coupling fields, obtained by scanning the frequency of the signal field. This measurement serves as a frequency reference and is used to identify the effective low-absorption EIT operating window for the subsequent binary-lattice experiments. It is separate from the near-field and far-field measurements and is not directly related to the A/B sublattice geometry. As the probe detuning increases, the relative absorption is reduced, leading to a higher peak transmission

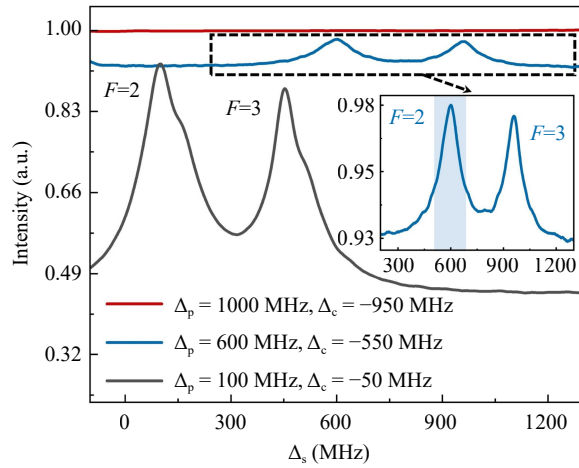


Fig. 2 V+ Ξ type four-level EIT spectra versus Δ_s for different Δ_p and Δ_c . The inset shows an enlarged view of the spectrum at $\Delta_p = 600$ MHz and $\Delta_c = -550$ MHz, where the blue-shaded region indicates the effective frequency window used in the experiment. The vertical axis denotes the transmitted intensity of the probe field, with 1.0 corresponding to full transmission (no absorption).

of the EIT signal. To ensure clear and well-resolved experimental observations in the subsequent measurements while keeping optical attenuation minimal, the probe and coupling detunings are set to $\Delta_p = 600$ MHz and $\Delta_c = -550$ MHz (blue curve). Within the effective EIT window (blue-shaded region in the inset, approximately 200 MHz wide), the probe transmission loss remains below about 6%, placing the system in a low-loss operating regime that is significantly reduced compared with near-resonant laser excitation.

3 Theory

Having established the binary photonic lattice, we next analyze how this spatial modulation reshapes the probe transmission behavior. The propagation of the probe field inside this V+ Ξ type four-level ^{85}Rb atomic system is governed by the spatially varying susceptibility, which can be written as [39]

$$\chi = \frac{iN_a|\mu_{21}|^2(\gamma_{32}\tilde{\Omega}_s^2 + \gamma_{41}\tilde{\Omega}_c^2 + 4\gamma_{32}\gamma_{41}\gamma_{43})}{\varepsilon_0\hbar G}, \quad (1)$$

where N_a is the atomic density, ε_0 is the permittivity of free space, μ_{21} is the transition dipole momentum between corresponding levels, and $G = \tilde{\Omega}_s^2\tilde{\Omega}_c^2 - \tilde{\Omega}_s^4 - \tilde{\Omega}_c^4 - 2\tilde{\Omega}_s^2(\gamma_{21}\gamma_{32} + \gamma_{41}\gamma_{43}) - 3\tilde{\Omega}_c^2(\gamma_{21}\gamma_{41} + \gamma_{32}\gamma_{43}) - 8\gamma_{21}\gamma_{32}\gamma_{41}\gamma_{43}$. Here, Ω_c and Ω_s are the Rabi frequency of corresponding fields, $\tilde{\Omega}_c = \Omega_c \sin(\pi x/d_c + \pi/2)$ and $\tilde{\Omega}_s = \Omega_s \sin(\pi x/d_s - \pi/2)$ are the effective Rabi frequency of two standing wave fields. The complex terms are defined as $\gamma_{21} = -\Gamma_2/2 + i\Delta_p$, $\gamma_{32} = -(\Gamma_3 + \Gamma_2)/2 + i(\Delta_s - \Delta_p)$, $\gamma_{41} =$

