

Realization of backward retrieval in a Stark-modulated spin-wave quantum memory

Zhenqi Xu^{1,2}, Mucheng Guo^{1,2}, Weiye Sun^{1,2}, Pengjun Liang^{3,4,5,6}, Zongquan Zhou^{3,4,5,6}, Fudong Wang^{2,*}, Shuping Liu^{2,†} and Manjin Zhong^{2,‡}

¹*Southern University of Science and Technology, Shenzhen 518055, China*

²*International Quantum Academy and Shenzhen Branch, Hefei National Laboratory, Shenzhen 518048, China*

³*CAS Key Laboratory of Quantum Information, University of Science and Technology of China, Hefei 230026, China*

⁴*Anhui Province Key Laboratory of Quantum Network, University of Science and Technology of China, Hefei 230026, China*

⁵*CAS Center for Excellence in Quantum Information and Quantum Physics,*

University of Science and Technology of China, Hefei 230026, China

⁶*Hefei National Laboratory, University of Science and Technology of China, Hefei 230088, China*



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We report an experimental realization of backward retrieval in a spin-wave quantum memory based on a Stark-echo-modulated protocol in $\text{Eu}^{3+}:\text{Y}_2\text{SiO}_5$. By using Stark control, we preserve the full optical depth of the ensemble while suppressing coherent noise, enabling conditional storage fidelities above 97%. Our analysis shows that the present backward-retrieval efficiency is mainly limited by technical imperfections rather than by fundamental constraints. With realistic engineering improvements, backward retrieval in this protocol could move beyond the reabsorption-limited forward-emission regime. The protocol is also compatible with cavity-enhanced operation, offering an additional route toward higher efficiencies. These findings establish Stark-echo modulation as a practical and scalable route to high-efficiency, long-lived solid-state quantum memories.

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I. INTRODUCTION

Quantum networks aim to distribute quantum states among distant nodes, forming the basis for secure communication, high-precision sensing, and distributed quantum computation [1–4]. A central component of such networks is the quantum memory, which synchronizes probabilistic photon events by mapping flying qubits onto long-lived matter states and retrieving them on demand [5–9]. Rare-earth-ion (REI) doped crystals combine long optical coherence at cryogenic temperatures with spin states that can maintain coherence for orders of magnitude longer. This combination provides a powerful platform for efficient absorption and reemission of photonic excitations, as well as for long-lived, on-demand quantum storage [10–14]. Their spectral multiplexing capability [15,16] and compatibility with photonic integration [17,18] further enhance their suitability for high-performance quantum memories.

A variety of quantum memory protocols have been developed to harness the long-lived spin states of REI-doped crystals [10,19–23]. Exploiting these spin states requires transferring optical excitations to spin waves while

maintaining sufficient optical depth to ensure high storage and retrieval efficiency [23–25]. However, reabsorption, stemming from the directionality of the emitted echo, fundamentally limits the theoretical storage efficiency to 54%. Backward retrieval avoids reabsorption and can, in principle, reach near-unity efficiency [26]. Leading schemes, such as the spin-wave atomic frequency comb (AFC) [22,27] and noiseless photon echo (NLPE) [28–30], either reduce effective optical depth during spectral tailoring or impose geometric constraints that hinder true backward emission. Consequently, a practical quantum memory protocol that preserves full optical depth and enables low-noise backward retrieval remains elusive.

In this work, we experimentally demonstrate a Stark-echo (SE) modulated spin-wave quantum memory in $\text{Eu}^{3+}:\text{Y}_2\text{SiO}_5$, with both forward and backward retrieval [31]. This protocol combines the NLPE scheme [30,32] with Stark-control techniques [33,34], allowing deterministic scrambling and rephasing of the ensemble phases without sacrificing optical depth [33]. As a result, it suppresses coherent noise and provides the geometric flexibility needed to realize backward emission. We measure conditional storage fidelities exceeding 97% for both retrieval directions, identify key factors limiting backward efficiency, and outline clear routes for substantial improvements. These results establish Stark-echo modulation as a practical foundation for high-efficiency, scalable REI-based quantum memories suitable for future quantum network applications [28,35].

