

Storage and retrieval of optical skyrmions with topological characteristics

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Optical skyrmions are topological structures of light whose defining property, the skyrmion number, is robust against perturbations. This makes them attractive for applications in quantum information storage, where resilience to decoherence is paramount. However, their preservation during coherent storage remains unexplored. We report the first experimental demonstration of storing and retrieving optical skyrmions in a cold ⁸⁷Rb vapor using a dual-path electromagnetically induced transparency memory. Crucially, we show that the skyrmion number remains invariant for storage times up to several microseconds, even when subjected to imbalanced loss between the two paths and substantial perturbations in control beam power. Our work demonstrates the survival of a nontrivial topological invariant in a quantum memory, marking a significant step towards topologically protected photonic technologies.

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I. INTRODUCTION

Optical skyrmions, as topological quasiparticles with distinct polarization textures, have become an active area of research in recent years [1–3]. These structures promise inherent stability against perturbations, making them potential candidates for photonic devices and information processing technologies [4–6]. Originally proposed in the context of particle physics, the concept of skyrmions has since then been translated to various physical systems, such as magnetic materials [7,8], ultracold atoms [9–11], acoustic systems [12], and fluid systems [13,14]. In optics, skyrmions manifest as vectorial textures of light fields, allowing the exploration of various skyrmionic structures [15–25]. Their preparation is comparatively straightforward, since optical systems are free from the crystal structure constraints inherent to condensed matter systems. Recent studies of optical skyrmions demonstrate that their topological information could resist noise and

defects in complex systems [26–28], with potential applications in light-matter interactions [29–33], as well as classical or quantum communication and computation [34–38]. For instance, quantum information could be encoded in a two-photon entangled state defined by a fixed skyrmion number, where the topological protection suppresses noise effects and keeps the global invariant [39].

Quantum networks with topological quasiparticles require quantum memories in addition to communication channels. So far, the optical storage of skyrmionic beams remains unexplored. Storing polarization-encoded information, such as polarization qubits [40–42] and vector beams [43–46] in atomic ensembles typically requires a dual-path interferometric configuration to cancel polarization sensitivity, which arises due to the presence of multiple Zeeman sublevels, although alternative methods exist [47,48]. This approach, however, demands highly balanced efficiency and phase stability between the two paths, as any imperfection severely degrades the output fidelity [49]. In light of these challenges, optical skyrmions emerge as a promising alternative information carrier that could overcome the limitations of conventional polarization storage schemes. Their topological properties could be preserved under unitary or nonunitary transformations [50–53] which are commonly encountered in optical systems, for instance, in scattering or propagation through noisy channels, making them resistant to decoherence and mode distortion. This suggested topological protection could provide a reliable stability that potentially cancels the effective noise arising from decoherence, efficiency imbalances,

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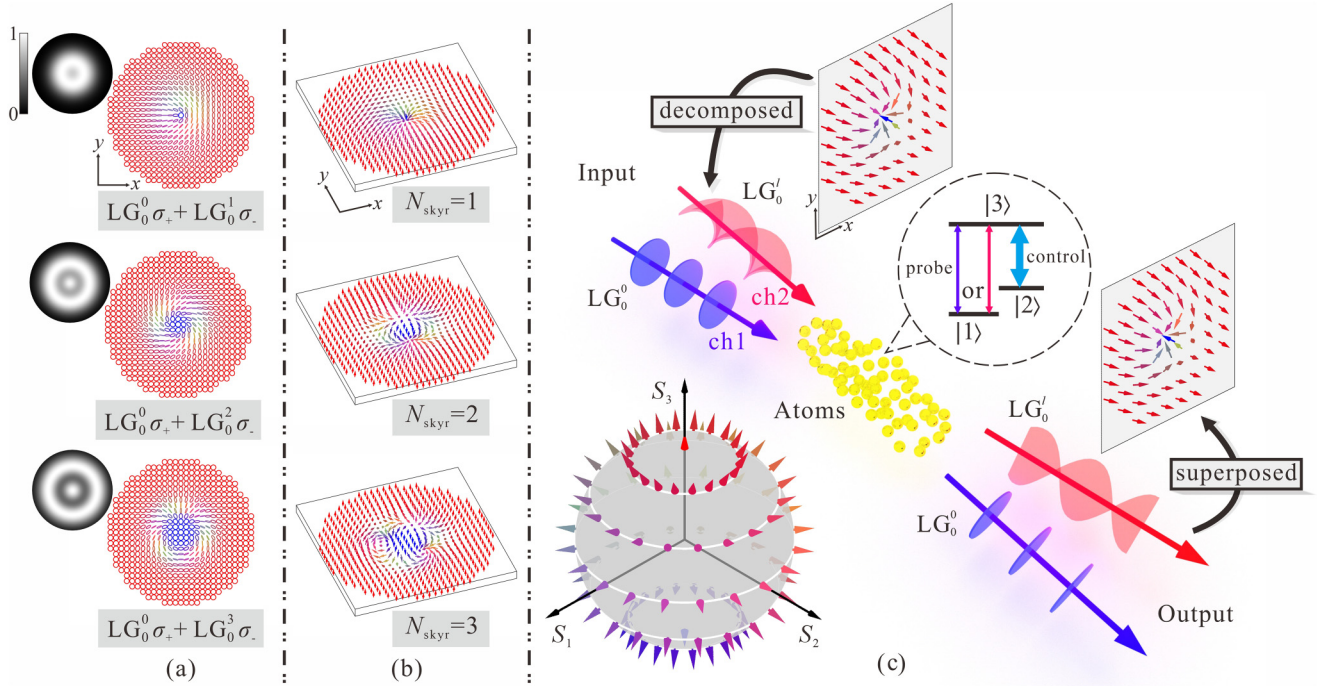


FIG. 1. (a) Simulated polarization and intensity distributions of paraxial optical skyrmions for $N_{\text{skyr}} = 1, 2, 3$ in Eq. (1). (b) Three-dimensional perspective of simulated Stokes vectors for optical skyrmions with the corresponding N_{skyr} values. (c) The schematic diagram of optical storage. The probe beam drives the level $|1\rangle$ to $|3\rangle$ in each path, and the control beam simultaneously drives the level $|2\rangle$ to $|3\rangle$ in both paths. We input the signals that need to be stored in the atomic storage medium, and obtain the output (retrieved) signals after storage. More details are in the main text. The left bottom inset shows the Stokes vector s on the unit Poincaré sphere. And the mapping relationship between arrows and polarizations are indicated by the same color.

and phase differences between the two paths, and still retains a global topological invariant and a robust quantum number.