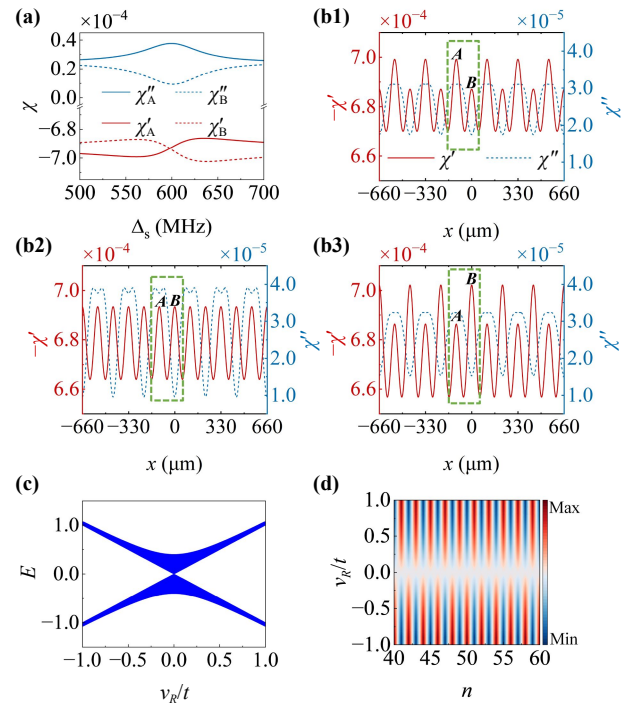


Fig. 3 Susceptibility modulation and AAH-model description of the binary photonic lattice. **(a)** Real (χ') and imaginary (χ'') parts of the susceptibility in sublattices A and B as functions of Δ_s . Here, $\Delta_p = 600$ MHz, $\Delta_c = -550$ MHz, $\Omega_c = 15$ MHz, and $\Omega_s = 5$ MHz. **(b1–b3)** Representative spatial profiles of $\chi'(x)$ and $\chi''(x)$ at three different values of $\Delta_s = 560, 600$, and 630 MHz, respectively. The dashed box indicates one unit cell of the binary lattice, consisting of sublattices A and B . **(c)** Energy spectrum of the model as a function of the modulation depth ν_R/t . **(d)** Site-resolved density of states (DOS) as a function of ν_R/t and the site index n .

$-\Gamma_4/2 + i(\Delta_p + \Delta_c)$, and $\gamma_{43} = -(\Gamma_4 + \Gamma_3)/2 + i(\Delta_c + \Delta_p - \Delta_s)$, where Γ_2 , Γ_3 , and Γ_4 are the decay rates from the upper levels.

Figure 3(a) shows the calculated real (χ') and imaginary (χ'') parts of the susceptibility in sublattices A and B as functions of Δ_s . Figures 3(b1)–(b3) show representative spatial susceptibility profiles at $\Delta_s = 560, 600$, and 630 MHz, respectively, from which the A - and B -sublattice values plotted in Fig. 3(a) are extracted at the fixed positions $x = -110$ and $x = 0$ μm , respectively. It is noted that although the signal field directly covers only the A sublattice, the coupling between adjacent sites allows variations in the signal-beam parameters to also influence the optical response of the B sublattice. For χ' , the two sublattices exhibit completely inverted behaviors on both sides of $\Delta_s = 600$ MHz, namely, $|\chi'_A| > |\chi'_B|$ for $\Delta_s < 600$ MHz, and $|\chi'_A| < |\chi'_B|$ for $\Delta_s > 600$ MHz. Although the χ'' remains slightly larger in sublattice A than in B , its magnitude is an order smaller than that of χ' under such far-detuned EIT conditions. As a result, the system effectively operates in

a Hermitian-like regime [40], where the optical response is dominated by the real part of the susceptibility. More specifically, the corresponding refractive index can be expressed as $n_{ref} = \sqrt{1 + \chi} \approx 1 + \chi/2$. Thus, the real part χ' gives rise to the refractive-index (phase) modulation, while the imaginary part χ'' describes the loss (amplitude) modulation [41]. The probe beam therefore evolves in a continuous complex optical lattice directly determined by the spatial susceptibility profile.

To capture the essential binary-lattice physics in a compact form, we further introduce an effective two-site periodic Aubry–André–Harper (AAH) model [42], whose Hamiltonian takes the form

$$H = t \sum_{n=0}^{N-2} (a_n^\dagger a_{n+1} + H.c.) + \sum_{n=0}^{N-1} \nu_R \sin(\pi n + \pi/2) a_n^\dagger a_n, \quad (2)$$

where $a_n^\dagger$ and a_n are creation and annihilation operators in site n , t is not an independently imposed parameter, but an effective coupling coefficient determined by the lattice itself, and $H.c.$ is the Hermitian conjugate. In our system, ν_R is associated with the refractive-index contrast between the two sublattices, i.e., $\nu_R \propto \Delta n_R \approx (|\chi'_A| - |\chi'_B|)/2$, which, in the present parameter regime, is equivalently characterized by the difference between the extracted A - and B -site values of χ' . Since both χ'_A and χ'_B remain negative over the considered detuning range, their absolute values are used to compare the corresponding modulation strengths. In this sense, the continuous susceptibility-based model provides the direct description of probe propagation and diffraction, while the effective AAH model serves as a physically intuitive minimal model for interpreting the detuning-controlled sublattice contrast and the associated mode reconfiguration. Changing the sign of ν_R reverses the relative on-site potentials of the two sublattices, effectively swapping the higher- and lower-index sites.

Figure 3(c) shows the calculated energy spectrum for a chain model with N sites under open boundary conditions as a function of the normalized modulation depth ν_R/t . When $\nu_R = 0$, the dispersion reduces to that of a uniform tight-binding lattice. As $|\nu_R|$ increases, the periodic on-site modulation lifts the sublattice degeneracy, causing the original band to split into two symmetric branches with an increasing gap—consistent with the enhanced contrast between the A and B sites. To further characterize the spatial distribution of the modes, we define the site-resolved density of states (DOS), or equivalently, the spatially resolved spectral density, as $D(E, n) = \sum_\alpha \delta(E - E_\alpha) |\psi_\alpha(n)|^2$, where E_α and $\psi_\alpha(n)$ are the eigenvalues and eigenvectors of the effective Hamiltonian, respectively. This quantity describes the spectral weight at energy E carried by site n , and the corresponding result is shown in Fig. 3(d). The alternating two-band pattern reflects the intrinsic two-site periodicity

of the binary lattice. As $|\nu_R|$ increases, the sublattice contrast becomes more pronounced, and reversing the sign of ν_R shifts the regions where the modes are more concentrated between the two sublattices, illustrating how tuning ν_R (experimentally via Δ_s) reshapes the modal distribution. In the present work, Figs. 3(c) and (d) are intended to illustrate the qualitative evolution of the effective model, and the relevant control parameter is therefore the normalized modulation depth ν_R/t . Accordingly, t is taken as the normalization unit in these calculations. Together, these results establish the AAH model as an effective description of the dispersion-driven binary photonic lattice and elucidate how refractive-index inversion controls the observed light-transport dynamics.