II. THEORY

The SE protocol exploits the inversion symmetry of the ensemble together with the linear Stark effect. In certain

*Contact author: fdwang.phys@foxmail.com

†Contact author: liushuping@iqasz.cn

‡Contact author: manjin.zhong@gmail.com

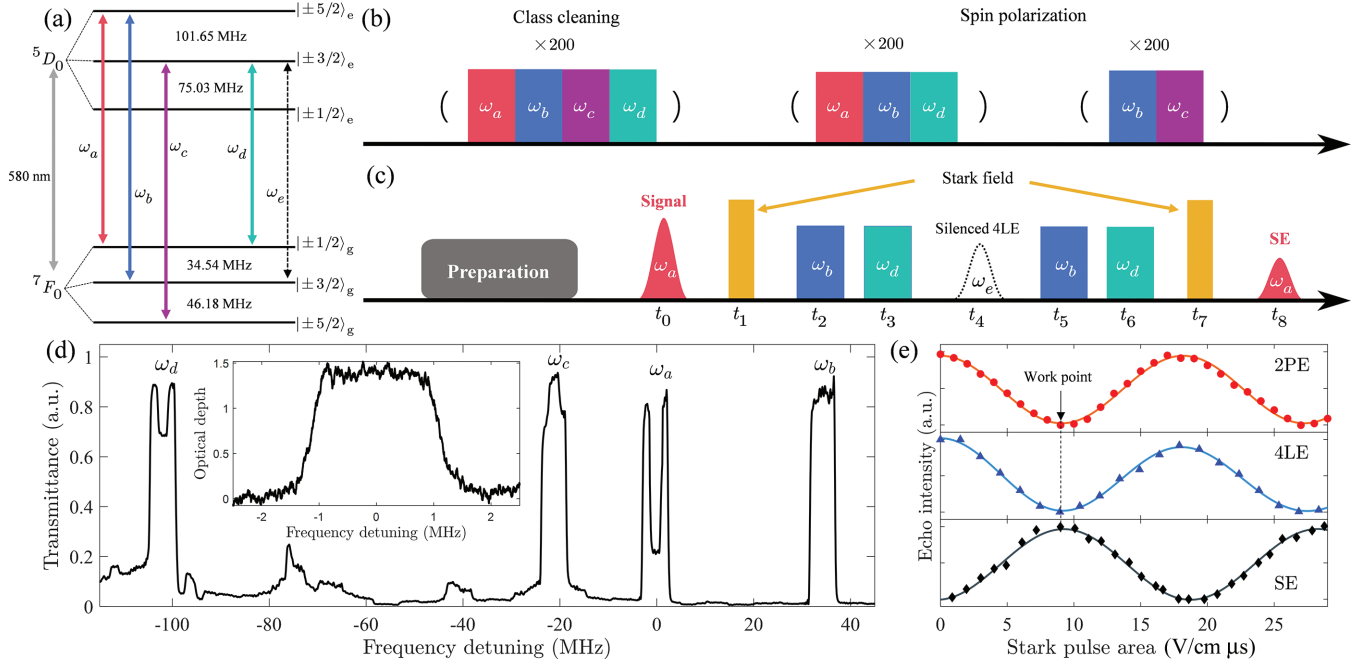


FIG. 1. (a) Energy-level structure of $^{151}\text{Eu}^{3+}$ used in the $\text{Eu}^{3+}:\text{Y}_2\text{SiO}_5$ crystal. (b) Preparation sequence including class cleaning and spin polarization. (c) Full quantum memory sequence showing the preparation process, four optical π pulses, and two electric-field pulses. (d) Optical spectrum after class cleaning and spin polarization; inset: the resulting absorption feature. (e) Stark modulation measurements via two-pulse photon echo (2PE), four-level photon echo (4LE), and the proposed Stark-echo (SE) method, respectively. Arrows and dashed lines highlight the electric-field pulse area of $9.25 \text{ V/cm } \mu\text{s}$ used for the memory operation. Data points indicate experimental values; solid curves are fitting results.