In this work we realize the optical storage of optical skyrmions by employing a dual-path storage scheme based on an electromagnetically induced transparency (EIT) protocol. The two spatial modes of the optical skyrmion are spatially separated and independently stored within cold ^{87}Rb atoms. Differences in the intensity distribution of the two spatial mode profiles lead to an unavoidable imbalance in storage efficiency and a phase difference due to dispersion accumulated along the two paths. In addition, the system experiences a further decoherence-induced loss of storage efficiency. Our experimental results demonstrate, that nevertheless the skyrmion number is conserved for a storage time of several microseconds. While EIT-based storage is well established, here, we demonstrate that the topological invariant of an optical skyrmion beam remains robust, which is not achievable with conventional polarization or orbital angular momentum (OAM) encoding. The results of our study not only provide experimental evidence for the stability of a topologically nontrivial photonic encoding space against realistic memory imperfections but also offer valuable insights into future storage devices based on skyrmionic structures [54–56].

II. CONCEPT AND MODEL

To study the storage of optical skyrmions, we begin by generating states of the form

$$|\Psi^l\rangle = \frac{1}{\sqrt{2}}(|LG_0^0\rangle|\sigma_+\rangle + |LG_0^l\rangle|\sigma_-\rangle), \quad (1)$$

where p (the lower index is set to 0) and l denote the radial and azimuthal mode number of Laguerre-Gaussian (LG) modes [57], and σ_+, σ_- are left and right circular polarizations. Thus, the polarization structure depends on the relative local intensity and phase between the two modes, and thereby on the azimuthal quantum number l of the LG mode [58]. The simulated polarization distributions and Stokes vectors of paraxial optical skyrmions with corresponding skyrmion numbers $N_{\text{skyr}} = 1, 2$, and 3 are shown in Figs. 1(a) and 1(b). More details about calculation of optical skyrmions can be found in Appendix A.

As illustrated in Fig. 1(c), the input optical skyrmion is converted from a spatial mode and polarization-correlated state to a spatial mode and path-correlated state using conventional polarization optics shown in Fig. 2(a). The state is now

$$|\Psi^l\rangle = \frac{1}{\sqrt{2}}(|LG_0^0\rangle|\text{ch}_1\rangle + |LG_0^l\rangle|\text{ch}_2\rangle)\sigma_+, \quad (2)$$

a format that lends itself to dual-path EIT storage. Each spatial mode encoded with σ_+ polarization to interface with the atomic Λ -system [59], matching the control beam. The associated Poincaré sphere picture would feature the two paths as North and South pole, spanning complex superpositions of modal components along the two paths. In each storage path, the LG_0^l mode with a Rabi frequency Ω_l drives the level $|1\rangle$ to $|3\rangle$, and one shared control beam simultaneously drives the level $|2\rangle$ to $|3\rangle$. Under the slowly varying envelope approximation, the Maxwell equation for the

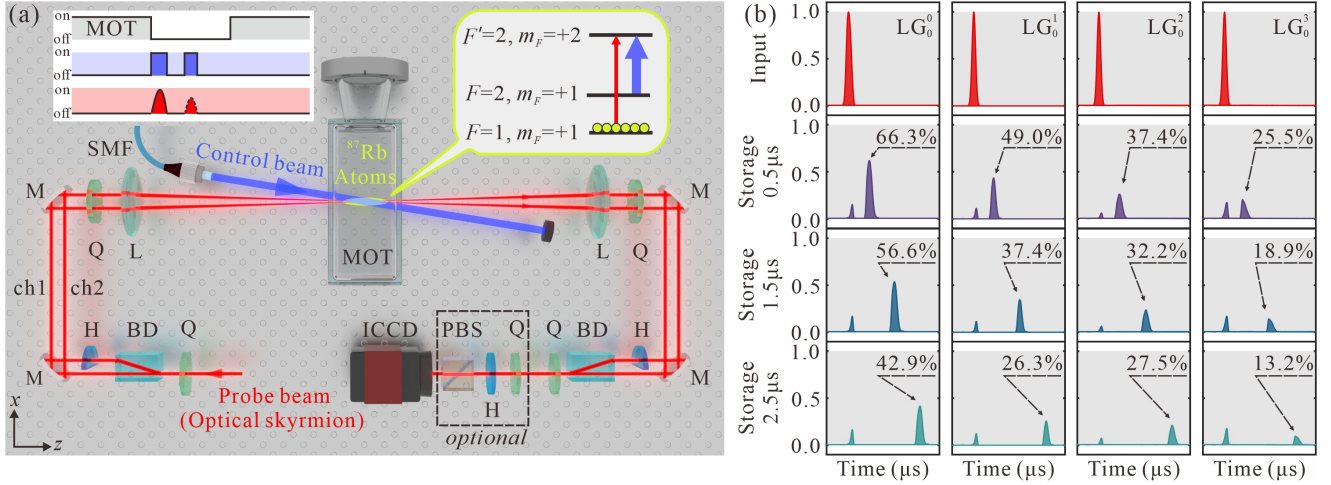


FIG. 2. (a) Experimental setup for the optical skyrmion storage. The control beam, with a waist size of 4 mm and a power of 33 mW, is incident on the atomic medium at a 1° angular separation relative to probe beams. By controlling the timing sequence of the MOT and optical switching (top left inset), as well as locking the frequency and phase of two beams (top right inset), the storage and retrieving processes are achieved. Full spatially resolved Stokes tomography (optional) is performed to characterize the polarization structure. Optical elements: H, half-wave plate; L, lens; M, mirror. (b) Temporal waveforms of the input LG beams, which both have Gaussian temporal profiles with $0.5 \mu\text{s}$ width, and their corresponding retrieved signals. The storage efficiencies for all paths are presented. The nonstored (smaller) remnants of the pulses are the leakage.

probe field is [59–61]

$$\frac{\partial \Omega_l}{\partial z} + \frac{1}{c} \frac{\partial \Omega_l}{\partial t} = i \frac{D_l \Gamma}{2L} \rho_{31}, \quad (3)$$

where Γ denotes the decay rate of $|3\rangle$, L is the length of the atomic medium, and ρ_{31} represents the atomic coherence between levels $|1\rangle$ and $|3\rangle$. D_l is the effective optical depth (OD) of the atomic medium for the LG_0^l mode. The beam waist of LG modes scales as $\sqrt{l+1}$, leading to a reduced OD and weaker light-atom interaction for higher order modes. It can be observed from Eq. (3) that D_l affects the storage performance of each path significantly, and consequently causes an imbalance in storage efficiency, acting as an effective noise source. After a specified storage time, the control beam triggers the read-out of these two components, which are recombined so that the topological structures of the (retrieved) output signals can be analyzed. By solving the Maxwell-Bloch equations, we derive the numerical relation between the storage efficiency and the l for each path. We further compare this with the concurrence [62] of vector beams, which are more resistant to perturbations than OAM modes [63], as detailed in Appendix C.

Although the dual-path storage protocol introduces asymmetric efficiency and hence is a nonunitary transformation which alters the superposition state of the skyrmions, the topological invariant N_{skyr} remains intact. This contrasts with conventional polarization-based storage, where efficiency imbalance or phase errors severely degrade reconstruction fidelity. The resilience of skyrmionic structures highlights their potential as information carriers compatible with imperfect photonic memory platforms. Based on this mechanism, we implement a topologically protected optical memory. The experimental setup and results are described in the following section.