4 Results and discussion

Figure 4 shows the near-field redistribution of probe energy between the two sublattices. The probe beam is intentionally chosen to be broad, so that the near-field pattern reflects the sublattice-dependent redistribution induced by the binary lattice [35, 36, 43]. As shown in Fig. 4(a), both the simulated and measured near-field intensity patterns exhibit a clear transition as Δ_s is tuned across the inversion point. Here, the simulated results are obtained from the continuous susceptibility-based propagation model, rather than directly from the effective tight-binding model. The CCD-extracted horizontal profiles ($y = 0$) in Fig. 4(b) provide a clearer view of how the intensity shifts during this transition. In particular, the probe field becomes predominantly concentrated on the A sites at $\Delta_s = 560$ MHz and on the B sites at $\Delta_s = 640$ MHz. The fitted Gaussian profile of the absorption-free input beam (gray shaded curves) indicating that the energy concentrated on the lattice sites remains nearly loss-free in both energy-concentrated configurations.

Moreover, the total normalized intensity $I_{AB} = I_A + I_B$ shown in Fig. 4(c), varies smoothly with Δ_s , indicating that the total energy stored in the lattice changes only weakly during the energy transition. The energy fluctuations (about 10%) arise primarily from slight differences in the Gaussian weights of the A and B sampling windows, caused by the finite Gaussian profile of the input probe beam, which gives the two sampling windows slightly different background intensities, as indicated by the gray shaded curve in Fig. 4(b). Such behavior highlights the intrinsically low-loss character of the dispersion-engineered lattice. The intensity ratios I_A/I_{AB} and I_B/I_{AB} in Fig. 4(d) display clear anticorrelated trends, demonstrating a reversible redistribution of the probe intensity between the two sublattices. This process enables controllable redistribution of optical energy between the two sublattices. Accordingly, the

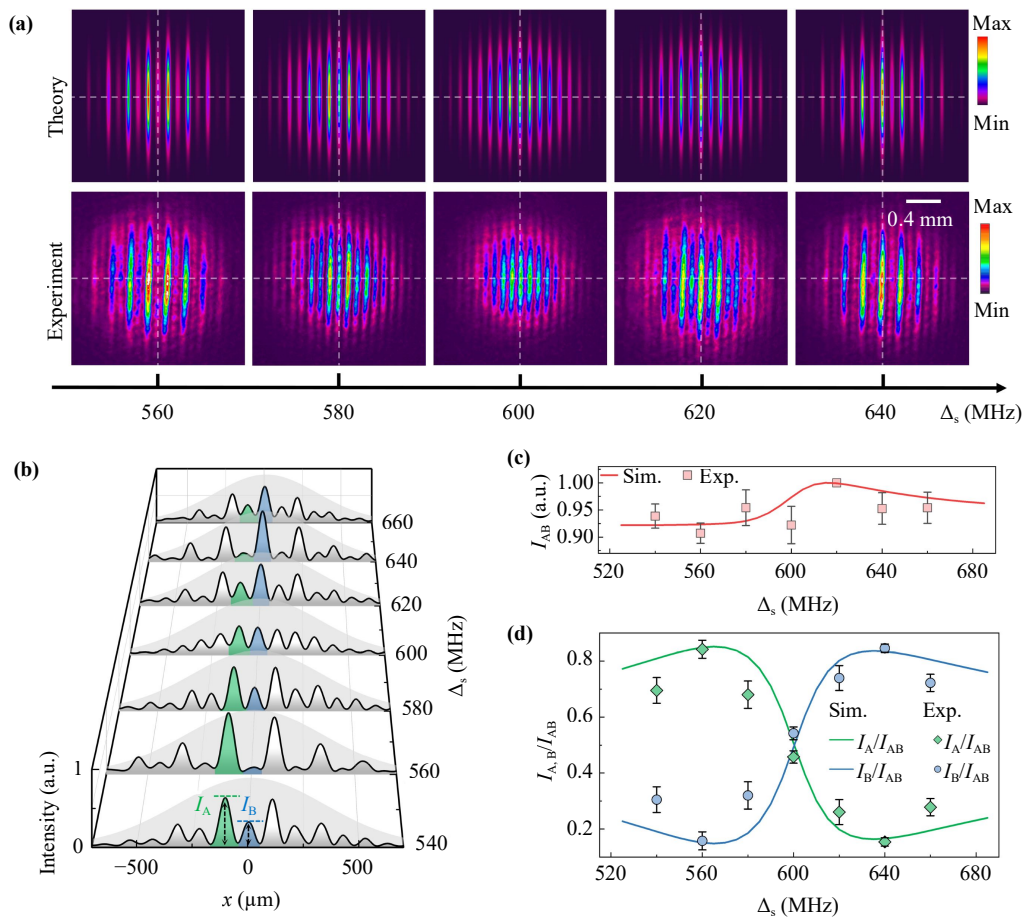


Fig. 4 Near-field energy redistribution in the binary photonic lattice. **(a)** Simulated and experimental near-field intensity distributions of the probe beam for different Δ_s . **(b)** Line profiles of the near-field intensity along the ($y = 0$) axis. The green and blue shaded curves denote the output intensities at sublattices A and B, respectively, and I_A and I_B represent their corresponding peak intensities. The gray shaded curves represent the fitted Gaussian profile of the absorption-free input beam ($\Delta_p = 1000$ MHz). **(c)** Total normalized intensity $I_{AB} = I_A + I_B$ as a function of Δ_s . **(d)** The detuning Δ_s induces energy exchanges between sublattices A and B during the propagation of probe.