low-symmetry solid-state systems, inversion-related subensembles can experience linear Stark shifts of opposite signs under an applied electric field [36,37]. In $\text{Eu}^{3+}:\text{Y}_2\text{SiO}_5$, this gives rise to two classes of Eu^{3+} ions whose optical transition frequencies shift in opposite directions:

$$\Delta\nu = -L \Delta\boldsymbol{\mu} \cdot \mathbf{E}, \quad (1)$$

where $\Delta\boldsymbol{\mu}$ denotes the difference in permanent electric dipole moments between the two levels and L is the local-field correction factor [38]. Under a Stark pulse of duration τ_S and electric field $\mathbf{E}_S$, the two inversion-related subensembles accumulate opposite phases, leading to a relative phase difference [34]

$$2\phi_S = 2 \frac{\Delta\boldsymbol{\mu} \cdot \mathbf{E}_S}{\hbar} \tau_S. \quad (2)$$

This relative phase controls the collective emission of the ensemble: When $2\phi_S = \pi$, the two subensembles interfere destructively, and the echo is silenced, whereas when $2\phi_S = 2\pi$, they rephase and the collective emission is restored.

The SE protocol requires a four-level scheme, which is implemented here using the Eu^{3+} energy-level structure in $\text{Eu}^{3+}:\text{Y}_2\text{SiO}_5$, as shown in Fig. 1(a). Figure 1(b) shows the preparation sequence used to initialize the ions in $|\pm 1/2\rangle_g$, thereby enabling the implementation of the protocol. The pulse sequence shown in Fig. 1(c) consists of four optical π pulses interleaved with two Stark pulses. The optical coherence created by the input signal is transferred by the control π pulses into spin coherence. During the interval from t_2 to t_3 , the coherence is stored on the ground-state spin transition

$|\pm 1/2\rangle_g \leftrightarrow |\pm 3/2\rangle_g$, and from t_5 to t_6 it is transferred to the excited-state spin transition $|\pm 3/2\rangle_e \leftrightarrow |\pm 5/2\rangle_e$. The ground-state spin coherence thus provides the degree of freedom for long-lived storage. A Stark pulse applied at t_1 induces a phase difference of $2\phi_S = \pi$ between the two classes of ions, thereby scrambling the ensemble coherence and silencing the unwanted four-level photon-echo (4LE) emission [32]. A second Stark pulse at t_7 rephases the accumulated shifts to $2\phi_S = 2\pi$, producing an SE at t_8 whose emission direction is determined by the phase-matching condition [30]:

$$\mathbf{k}_{\text{echo}} = \mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_3 + \mathbf{k}_5 + \mathbf{k}_6, \quad (3)$$

where $\mathbf{k}_i$ denotes the wave vector of the optical pulse applied at time t_i . The emission time of the retrieved echo is given by $t_8 = t_0 - t_2 - t_3 + t_5 + t_6$.

If we adjust the propagation direction of the π pulse at t_5 or t_6 , the retrieved echo can be directed either forward or backward. Neglecting the inhomogeneity of the Stark pulses, we adopt an effective model based on Ref. [30] to describe the experimental data. In this simplified treatment, finite-bandwidth, spectral-dispersion, and propagation effects are neglected [39–41], and the total memory efficiency is given by

$$\eta_{\text{SE}} = \eta_{\text{retrieval}} \eta_{\text{pm}} \eta_{\text{control}}^4 \eta_{\text{decay}}, \quad (4)$$