III. EXPERIMENTAL SETUP

Various techniques have been developed for generating paraxial optical skyrmions [64–67]. Here we use a configuration comprising a segmented-screen spatial light modulator and an interferometer, as detailed in Appendix A. Our simplified experimental setup for optical storage is illustrated in Fig. 2(a). The generated optical skyrmions are then stored in a cigar-shaped cloud of cold ^{87}Rb atoms, with 3 cm length and 1.5 mm diameter, prepared within a two-dimensional magneto-optical trap (MOT) [68]. The relevant atomic Λ -system for the reversible mapping is based on three Zeeman sublevels of the D_1 line. Before storage, atoms are prepared in $|F=1, m_F=+1\rangle$ to maximize the absorption of the probe beams. Dual-path storage is realized by converting the optical skyrmion into a superposition of separate paths, as given in Eq. (2), where the usual orthogonal circular polarizations are unified to σ_+ and the entanglement is instead encoded in the paths. To achieve this, a quarter-wave plate (Q) and a polarizing beam displacer (BD) separate the spatial modes LG_0^0 and LG_0^l . Additional polarization operations allow us to obtain the desired polarization state. We also calibrate and optimize our system for equal mode profiles to ensure balanced OD in each path before storage, as detailed in Appendix E.

During the optical storage process, the temporal waveform of the probe pulse is tailored to a Gaussian shape with $0.5 \mu\text{s}$ duration, and the control beam is adiabatically switched off once the probe pulse fully enters the atomic ensemble, allowing the optical states to be coherently transferred into a collective spin excitation, as shown in the top left inset of Fig. 2(a). Upon reactivation of the control beam after a specific storage time, both components of the probe beam in the two paths are retrieved with their preserved spatial phase and intensity distribution, facilitated by a collective enhancement effect [69–71]. The temporal waveforms of both

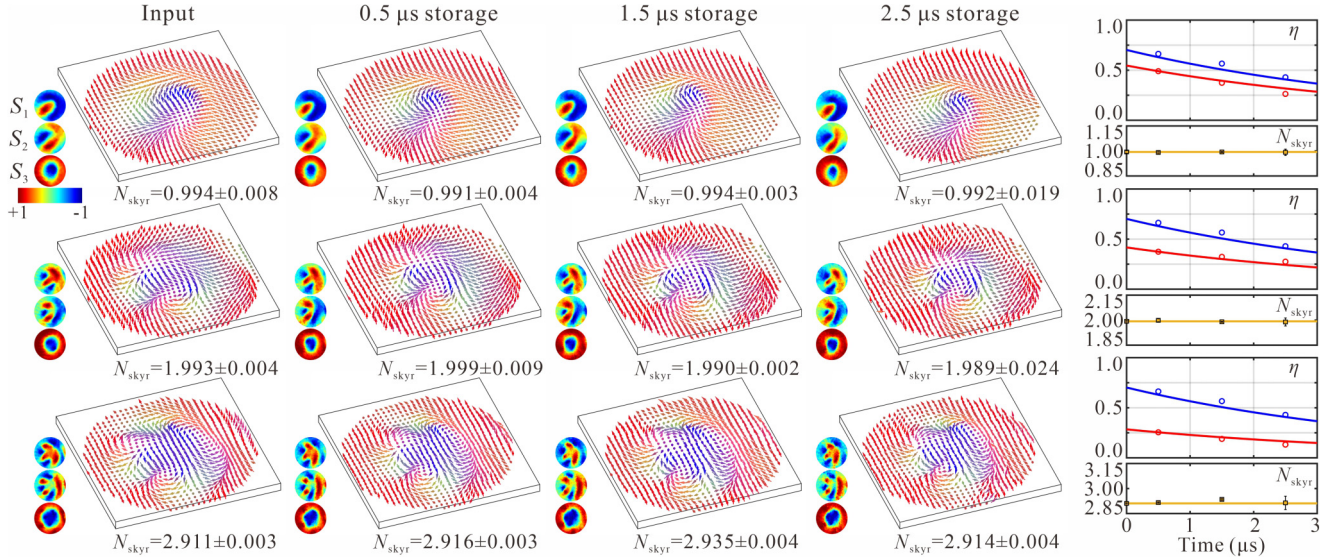


FIG. 3. Left: The measured topological textures of input and output optical skyrmions with $N_{\text{skyr}} = 1, 2$, and 3 (from top to bottom). The experimentally calculated N_{skyr} are presented with respect to each subfigure, with error bars representing the standard deviation of five runs. The insets on the side are the normalized local Stokes parameters. Right: Comparison between experimental data and theoretical simulation. The blue (red) solid lines represent the relationship of storage efficiency of the LG_0^0 (LG_0^1) mode against storage time, while the blue (red) circles represent the experimental data of the respective modes. The yellow solid line represents the measured input N_{skyr} , and the data points represent the experimental data after storage.

paths are obtained using photodetectors (not shown) to assess the storage efficiency. Lastly, by employing optical elements that unitarily reverse the decomposition process of the optical skyrmion, the retrieved skyrmion is reassembled. Its spatial profile is subsequently reconstructed and imaged on an intensified charge-coupled device (ICCD).

IV. EXPERIMENTAL RESULTS

Figure 2(b) shows the temporal waveforms and storage efficiencies for the different modes. As expected, the efficiency is highest for the LG_0^0 mode and decreases for higher-order modes (larger l) due to reduced light-atom overlap. The efficiency in each path decays with storage time due to spin-wave decoherence, which has been shown to substantially degrade the fidelity of stored information [72–74]. As a first step, we perform fully spatially resolved Stokes tomography on the states retrieved from the atomic memory. The Stokes parameters are recorded to reconstruct the spatial polarization profiles, as shown in Appendix F. The resulting Stokes vector profiles are shown in Fig. 3.

We calculate N_{skyr} of the input mode in the absence of the MOT, as well as the corresponding retrieved results after storing $0.5 \mu\text{s}$, $1.5 \mu\text{s}$, and $2.5 \mu\text{s}$. The experimental results show a minor deviation of N_{skyr} from their input values, and demonstrate the topological robustness of optical skyrmions under decoherence and the imbalanced storage efficiency between the paths, underscoring the intrinsic topological protection of skyrmionic configurations. Notably, slight deviations could be attributed to several factors, such as the imaging system, the background noise of our ICCD image, and the fluorescence noise of the atoms, yet the overall trend confirms resilience of N_{skyr} against decoherence over practical storage timescales. This contrasts with the decay of concurrence

observed in vector beams under similar perturbations, highlighting the superior topological protection afforded by the skyrmion number, as detailed in Appendix C. During storage, the optical state is mapped into a stationary atomic spin wave, as a result, there is no free-space propagation and hence no accumulation of relative Gouy phase [75]. As a result, the retrieved skyrmion retains the original topological texture without changing to a different skyrmion configuration. On the right side of Fig. 3, we also present a comparison between the theoretical simulation and the experimental results, as detailed in Appendix C.