near-field pattern reveals the direct real-space signature of the detuning-controlled reconfiguration of the binary lattice and provides an effective visualization of its refractive-index modulation.

Unlike the near-field diffraction that reflects the real-space distribution of the lattice, the far-field diffraction arises from grating-induced beam splitting and reveals how the probe's modal energy is redistributed among different momentum components [43]. In atom-based lattice systems, far-field diffraction has become one of the key directions of study, since enhancing higher-order diffraction is of direct importance for applications such as beam splitting, optical routing, and controllable diffraction-channel engineering [33, 44–47]. Figure 5 shows how the real-space sublattice switching manifests in reciprocal space through the reorganization of diffraction orders. Figure 5(a) shows the simulated and experimentally measured far-field diffraction patterns, where the simulations are calculated from the output field of the continuous susceptibility-based model and its corresponding far-field

diffraction. To make the first- and second-order diffraction features more clearly visible, the images are intentionally overexposed. Under these conditions, the diffraction pattern undergoes a reversible reorganization: energy initially concentrated in the ± 2 nd orders (or even 0th order) flows toward the ± 1 st orders as Δ_s approaches the inversion point, and then returns to the ± 2 nd orders as the detuning increases further. The absence of visible 1st-order diffraction at larger detunings is caused by Doppler broadening, residual atomic absorption, the finite CCD detection threshold, and other factors not fully included in the present model, rather than by a new physical effect. This reversible migration of modal weight directly reflects the detuning-controlled modulation of the real-part contrast in the binary lattice.

Figure 5(b) compares the simulated and experimentally measured normalized diffraction intensities as functions of Δ_s . The relatively large error bars at small and large detunings arise because the diffraction intensities are more sensitive to detuning fluctuations in these regions,

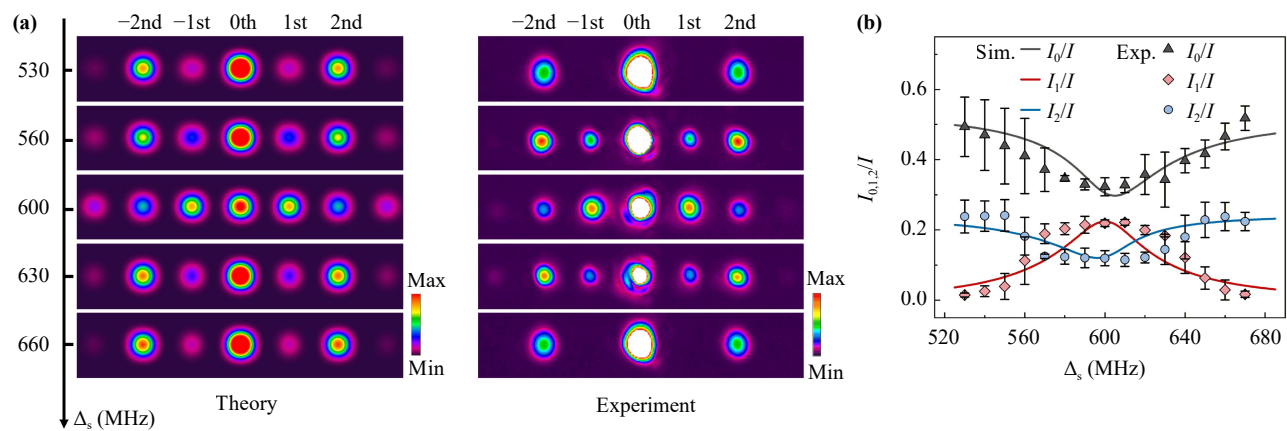


Fig. 5 Far-field energy redistribution in the binary photonic lattice. **(a)** Simulated and experimentally observed far-field diffraction patterns of the probe beam as a function of Δ_s . **(b)** Normalized diffraction intensities I_j/I as functions of Δ_s , where I_j ($j = 0, 1, 2$) denote the intensities of the 0th, 1st, and 2nd diffraction orders, and $I = I_0 + 2I_1 + 2I_2$ represents the total transmitted intensity (neglecting higher-order contributions).

whereas this sensitivity is reduced near the balance point around 600 MHz and the corresponding uncertainties become smaller. In addition, since the present theory is based on a simplified model that does not include effects such as Doppler broadening and finite laser linewidth, some data points near 600 MHz may deviate from the theoretical curves. Nevertheless, the overall experimental trend still follows the theoretical prediction reasonably well, revealing two successive energy-exchange processes. Notably, under the far-detuned conditions used in this study, the normalized diffraction intensities of both orders exceed 20% of the total transmitted diffraction power, and the system can even be tuned into a regime where the 2nd-order diffraction becomes stronger than the 1st-order. In contrast to most existing related studies, which mainly focus on enhancing the 1st-order diffraction efficiency [45, 48–51], this behavior highlights the binary-lattice nature of the system and provides an additional degree of freedom for controlling the dominant diffraction channel, which is beneficial for flexible beam splitting and tunable light routing in momentum space. As a proof-of-principle demonstration, the present efficiencies may not yet have reached the highest values and could be further improved by optimizing the detuning, interaction length, beam overlap, and standing-wave visibility. This reconfigurable efficiency of a higher-order channel constitutes a practical performance advantage of the dispersion-engineered binary lattice and provides enhanced flexibility for tunable light-field routing in momentum space.