where $\eta_{\text{retrieval}}$ is the retrieval efficiency, taking the form $\eta_{\text{retrieval}} = d^2 e^{-d}$ for forward retrieval and $\eta_{\text{retrieval}} = (1 - e^{-d})^2$ for backward retrieval, with $d = \alpha L$ being the optical depth. The factors η_{pm} , η_{control} , and η_{decay} represent the phase-matching efficiency, the population-transfer efficiency

of each π pulse, and the decoherence term, respectively. Decoherence includes both optical and hyperfine contributions and is used here as a practical fitting model for the measured decay curves:

$$\eta_{\text{decay}} = e^{-\frac{\Gamma_{13}^2(t_5-t_2)^2}{2\ln 2/\pi^2}} e^{-\frac{\Gamma_{35}^2(t_6-t_3)^2}{2\ln 2/\pi^2}} - 2\gamma(t_6-t_3), \quad (5)$$

where Γ_{13} and Γ_{35} are the inhomogeneous broadenings of spin transitions $|\pm 1/2\rangle_g \leftrightarrow |\pm 3/2\rangle_g$ and $|\pm 3/2\rangle_e \leftrightarrow |\pm 5/2\rangle_e$, respectively, and γ is the effective optical decoherence rate.

III. EXPERIMENT

The experiment was performed on a 0.1% $\text{Eu}^{3+}:\text{Y}_2\text{SiO}_5$ crystal with dimensions of $5 \times 5 \times 10 \text{ mm}^3$ along the D_1 , D_2 , and C_2 axes. We used the Eu^{3+} ions located at site 1 of the crystal, and the relevant hyperfine structure is shown in Fig. 1(a). All the optical transitions were excited with light propagating along the C_2 axis with the polarization along the D_1 axis. The electric fields were applied parallel to the D_1 axis via copper electrodes attached to the crystal using silver paint. The crystal was cooled to 1.8 K in a cryostat. The experiment employed a collinear counterpropagating configuration in which a fiber coupler was used to facilitate alignment. An acousto-optic modulator (AOM) was used to gate the detection window of the avalanche photodiode (APD), while two separate AOMs were used to generate and modulate the optical pulses. More details of the experimental setup are presented in Appendix A.

The preparation sequence, including class cleaning and spin polarization [42], is illustrated in Fig. 1(b). Using 1-ms frequency-chirped pulses with a peak power of 100 μW , this sequence produced the absorption feature shown in the inset of Fig. 1(d), with an optical depth of 1.3 and a bandwidth of 2 MHz. The linear Stark shift was then determined using two-pulse echoes (2PEs) [34] and 4LE sequences [32], as shown in Fig. 1(e). The measured Stark coefficient, $27.5 \pm 0.2 \text{ kHz}/(\text{V cm}^{-1})$, agrees well with previous reports [34,43]. SE rephasing was further verified by applying a fixed Stark pulse at t_1 and varying the area of the second pulse at t_7 [bottom panel of Fig. 1(e)].

IV. RESULTS AND DISCUSSION

A. Echo retrieval

To verify the operation of the SE protocol, forward retrieval was first experimentally demonstrated by direct detection of the SE intensity with an APD, as shown in Fig. 2(a). At early times, the SE was primarily obscured by the 4LE background from the input signal and π pulses, while the 2PE noise induced by the control pulses was suppressed by the electric field. By adjusting the relative timing of the π pulses, we separated the SE from the background and measured a memory efficiency of $4.8\% \pm 0.06\%$ at 29 μs . The relevant noise processes are analyzed in the next section.