We also study the impact of the control beam power on the stored optical skyrmions, illustrated by an example with $N_{\text{skyr}} = 2$ and a storage time of $0.5 \mu\text{s}$, as shown in Fig. 4. Changes in control beam intensity modify the medium's refractive response, which, owing to the differing spatial mode profiles of the two paths, results in distinct EIT responses. This alters the spatial polarization of the retrieved optical skyrmion, as detailed in Appendix D. As the control beam power increases, the stored skyrmion number decreases slightly, but the topological texture remains clear and intact, highlighting its inherent stability. Under the strong coupling regime, the loss in each path decreases and the delay between the paths shortens [76], while the imbalance in storage efficiency increases. The maximum power of 53 mW ($\approx 48.3 \text{ MHz}$) used in our experiments represents a practical upper bound set by our system constraints, and it also approaches the theoretical limit of the optimal control strength, as detailed in Appendix D. In theory, N_{skyr} remains quantized until the storage process completely fails for one path, beyond which the superposition is no longer coherent. This perturbation, while causing changes in local ellipticity, polarization rotations, and minor variations in skyrmion number, does not disrupt the overall topological structure [77]. We have

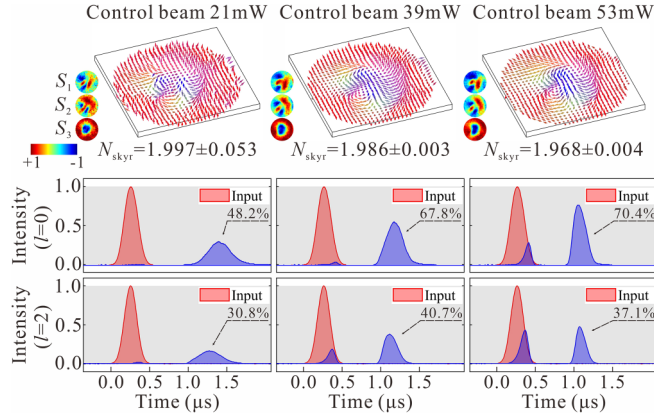


FIG. 4. Experimental results of storage under different control beam power, for an optical skyrmion with $N_{\text{skyr}} = 2$ and a storage time of $0.5 \mu\text{s}$. Top: The measured topological textures of input optical skyrmions with $N_{\text{skyr}} = 2$. The normalized local Stokes parameters and calculated N_{skyr} are in the legends on the left and below, respectively. The error bars represent the standard deviation of five runs. Bottom: The input and retrieved temporal waveforms of the input LG_0^0 and LG_0^2 modes. Additionally, the storage efficiencies are calculated and presented. The nonstored remnants of the pulses are the leakage.

therefore demonstrated that even large perturbations to the storage systems, by changing the control beam power by more than 100%, do not significantly perturb the topological structure of the retrieved skyrmion.

V. CONCLUSIONS

In summary, we have presented the first experimental demonstration that the topological characteristics of an optical skyrmion are preserved during coherent storage and retrieval in an optical memory. By using a dual-path EIT scheme, we directly demonstrated that the skyrmion number remains robust against experimental imperfections. This resilience is not merely the conservation of a classical number, but rather a direct manifestation of the topological protection inherent to the skyrmionic structure, which prevents the stored state from degenerating into a trivial configuration. Our work thus provides critical validation that the topological protection of skyrmions is not compromised by the complex light-matter interactions and decoherence processes central to quantum memories as long as the two-path superposition remains coherent.

Our primary focus in this work is to establish a proof-of-principle demonstration of skyrmion storage and to investigate the preservation of its topological invariant. Further work could optimize the memory performance itself or investigate the stability of skyrmion textures under varying control beam parameters. In particular, exploring the potential disruption mechanisms at high-power levels could provide valuable insights into the robustness limits of topological optical memory.

Looking forward, the significance of this result extends beyond the storage of a single topological charge. It establishes that the fixed skyrmion number defines a topologically protected subspace within which quantum information can be encoded. Future applications can exploit this stability by

encoding qubits in superpositions of states residing within this subspace, for instance, using different azimuth or radial indices while maintaining a fixed total skyrmion number [78–80], or even in superpositions of distinct skyrmion numbers themselves [81–83]. Our findings thereby open a pathway toward robust quantum photonic technologies by combining topological photonics and quantum memory. The demonstrated resilience against realistic noise in EIT-based storage marks a foundational step to leverage topological invariance as a resource for protecting quantum information [35–37,39].

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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.

APPENDIX A: GENERATION OF OPTICAL SKYRMIONS

A Gaussian beam with diagonal linear polarization and a waist size of 3 mm is split into its equal horizontal and vertical polarization components by a polarized beam displacer (BD). Since the spatial light modulator (SLM) can modulate only horizontally polarized light, the polarization of the vertically polarized beam is rotated by a half-wave plate (H) with the fast axis set to 45° . After passing through a further nonpolarizing beam splitter (BS), the spatial mode of each beam is individually shaped into the desired phase and amplitude profile by the SLM. Although such a configuration will result in a loss of light intensity, the outcome still meets our requirements for use. To achieve this, the SLM (HDSLM80R-Plus) is operated in a segmented-screen configuration, displaying two independently amplitude- and phase-modulated holograms [84], as