5 Conclusion

In summary, we have realized a reconfigurable binary photonic lattice in a four-level ^{85}Rb atomic vapor and uncovered its fundamental light-transport behavior

under tunable refractive-index contrast between adjacent sublattices. By tuning the signal-field detuning, we continuously invert the effective on-site potentials experienced by the probe field, thereby inducing a clear modal reconfiguration. Experimentally, this behavior manifests as a low-loss energy transfer between the *A* and *B* sublattices in real space, together with a reciprocal-space energy reorganization characterized by alternating enhancement of the first- and second-order diffraction channels. Notably, the system can even be tuned into a regime in which the 2nd-order diffraction becomes stronger than the 1st-order diffraction. In contrast to most related studies, which mainly focus on enhancing the 1st-order diffraction efficiency, this feature provides an additional degree of freedom for controlling the dominant diffraction channel and highlights the binary-lattice nature of the system. These results, corroborated by the two-site periodic AAH model, reveal that the system behaves as a dispersion-dominated lattice in which the real part of the susceptibility dominates the light-transport dynamics, whereas the imaginary part becomes negligible under far-detuned EIT conditions. As a proof-of-principle demonstration, the present results already show the potential of this platform for flexible beam splitting, optical routing, and momentum-space light-field control. This work provides a simple yet effective route to low-loss, momentum-selective light routing and coherent energy redistribution in photonic lattices, enabling reconfigurable photonic band engineering and optical information processing.

Declarations The authors declare that they have no competing interests and there are no conflicts.

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Appendix A: Experimental details

Figure A1 schematically illustrates the experimental configuration. Three external-cavity diode lasers (DL pro, Toptica) provide the probe, coupling, and signal beams, respectively. The three beams are first divided into two orthogonally polarized components by a polarizing beam splitter (PBS). The vertically polarized components are then completely overlapped inside the atomic vapor cell to form a standard EIT probing channel, in which E_c and E_s co-propagate, while E_p counter-propagates with respect to them. A photodetector (PD, PDA36A, Thorlabs) is used to record the transmitted intensity of E_p when scanning the frequency of E_s for frequency calibration and monitoring, as shown in Fig. 2.

The horizontally polarized component of E_c and E_s are expanded by a beam expander (BE) and subsequently split into two beams via a non-polarizing beam splitter (BS). The two split beams of E_c and E_s then interfere at different angles by a BS, forming two standing-wave fields with a period ratio of 1:2. These two standing-wave fields are superimposed at their antinodes, thereby creating a binary optical field inside the atomic vapor cell. The horizontally polarized component of E_p then counter-propagates with the binary optical field inside the vapor cell. The powers of the probe, coupling, and signal beams at the entrance of the vapor cell are about 1 mW, 45 mW, and 15 mW, respectively. Within the cell, it experiences a spatially modulated refractive index and, together with E_c and E_s under EIT conditions, forms a proposed one-dimensional binary photonic lattice, as illustrated in Fig. 1.

The near-field diffraction pattern of the transmitted E_p is imaged onto a charge-coupled device (CCD, SP928, Ophir-Spiricon) using an imaging lens with a focal length of $f = 100$ mm, with both the lens-to-CCD and lens-to-cell distances set to twice the focal length. The far-field distribution is simultaneously recorded by another CCD through a Fourier lens with the same focal length, where both distances are set to one focal length. Two CCDs are thus employed for real-time, synchronous monitoring of the E_p evolution in both real and reciprocal spaces.

Appendix B: Model and equations of the $V+\Xi$ type energy-level system

In this work, we consider a four-level $V+\Xi$ atomic

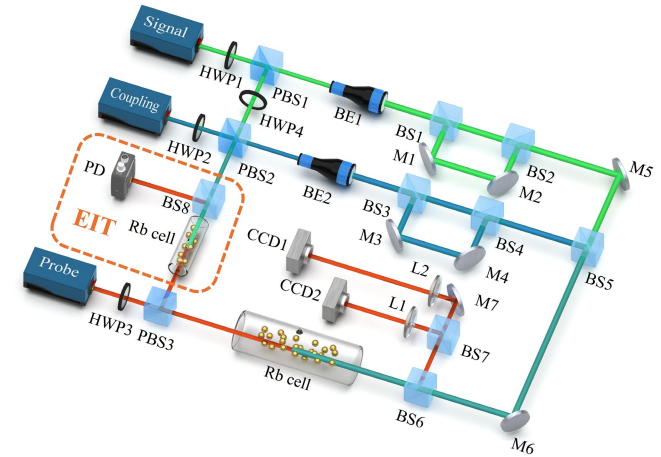


Fig. A1 Experimental setup. HWP: half-wave plate; BE: beam expander; PBS: polarization beam splitter; BS: non-polarizing beam splitter; L: lens; M: mirror; CCD: charge-coupled device (SP928, Ophir-Spiricon); PD: photodetector (PDA36A, Thorlabs).