Coherence dynamics were characterized by recording the SE decay curves via heterodyne detection, as shown in Fig. 2(b). Fitting the data using Eqs. (4) and (5) yielded linewidths of $\Gamma_{35} = 21.9 \pm 1.0 \text{ kHz}$ and $\Gamma_{13} = 17.4 \pm 1.2 \text{ kHz}$. The obtained Γ_{13} is consistent with the

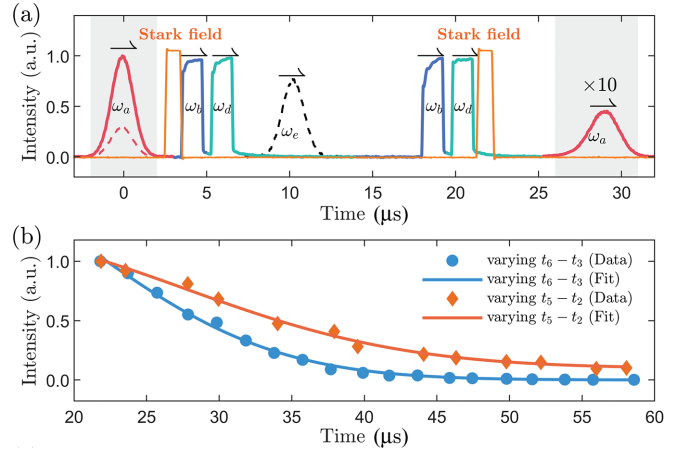


FIG. 2. Forward retrieval. (a) Experimental sequence with the applied propagation direction (shown as arrows) of each pulse. The black dashed curve represents the suppressed 4LE. The output signal in the shaded area is magnified by a factor of 10 for clarity. (b) SE decay curves measured by varying $t_6 - t_3$ (blue) and $t_5 - t_2$ (orange).

result of $16.1 \pm 0.9 \text{ kHz}$ determined from our independent all-optical measurements. The optical decoherence rate obtained from the 2PE measurement was $\gamma = 11.0 \pm 1.6 \text{ kHz}$, likely including contributions from spectral diffusion induced by the strong control pulses. Substituting these parameters and the measured efficiency into Eq. (4) and assuming $\eta_{\text{pm}} = 1$ for the copropagating geometry gives a control efficiency of $\eta_{\text{control}} = 82.8\%$.

To realize backward retrieval, we adopted the pulse sequence shown in Fig. 3(a). The propagation direction of the π pulse at t_5 was reversed to satisfy the phase-matching condition for backward retrieval in Eq. (3). In addition, the two π pulses at t_5 and t_6 were swapped, and the sequence timing was adjusted so that the background 4LEs were separated from the SE through a combination of spatial and temporal

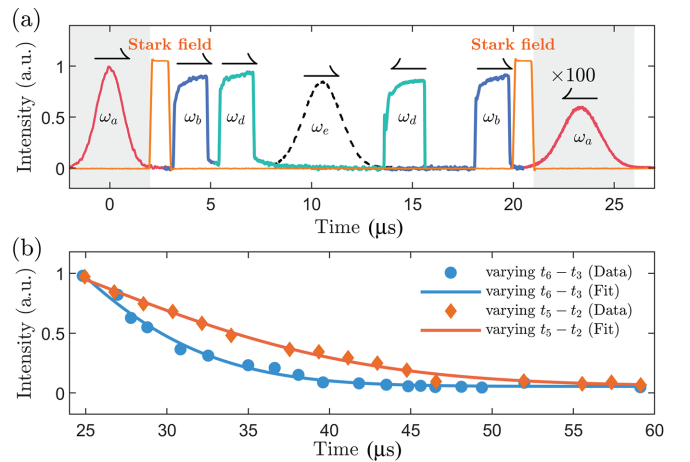


FIG. 3. Backward retrieval. (a) Experimental sequence with the propagation directions (shown as arrows) of different pulses and the silenced 4LE noise, shown by the black dashed line. The output signal in the shaded area is magnified by a factor of 100 for clarity. (b) SE decay profiles obtained by varying $t_6 - t_3$ (blue) and $t_5 - t_2$ (orange).