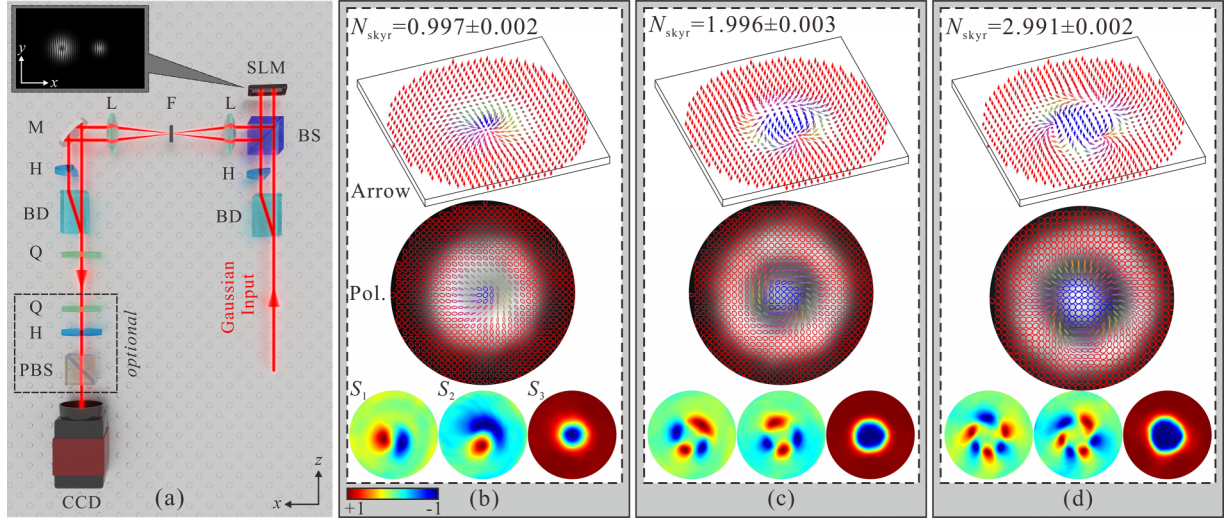


FIG. 5. (a) Experimental setup. BD: polarized beam displacer; H: half-wave plate; BS: beam splitter; SLM: spatial light modulator; L: lens with a 75 mm focal length; F: spatial filter; M: mirror; Q: quarter-wave plate; PBS: polarized beam splitter; CCD: charge-coupled device. The top left inset shows a representation of the holograms pattern displayed on the SLM for generating LG_0^0 and LG_0^1 modes. (b)–(d) The experimental results for $\frac{1}{\sqrt{2}}(LG_0^0\sigma_+ + LG_0^1\sigma_-)$, $\frac{1}{\sqrt{2}}(LG_0^0\sigma_+ + LG_0^2\sigma_-)$ and $\frac{1}{\sqrt{2}}(LG_0^0\sigma_+ + LG_0^3\sigma_-)$, respectively. Each part from bottom to top: Measured normalized Stokes parameters with corresponding colormap (left bottom); polarization (pol.) and intensity profiles; 3D view of equivalent Stokes vectors (arrow); The corresponding values of $\pm N_{\text{skyr}}$ in the top are the experimentally calculated skyrmion numbers. We use the colormap same as in Fig. 1 to reflect the mapping relationship between vector arrows and polarizations.

illustrated in Fig. 5(a). The shaped beams pass again through the BS and are then spatially filtered by a pinhole at the focus of a $4f$ lens system to remove undesired higher order modes. The two beams now have the desired orthogonal Laguerre-Gaussian (LG) modes, but are both horizontally polarized. In a reversal of the initial preparation steps, the polarization of one of the beams is rotated to be vertical by a suitable half-wave plate (H) and the two components are superimposed with a BD. A final quarter-wave plate (Q) transforms the polarization structure into the superposition in Eq. (1).

To confirm the successful generation of the desired skyrmion beams, We carry out a full Stokes analysis using combinations of a Q, an H and a polarizing beam splitter (PBS) shown as an optional part in Fig. 5(a). For appropriate settings of the polarization optics, we record the intensities in all three unbiased bases with a CCD camera, from which we obtain the Stokes parameters (S_0, S_1, S_2, S_3), where S_0 is the total intensity, S_1 describes the intensity difference between horizontal and vertical linear, S_2 between diagonal and anti-diagonal, and S_3 between right and left circular polarization components. The corresponding polarization ellipses of the light field can be obtained from the above Stokes parameters: The ellipticity is given by $\chi = \frac{1}{2} \arctan(S_3/\sqrt{S_1^2 + S_2^2})$, and the orientation angle by $\Psi = \frac{1}{2} \arctan(S_2/S_1)$. The spatially varying normalized Stokes vectors $\mathbf{s} = (S_1, S_2, S_3)$ play the role of an effective magnetization, with S_1, S_2 , and S_3 denoting the normalized local Stokes parameters. Formally, the skyrmion number N_{skyr} can be calculated from $\mathbf{s} = \mathbf{s}(x, y)$ as

$$N_{\text{skyr}} = \frac{1}{4\pi} \iint_{\sigma} \mathbf{s} \cdot \left(\frac{\partial \mathbf{s}}{\partial x} \times \frac{\partial \mathbf{s}}{\partial y} \right) dx dy, \quad (\text{A1})$$

for a quasiparticle confined to the surface σ , i.e., fully coherent light. The skyrmion number N_{skyr} is an integer that counts

how many times the vector $\mathbf{s}$ wraps around the unit sphere. Based on the mapping relation, Fig. 1(a), the polarization distribution and corresponding Stokes vector $\mathbf{s}$ profiles (topological texture) of arbitrary optical skyrmions can be obtained. In Figs. 5(b)–5(d) we display paraxial optical skyrmions generated in the experiment for $l = 1, 2, 3$, following Eq. (1).

The experimental results presented in here are acquired directly after the generation setup. The profile of the generated optical skyrmions exhibit excellent agreement with the numerical simulations as shown in Fig. 1. In contrast, the experimental results display in the main text are obtained after propagating along the optical paths required for storage and exhibit slight distortions after passing through various optical components, primarily due to distortions introduced by lenses, mirrors, and the vacuum chamber. However, as the errors indicate, the calculated skyrmion number N_{skyr} remains almost the same, with only minor discrepancies across all experimental measurements.

APPENDIX B: EIT AND ITS STORAGE APPLICATION

Electromagnetically induced transparency (EIT) is a seminal quantum interference phenomenon observed in coherently prepared atomic media [60]. It manifests as a narrow, highly transparent spectral window within an otherwise broad absorption profile of a material. This dramatic reduction in absorption is accompanied by an extremely steep normal dispersion, leading to a significant reduction in the group velocity of a propagating light pulse, an effect known as “slow light” [76]. The canonical EIT system involves a three-level atomic ensemble, typically in a Λ -configuration, as shown in Fig. 1(a). Two laser fields interact with this ensemble: a weak probe field, resonant with the transition from the ground state $|1\rangle$ to an excited state $|3\rangle$, and a strong control field, coupling

the metastable state $|2\rangle$ to the same excited state $|3\rangle$. Under conditions of two-photon resonance, the system evolves into a coherent superposition of the two ground states $|1\rangle$ and $|2\rangle$, known as a dark state [85]. This state has a vanishing dipole moment for the $|1\rangle$ to $|3\rangle$ transition, making the atoms transparent to the probe field. The destructive quantum interference between the two excitation pathways (direct excitation via the probe and indirect excitation via the control field) prevents the absorption of the probe photons.

The mechanism of light storage via EIT is a coherent process that enables the transfer of optical information into and out of a coherent atomic state [86]. This process begins with the propagation of a weak probe pulse, carrying the information (in our case LG_0^0 and LG_0^l modes) to be stored into an atomic medium prepared in a three-level Λ -configuration under the influence of a strong, continuous control field. The EIT condition creates a narrow transparency window, allowing the probe pulse to enter without absorption while its group velocity is drastically reduced, compressing its spatial extent within the atomic ensemble. Critically, the optical information is mapped onto the medium by adiabatically switching off the control field once the entire probe pulse is imprinted onto it. The photonic state is coherently converted into a collective atomic excitation, specifically a spin-wave coherence—a specific superposition of the atomic ground states that encodes the amplitude and phase of the original light pulse. To retrieve the information, the control field is reactivated after a programmable storage time, which reestablishes the EIT conditions and reverses the mapping process. Consequently, the spin-wave coherence acts as a source term, regenerating the probe pulse that then propagates out of the medium, faithfully preserving its quantum characteristics and temporal shape, thus completing the cycle of optical storage and retrieval. This retrieved pulse is a temporal and spatial replica of the input pulse at the moment it was stored, minus losses due to decoherence. It does not represent the field that would have resulted from propagating the input pulse through the length of the medium and then for an additional storage time. Thereby, the Gouy phase accumulation ceases during the storage time.