system in ^{85}Rb medium, as illustrated in Fig. 1(b). The optical response of the probe field is controlled by two standing-wave fields with different transverse periods, generated independently by the coupling and signal beams. Since each standing wave is formed by two symmetric beams derived from the same laser frequency, its spatial envelope is stationary. The corresponding standing-wave fields can be written as

$$\begin{aligned} E_c^{\text{SW}}(x, z, t) &\propto 2E_c \cos(k_c x \sin \theta_c) \exp[i(k_c z \cos \theta_c - \omega_c t)], \\ E_s^{\text{SW}}(x, z, t) &\propto 2E_s \cos(k_s x \sin \theta_s) \exp[i(k_s z \cos \theta_s - \omega_s t)], \end{aligned} \quad (\text{B1})$$

where θ_c and θ_s are the half crossing angles of the two beam pairs. The z -dependent factors correspond only to common propagation phases, while the spatial envelopes depend only on the transverse coordinate x . Therefore, the resulting binary superlattice is stationary and does not acquire an additional modulation along the z direction.

Under the rotating-wave approximation, the Hamiltonian of the system is given by

$$\begin{aligned} H = & \hbar[\Delta_p|2\rangle\langle 2| + \Delta_s|3\rangle\langle 3| + (\Delta_p + \Delta_c)|4\rangle\langle 4|] \\ & + \frac{\hbar}{2}[\Omega_p|1\rangle\langle 2| + \tilde{\Omega}_s|1\rangle\langle 3| + \tilde{\Omega}_c|2\rangle\langle 4|] + H.c., \end{aligned} \quad (\text{B2})$$

where $\Omega_p = \mu_{12}E_p/\hbar$, $\Omega_c = \mu_{24}E_c/\hbar$, $\Omega_s = \mu_{13}E_s/\hbar$ are the Rabi frequencies of the probe, coupling, and signal fields, respectively, and $\Delta_p = \omega_p - \omega_{21}$, $\Delta_c = \omega_c - \omega_{42}$, $\Delta_s = \omega_s - \omega_{31}$ are the corresponding detunings. Here, ω_p , ω_c , and ω_s stand for the frequencies of probe, coupling and signal fields, respectively, and μ_{ij} , ω_{ij} ($i, j = 1, 2, 3, 4$) are respectively the transition dipole moment and frequency differences between levels $|i\rangle$ and level $|j\rangle$. The

notation $H.c.$ denotes the Hermitian conjugate.

To describe the transverse standing-wave modulation, we introduce the effective Rabi frequencies

$$\begin{aligned}\tilde{\Omega}_c &= \Omega_c \sin\left(\frac{\pi x}{d_c} + \frac{\pi}{2}\right), \\ \tilde{\Omega}_s &= \Omega_s \sin\left(\frac{\pi x}{d_s} - \frac{\pi}{2}\right),\end{aligned}\quad (B3)$$

where d_c and d_s are the transverse periods of the two standing waves. We note that these effective Rabi frequencies are allowed to change sign because they represent the local coherent field amplitudes of the

standing waves, rather than optical intensities. In other words, the sign change reflects the local phase reversal of the standing wave, whereas the optical intensity remains nonnegative.

To characterize the optical response of the system, we employ the density-matrix formalism. The time evolution of the density matrix is governed by the Liouville equation

$$\frac{\partial \rho}{\partial t} = -\frac{i}{\hbar}[H, \rho] + \hat{L}(\rho), \quad (B4)$$

where the Lindblad form of Liouvillian matrix $\hat{L}(\rho)$ describing the relaxation by spontaneous decay is expressed as

$$\hat{L}(\rho) = \begin{pmatrix} L_{11} & -\frac{\Gamma_2}{2}\rho_{12} & -\frac{\Gamma_3}{2}\rho_{13} & -\frac{\Gamma_4}{2}\rho_{14} \\ -\frac{\Gamma_2}{2}\rho_{21} & L_{22} & -\frac{\Gamma_2 + \Gamma_3}{2}\rho_{23} & -\frac{\Gamma_2 + \Gamma_4}{2}\rho_{24} \\ -\frac{\Gamma_3}{2}\rho_{31} & -\frac{\Gamma_3 + \Gamma_2}{2}\rho_{32} & -\Gamma_3\rho_{33} & -\frac{\Gamma_3 + \Gamma_4}{2}\rho_{34} \\ -\frac{\Gamma_4}{2}\rho_{41} & -\frac{\Gamma_4 + \Gamma_2}{2}\rho_{42} & -\frac{\Gamma_4 + \Gamma_3}{2}\rho_{43} & -\Gamma_4\rho_{44} \end{pmatrix}, \quad (B5)$$

where $L_{11} = \Gamma_2\rho_{22} + \Gamma_3\rho_{33}$, $L_{22} = -\Gamma_2\rho_{22} + \Gamma_4\rho_{44}$, and Γ_i ($i = 1, 2, 3, 4$) are the decay rate from the corresponding levels. The corresponding decay rates are $\Gamma_1 = 0$ (since level $|1\rangle$ is ground state), $\Gamma_2 = 2\pi \times 6.07$ MHz, $\Gamma_3 = 2\pi \times 5.75$ MHz, and $\Gamma_4 = 2\pi \times 0.67$ MHz, respectively. The equations of motion for the density-matrix elements are given by