APPENDIX C: THEORETICAL ANALYSIS OF OPTICAL STORAGE VIA EIT

We consider an ensemble of Λ -type three-level atoms interacting with two classical laser fields (probe and control) in the dipole and rotating-wave approximations for each path. Following our definition in the main text, the Hamiltonian in the interaction picture is

$$\begin{aligned} \hat{H}_{\text{int}} = & -\frac{\hbar}{2}(\Omega_l|3\rangle\langle 1| + \Omega_c|3\rangle\langle 2| + \text{H.c.}) \\ & -\hbar\Delta_l|3\rangle\langle 3| - \hbar(\Delta_l - \Delta_c)|2\rangle\langle 2|, \end{aligned} \quad (\text{C1})$$

where Ω_l and Ω_c are the Rabi frequencies of the probe and control beams. Under two-photon and single-photon resonance ($\Delta_l = \Delta_c = 0$), the system evolves into a coherent superposition known as the dark state $|\text{Dark}_l\rangle$,

$$|\text{Dark}_l(t)\rangle = \frac{1}{\sqrt{|\Omega_c|^2 + |\Omega_l|^2}}(\Omega_c|1\rangle - \Omega_l|2\rangle). \quad (\text{C2})$$

This state is decoupled from the excited state $|3\rangle$, leading to a narrow transparency window for the probe beam. The steep dispersion associated with this transparency window dramatically reduces the group velocity, given by $v_g \approx 2|\Omega_c|^2/\alpha\Gamma$, where α is the linear intensity absorption coefficient of the medium at resonance ($\Delta_l = 0$) in the absence of the control field ($\Omega_c = 0$), and Γ is the decay rate of the excited state $|3\rangle$. Under the slowly varying envelope approximation ($\Omega_c \gg \Omega_l$), the Maxwell-Bloch equations can then be written as

$$\frac{d\rho_{31}}{dt} = -\gamma_{31}\rho_{31} + \frac{i}{2}\Omega_c\rho_{21} + \frac{i}{2}\Omega_l, \quad (\text{C3})$$

$$\frac{d\rho_{21}}{dt} = -\gamma_{21}\rho_{21} + \frac{i}{2}\Omega_c^*\rho_{31}, \quad (\text{C4})$$

$$\left(\frac{1}{c}\partial_t + \partial_z\right)\Omega_l = i\frac{D_l\Gamma}{2L}\rho_{31}. \quad (\text{C5})$$

Here ρ_{ij} are the slowly varying elements of the atomic density matrix (atomic coherence between levels $|i\rangle$ and $|j\rangle$), γ_{ij} is the decoherence rate of $|j\rangle$ to $|i\rangle$, and L is the length of the medium. Equation (C5) indicates that the spatial evolution of the probe field is driven by the atomic coherence ρ_{31} it induces. Thus, the efficiency η of this coherent mapping process is fundamentally limited by the effective optical depth D_l of the atomic ensemble. This is equivalent to the coupling strength between different LG_0^l beams and atomic clouds in our work. In the ideal adiabatic limit, the efficiency is given by [59,61]

$$\eta = \frac{\exp(-2\gamma_{21}D_l\Gamma/\Omega_c^2)}{\sqrt{1 + 32\ln 2 \frac{\gamma_{31}D_l\Gamma}{(T_p\Omega_c^2)^2}}} \times \frac{1}{2}(A + B), \quad (\text{C6})$$

$$A = \text{erf}(2\sqrt{\ln 2}\kappa), \quad (\text{C7})$$

$$B = \text{erf}\left(\frac{2\sqrt{\ln 2}(D_l\Gamma/(T_p\Omega_c^2) - \kappa)}{\sqrt{1 + 32\ln 2 \frac{\gamma_{31}D_l\Gamma}{(T_p\Omega_c^2)^2}}}\right), \quad (\text{C8})$$

where T_p denotes the full width at half maximum (FWHM) of signal pulse duration. κ is the proportional to the T_p when the control beam is switched off. This result underscores the importance of complete overlap and interaction between the light and the atomic ensemble.

As depicted in the main text, we assume a Gaussian transverse profile $n(r)$ for the atomic ensemble and an LG_0^l mode with $u_0^l(r)$ are given by

$$n(r) = n_0 \exp\left(-\frac{r^2}{w_a^2}\right), \quad (\text{C9})$$

$$u_0^l(r, \theta) = \frac{\sqrt{2^{1+|l|}}}{\sqrt{\pi}w_L|l|!}\left(\frac{r}{w_L}\right)^{|l|} \exp\left(-\frac{r^2}{w_L^2}\right) \exp(i l \theta), \quad (\text{C10})$$

where w_a represents the atomic cloud radius, w_L represents the radius of the LG_0^l mode, and n_0 denotes the peak atomic density. The storage efficiency is then fundamentally governed by the mode overlap integral A_{overlap} between the optical field and atomic distribution:

$$A_{\text{overlap}} = \int_0^{2\pi} \int_0^\infty n(r) |u_0^l(r, \theta)|^2 r dr d\theta. \quad (\text{C11})$$

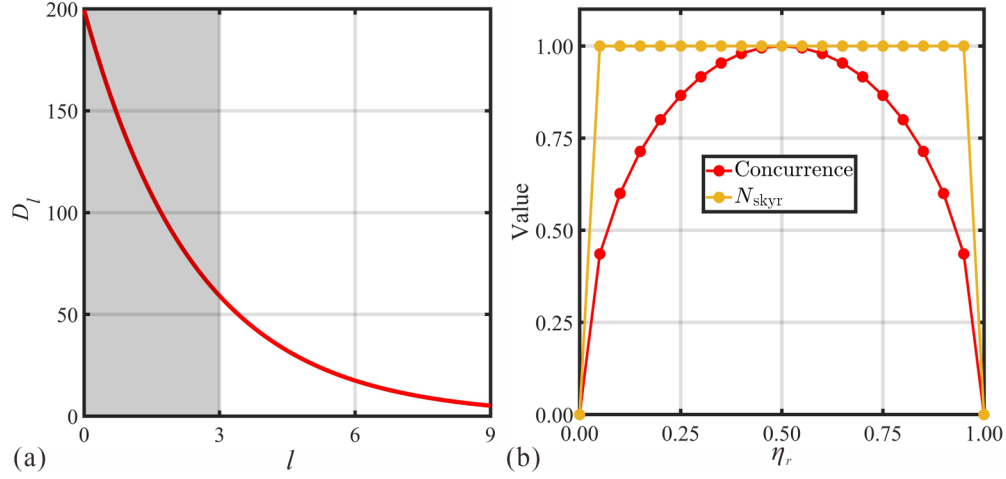


FIG. 6. (a) The simulation results of effective optical depth D_l as a function of l . We set $\beta = 1$. (b) The simulated normalized N_{skyr} and concurrence values vary with the relative storage efficiency η_r .