$$\begin{aligned}\dot{\rho}_{11} &= \Gamma_2\rho_{22} + \Gamma_3\rho_{33} + \frac{i}{2}(\rho_{21} - \rho_{12})\Omega_p + \frac{i}{2}(\rho_{31} - \rho_{13})\tilde{\Omega}_s, \\ \dot{\rho}_{22} &= -\Gamma_2\rho_{22} + \Gamma_4\rho_{44} + \frac{i}{2}(\rho_{12} - \rho_{21})\Omega_p + \frac{i}{2}(\rho_{42} - \rho_{24})\tilde{\Omega}_c, \\ \dot{\rho}_{33} &= -\Gamma_3\rho_{33} + \frac{i}{2}(\rho_{13} - \rho_{31})\tilde{\Omega}_s, \\ \dot{\rho}_{44} &= -\Gamma_4\rho_{44} + \frac{i}{2}(\rho_{24} - \rho_{42})\tilde{\Omega}_c, \\ \dot{\rho}_{21} &= \gamma_{21}\rho_{21} - \frac{i}{2}(\rho_{22} - \rho_{11})\Omega_p + \frac{i}{2}\rho_{41}\tilde{\Omega}_c - \frac{i}{2}\rho_{23}\tilde{\Omega}_s, \\ \dot{\rho}_{31} &= \gamma_{31}\rho_{31} - \frac{i}{2}\rho_{32}\Omega_p - \frac{i}{2}(\rho_{33} - \rho_{11})\tilde{\Omega}_s, \\ \dot{\rho}_{41} &= \gamma_{41}\rho_{41} - \frac{i}{2}\rho_{42}\Omega_p + \frac{i}{2}\rho_{21}\tilde{\Omega}_c - \frac{i}{2}\rho_{43}\tilde{\Omega}_s, \\ \dot{\rho}_{32} &= \gamma_{32}\rho_{32} - \frac{i}{2}\rho_{31}\Omega_p - \frac{i}{2}\rho_{34}\tilde{\Omega}_c + \frac{i}{2}\rho_{12}\tilde{\Omega}_s, \\ \dot{\rho}_{42} &= \gamma_{42}\rho_{42} - \frac{i}{2}\rho_{41}\Omega_p - \frac{i}{2}(\rho_{44} - \rho_{22})\tilde{\Omega}_c, \\ \dot{\rho}_{43} &= \gamma_{43}\rho_{43} + \frac{i}{2}\rho_{23}\tilde{\Omega}_c - \frac{i}{2}\rho_{41}\tilde{\Omega}_s.\end{aligned}\quad (B6)$$

Here, $\gamma_{21} = -\Gamma_2/2 + i\Delta_p$, $\gamma_{31} = -\Gamma_3/2 + i\Delta_s$, $\gamma_{41} = -\Gamma_4/2 + i(\Delta_p + \Delta_c)$, $\gamma_{32} = -(\Gamma_3 + \Gamma_2)/2 + i(\Delta_s - \Delta_p)$, $\gamma_{42} = -(\Gamma_4 + \Gamma_2)/2 + i\Delta_c$, and $\gamma_{43} = -(\Gamma_4 + \Gamma_3)/2 + i(\Delta_p + \Delta_c - \Delta_s)$.

Under the EIT condition, the system is driven by a

weak probe field together with relatively strong coupling and signal fields. We therefore adopt the weak-probe approximation, $\Omega_p \ll \Omega_c, \Omega_s$, under which the atomic population remains predominantly in the ground state, i.e., $\rho_{11} \approx 1$, $\rho_{22} \approx \rho_{33} \approx \rho_{44} \approx 0$.

Meanwhile, under the steady-state condition $\dot{\rho}_{ij} = 0$, the probe coherence ρ_{21} can be solved analytically as

$$\rho_{21} = \frac{i\Omega_p \left(\gamma_{32}\tilde{\Omega}_s^2 + \gamma_{41}\tilde{\Omega}_c^2 + 4\gamma_{32}\gamma_{41}\gamma_{43} \right)}{G}, \quad (B7)$$

where $G = \tilde{\Omega}_s^2\tilde{\Omega}_c^2 - \tilde{\Omega}_s^4 - \tilde{\Omega}_c^4 - 2\tilde{\Omega}_s^2(\gamma_{21}\gamma_{32} + \gamma_{41}\gamma_{43}) - 3\tilde{\Omega}_c^2(\gamma_{21}\gamma_{41} + \gamma_{32}\gamma_{43}) - 8\gamma_{21}\gamma_{32}\gamma_{41}\gamma_{43}$.