By defining the peak optical depth as $D_0 = \alpha n_0 L$ and appropriate normalization, the final expression is

$$D_l(l) = D_0 \left(\frac{2}{2 + \beta^2} \right)^{1+|l|}, \quad (\text{C12})$$

where $\beta = w_a/w_L$ represents the mode matching parameter. Combining Eqs. (C6) and (C12), we can calculate the effective optical depth of different OAM modes and the storage efficiency that changes with the storage time. The parameters used for Fig. 3 are $\Gamma = 2\pi \times 6.07$ MHz, $\gamma_{31} = \Gamma/2$, $\gamma_{21} = 5\Gamma \times 10^{-3}$, $\Omega_c = 3\Gamma$, $T_p = 300$ ns, $\kappa = 0.8$, and $\beta = 1$. Here we show the simulation results of effective optical depth D_l as a function of OAM, as illustrated in Fig. 6(a). The imbalance in storage efficiency corresponds to the change in the coefficients between the two modes [28], as $|\Psi^l\rangle = a|\text{LG}_0^0\rangle|\sigma_+\rangle + b|\text{LG}_0^l\rangle|\sigma_-\rangle$

($a = \sqrt{\eta_r}$, $b = \sqrt{1 - \eta_r}$, $\eta_r \in [0, 1]$ represents the relative storage efficiency). We calculated the N_{skyr} by changing the coefficients, and compared it with the concurrence (which is more resistant to perturbations compared to OAM [63]) of the vector beams, as illustrated in Fig. 6(b). This verified the topological protection characteristics of the N_{skyr} , as well as during optical storage.

APPENDIX D: IMPACT OF THE CONTROL BEAM POWER

The EIT storage requires that the control field variation satisfies the adiabatic condition [76]

$$\left| \frac{d\Omega_c}{dt} \right| \ll \frac{|\Omega_c|^2}{\gamma_{31}\sqrt{D_l}}. \quad (\text{D1})$$

The Rabi frequency of the control field Ω_c exhibits an optimal value, which can be determined by solving $\partial\eta/\partial\Omega_c = 0$ for

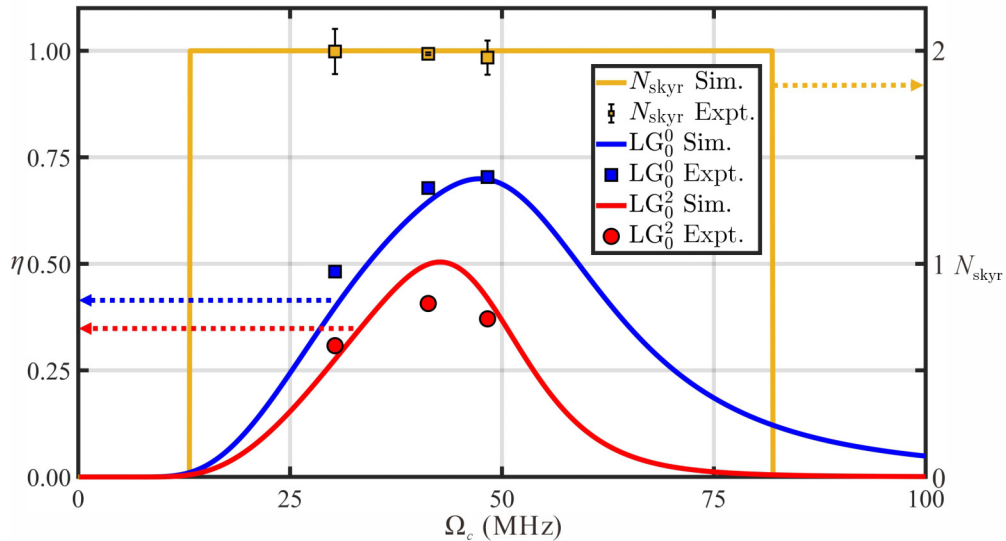


FIG. 7. Comparison between experimental data and theoretical simulation as a function of Ω_c . The blue (red, yellow) solid lines represent the relationship of LG_0^0 mode storage efficiency (LG_0^2 mode storage efficiency, N_{skyr}) against Ω_c , while marked points represent the experimental data.

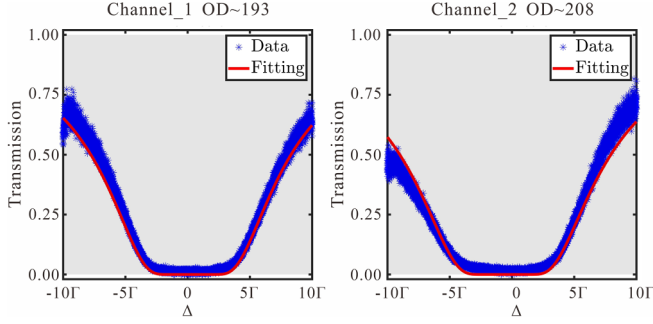


FIG. 8. Measured transmission (T) spectra for each channels versus probe detuning Δ from the atomic resonance $|5S_{1/2}, F=1\rangle \rightarrow |5P_{1/2}, F'=2\rangle$. The corresponding values in the top are obtained from the experimental fits according to Eq. (1).

Eq. (C6):

$$\Omega_c^{\text{opt}} \approx \sqrt{\frac{2\gamma_{31}D_l}{\tau}}. \quad (\text{D2})$$

In our storage experiment, a single control beam is used. Therefore, different EIT responses in the two paths occur due to spatial intensity profiles of the control beam. The control beam itself has a Gaussian spatial profile (waist of 4 mm). The two probe components (e.g., LG_0^0 and LG_0^1) have different intensity distributions and spatial overlaps with this control beam. The strength of the light-atom interaction for EIT, specifically, the Rabi frequency Ω_c experienced by atoms, is proportional to the local intensity of the control beam at the atom's position. Therefore, atoms located in different parts of

the cloud experience different Ω_c . The LG_0^0 probe mode, being more localized near the center of the control beam, interacts with atoms that, on average, experience a higher control Rabi frequency Ω_c . The larger LG_0^1 mode samples regions where Ω_c is weaker. The width of the EIT transparency window and the storage bandwidth scale with Ω_c . A higher Ω_c leads to a broader EIT window and generally higher efficiency. This is the fundamental reason for the differential response to a change in global control power: the two probe modes are effectively coupled into regions of the atomic ensemble with different effective Ω_c .

A recent publication [87] has systematically studied the impact of control beam on optical storage. Insufficient Ω_c violates the adiabatic condition, while increasing the control beam power significantly leads to two detrimental effects that prevent meaningful storage and retrieval. First, intense control beam causes substantial heating and optical pumping effects, effectively disrupting the cold atomic cloud required for coherent storage. This drastically reduces the optical depth and destroys the spin-wave coherence. Second, at extremely high control Rabi frequencies, additional atomic levels and nonlinear effects become relevant, and the simple Λ -system model breaks down. As a result, the EIT memory protocol itself ceases to function efficiently. It is therefore essential to optimize the control beam power within a reasonable range to balance storage efficiency and system stability. The maximum power of 53 mW ($\Omega_c \approx 48.3$ MHz) used in our experiments represents a practical upper bound imposed by our system constraints, while it also reaches the theoretical limit of optimal strength, as shown in Eq. (D2).

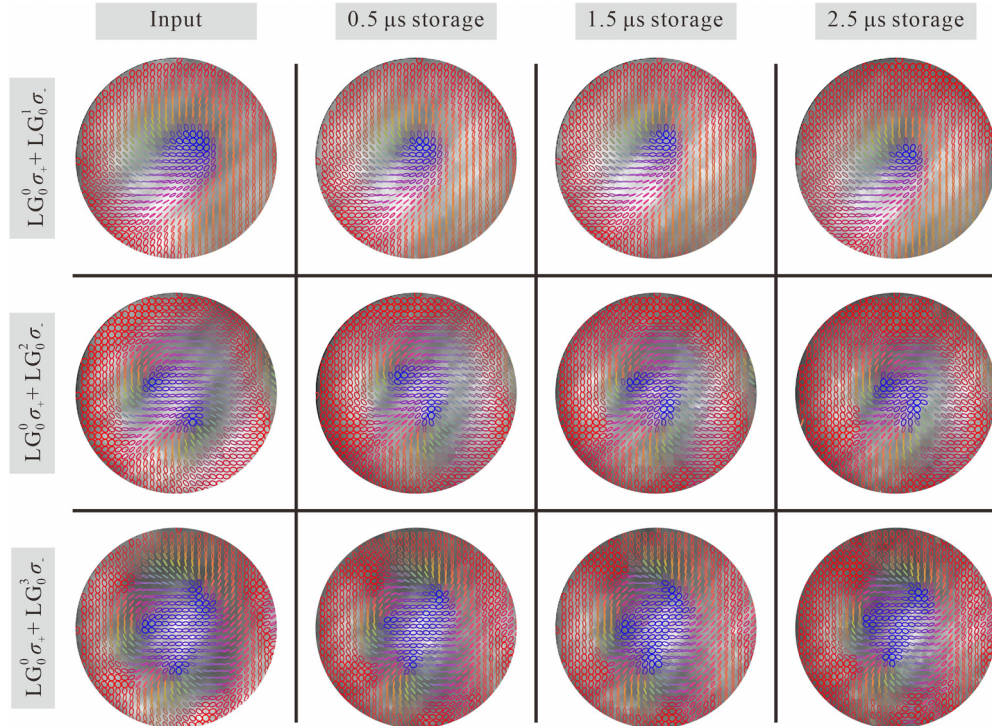


FIG. 9. The measured polarization distributions of input and after storage for 0.5 μs , 1.5 μs , and 2.5 μs for $\frac{1}{\sqrt{2}}(\text{LG}_0^0\sigma_+ + \text{LG}_0^1\sigma_-)$, $\frac{1}{\sqrt{2}}(\text{LG}_0^0\sigma_+ + \text{LG}_0^2\sigma_-)$ and $\frac{1}{\sqrt{2}}(\text{LG}_0^0\sigma_+ + \text{LG}_0^3\sigma_-)$, respectively. These results correspond to the Fig. 3, and the color map is the same as in Fig. 1.

Although this regime is not experimentally accessible, we can simulate it based on standard EIT theory (see Fig. 7). At powers exceeding those employed here, the storage efficiency imbalance between the two paths would be more extreme (as the differential response to Ω_c intensifies). The retrieved field therefore becomes a strongly unbalanced superposition of two modes, leading to severe distortion of the local polarization texture. Nevertheless, the topological skyrmion number N_{skyr} , as a global invariant robust to continuous deformations, remains quantized until the storage process completely fails (e.g., when the efficiency for one path drops to zero and the superposition is no longer coherent). Ultimately, at sufficiently high power, e.g., far beyond the theoretical optimum given in Eq. (D2), the memory mechanism itself would be destroyed, thus no coherent signal would be retrieved. Future work could systematically explore the relationship between the control beam parameters and the skyrmion stability, particularly the potential disruption mechanisms at elevated power levels. Achieving this will require an upgraded apparatus capable of safely exceeding the present power limits while maintaining precise control of other experimental parameters.

APPENDIX E: OPTICAL DEPTH MEASUREMENT

Optical depth (OD) serves as a fundamental parameter that quantitatively characterizes the strength of light-atom interactions. To ensure optimal overlap between the two probe channels and the atomic ensemble, we measure the OD using an LG_0^0 mode for each probe channel prior to the optical storage experiment. We adjust the beam alignment to maximize OD, thereby maximizing coupling efficiency.

The relationship between probe transmission T and OD can be expressed as

$$T = \exp \left[\text{Im} \left(\frac{\text{OD} \Delta \Gamma}{2\Delta(\Delta + i\Gamma/2)} \right) \right], \quad (\text{E1})$$

where Δ is the detuning of the probe beam, and $\Gamma = 2\pi \times 6 \text{ MHz}$ is the excited state decay rate. By scanning Δ from -60 MHz to $+60 \text{ MHz}$, we obtain the absorption curves in Fig. 8. For both channels, the measured OD values

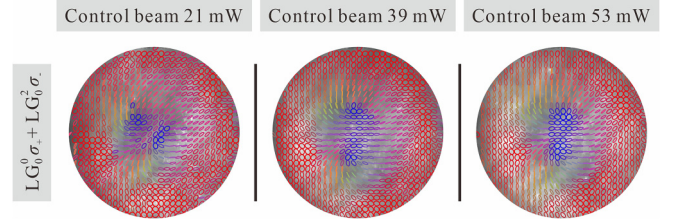


FIG. 10. The measured polarization distributions under different control beam power for $\text{LG}_0^0\sigma_+ + \text{LG}_0^2\sigma_-$ and a storage time of $0.5 \mu\text{s}$. These profiles correspond to the results presented in Fig. 4.

are typically around 200, confirming the strong light-atom coupling necessary for efficient EIT-based storage [59].

APPENDIX F: SPATIAL POLARIZATION PROFILES

Figure 9 shows the measured polarization distributions using recorded Stokes parameters after the storage optical path. Compared with input, some variations in both intensity and polarization distributions are observed between the input mode and the retrieved mode after optical storage. These modifications arise from two primary factors: first, spatial mode distortion induced by optical components (e.g., mirrors or lenses misalignment) in the storage path, which perturb the polarization structure of the optical skyrmion; second, interference fringes caused by the antireflection coating on the ICCD surface, introducing perturbation fringes in the detected profiles. However, even under experimental noise sources (e.g., ICCD background), the skyrmion number deviates minimally, confirming resilience.

Figure 10 shows the polarization distributions under different control beam power for $\frac{1}{\sqrt{2}}(\text{LG}_0^0\sigma_+ + \text{LG}_0^2\sigma_-)$ and a storage time of $0.5 \mu\text{s}$. Variations in the control beam intensity induce different EIT responses in the two storage channels due to their orthogonal spatial modes. This differential EIT effect modifies the atomic refractive index [88], thereby altering the spatial polarization structure of optical skyrmions. While such control-beam-induced optical activity can significantly perturb the polarization states of the probe beam, compromising the fidelity of encoded polarization information, the topological characteristics of the optical skyrmion remain remarkably preserved.

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