To establish the connection between the microscopic atomic response and the macroscopic optical susceptibility, we derive the polarization from the induced atomic dipole moment. The macroscopic polarization of the atomic ensemble is

$$P = N_a\mu_{21}\rho_{21}, \quad (B8)$$

where N_a is the atomic density. On the other hand, the polarization is related to the probe electric field through

$$P = \epsilon_0\chi E_p. \quad (B9)$$

Combining Eqs. (B8) and (B9), the susceptibility experienced by the probe beam is obtained as

$$\chi = \frac{iN_a|\mu_{21}|^2 \left(\gamma_{32}\tilde{\Omega}_s^2 + \gamma_{41}\tilde{\Omega}_c^2 + 4\gamma_{32}\gamma_{41}\gamma_{43} \right)}{\epsilon_0\hbar G}. \quad (B10)$$

The refractive index is related to the susceptibility through



$$n_{ref} = \sqrt{1 + \chi} \approx 1 + \chi/2. \quad (\text{B11})$$

Writing $\chi = \chi' + i\chi''$, and $n_{ref} = n_0 + n_R + i n_I$ ($n_0 \approx 1$ is the background index of the medium), the real and imaginary components of the refractive index can be written as $n_R \approx \chi'/2$, and $n_I \approx \chi''/2$, respectively. Under the far-detuned EIT condition considered in the main text, χ' dominates over χ'' , so the optical lattice is mainly dispersion-controlled.

Under the slowly varying envelope and paraxial approximations, the propagation of the probe field through the atomic medium is governed by the following Schrödinger-like equation:

$$i \frac{\partial \psi(x, y, z)}{\partial z} = -\frac{1}{2k_p} \nabla_{\perp}^2 \psi(x, y, z) - \frac{k_p}{n_0} \Delta n(x) \psi(x, y, z), \quad (\text{B12})$$

where $\psi(x, y, z)$ denotes the slowly varying electric-field envelope of the probe field E_p , $\nabla_{\perp}^2 = \partial^2/\partial x^2 + \partial^2/\partial y^2$ is the transverse Laplacian, z is the propagation direction, $k_p = 2\pi n_0/\lambda_p$ is the wave number of the probe field in the medium, and $\Delta n(x) = n_R(x) + i n_I(x)$ represents the complex refractive-index modulation induced by the atomic medium. For the parameters used in our simulations, the real part of the susceptibility is approximately one order of magnitude larger than its imaginary part, indicating that the probe-field evolution is dominated by the dispersive refractive-index modulation. Under this dispersion-dominated condition, equation (B12) provides a continuous propagation framework for simulating the transverse evolution and diffraction of the probe field in the atomic lattice.

For the calculations presented here, the subsequent free-space evolution of the probe field into the far field must also be considered. To obtain the exit-plane field required for the subsequent analytical calculation of the far-field distribution, we adopt the thin-grating approximation. Within this approximation, the spatially varying complex susceptibility determines the local amplitude attenuation and phase accumulation acquired by the probe field over the interaction length. The corresponding propagation equation is

$$\frac{\partial E_p}{\partial z} = (-k_p \chi'' + i k_p \chi') E_p(x, y, z). \quad (\text{B13})$$

Since the susceptibility modulation is transverse and stationary, the transmission function after an interaction length L can be written as

$$T(x, y) = E_p(x, y, 0) \exp(-k_p \chi'' L) \exp(i k_p \chi' L). \quad (\text{B14})$$

Here, the first exponential term describes the amplitude modulation, while the second term describes the phase modulation induced by the atomic medium.

Under the framework of the Fresnel–Kirchhoff diffraction theory and within the paraxial approximation, the intensity profile of the probe beam at an observation plane located a distance z from the output surface can be described as

$$I_{\text{near}}(X, Y, Z) = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} T(x, y) \exp \left[i k_p \left(2Z + \frac{x^2 + y^2}{2Z} - \frac{xX + yY}{2Z} + \frac{X^2 + Y^2}{2Z} \right) \right] dx dy, \quad (\text{B15})$$

where (x, y) and (X, Y) are the coordinates in the object and observation planes, respectively. In order to solve the integral, $T(x, y)$ in equation is represented as a two-dimensional Fourier series.

$$T(x, y) = \sum_{m, h=-\infty}^{\infty} C_{mh} \exp \left[-i 2\pi \left(\frac{m}{a} x + \frac{h}{b} y \right) \right], \quad (\text{B16})$$

where C_{mh} is a 2D Fourier coefficient, while a and b are the spatial periods along the X and Y directions, respectively. By substituting equation (B16) into equation (B15), we can obtain the diffraction intensity distribution along z as

$$I_{\text{near}}(X, Y, Z) = C_0 \sum_{m, h=-\infty}^{\infty} C_{mh} \exp \left[-i \pi \lambda_p Z \left(\frac{m^2}{a^2} + \frac{h^2}{b^2} \right) \right] \times \exp \left[i 2\pi \left(\frac{m}{a} X + \frac{h}{b} Y \right) \right], \quad (\text{B17})$$

where the irrelevant terms have been absorbed into C_0 . Since the refractive index information is encoded in the two-dimensional Fourier coefficients, the discrete diffraction patterns of the probe beam can be controlled by tuning the parameters of the binary photonic lattice.

The far-field diffraction pattern can be obtained by performing the Fourier transform of the output field. Within the Fraunhofer diffraction approximation, the far-field intensity distribution for far-field diffraction is given by

$$I_{\text{far}}(X, Y) = \frac{\exp(i k_p z)}{i \lambda_p z} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} T(x, y) \times \exp \left[i \frac{k_p}{2z} (X^2 + Y^2) \right] dx dy. \quad (\text{B18})$$

These expressions are used in the main text to simulate the near-field and far-field diffraction patterns shown in Figs. 4 and 5.

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