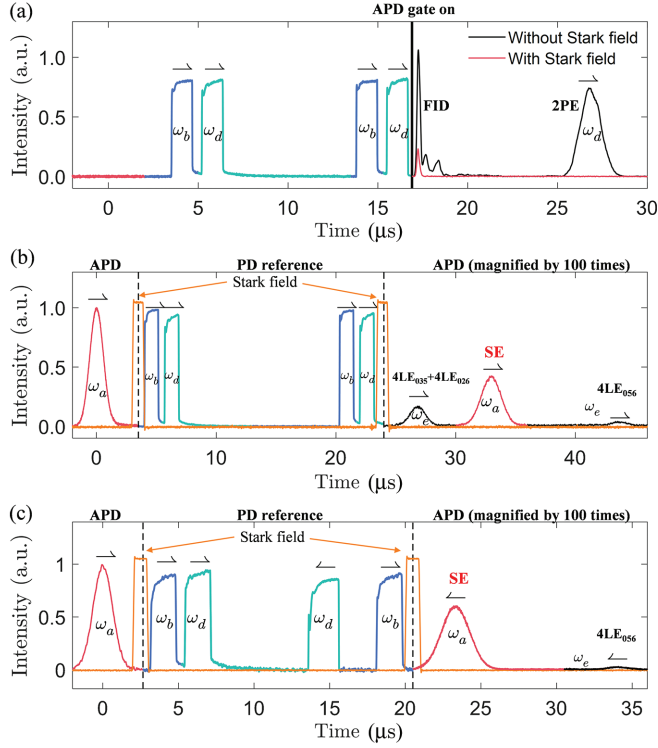


FIG. 4. Experimental noise analysis. (a) FID and 2PE noise produced without an input signal. (b) and (c) The measured 4LE noise in the forward and backward schemes, respectively.

discrimination. The corresponding noise pathways are analyzed in detail in the next section. The measured backward efficiency is $0.6\% \pm 0.01\%$ at 23 μ s. Fitting the echo-decay curves in Fig. 3(b) yields $\Gamma_{35} = 24.3 \pm 1.2$ kHz and $\Gamma_{13} = 16.6 \pm 1.1$ kHz, consistent with the corresponding values for forward retrieval. This agreement indicates that the decoherence dynamics are essentially independent of the retrieval direction. In this configuration, phase matching cannot be assumed to be perfect, meaning that η_{pm} does not necessarily reach unity. As a result, $\eta_{control}$ cannot be independently evaluated, and we can extract only the combined quantity $\eta_{pm}\eta_{control}^4 = 13.4\%$.

B. Noise analysis

In the absence of an input signal, both 2PEs and free-induction-decay (FID) noise originate solely from the applied π pulses. The 2PEs arise from pulse imperfections associated with the ω_b - ω_b and ω_d - ω_d pulse pairs. However, because the population on the ω_b transition is largely depleted during the state-preparation stage, the corresponding echo contribution is negligible. Applying an electric-field (Stark) pulse at t_7 strongly suppresses these coherent noise components. As shown in Fig. 4(a), the pulse nearly eliminates the 2PE noise and reduces the FID noise induced by the final π pulse to below 20% of its original level.

In practice, the control π pulses cannot achieve unit transfer efficiency, i.e., $\eta_{control} < 1$. As a result, the coherence created by the input signal is not completely transferred, and the residual coherence can be further driven by subsequent control pulses, giving rise to additional background. Apart

from the intended 4LE generated at t_4 , the input signal excitation can combine with any ω_b - ω_d or ω_d - ω_b π -pulse pair to form parasitic 4LE sequences. These unwanted 4LEs act as coherent noise and have optical frequencies different from those of the SE. Their emission directions satisfy the phase-matching condition [32]

$$\mathbf{k}_{4LE} = -\mathbf{k}_0 + \mathbf{k}_{\pi_1} + \mathbf{k}_{\pi_2}, \quad (6)$$

where $\mathbf{k}_{\pi_1}$ and $\mathbf{k}_{\pi_2}$ denote the wave vectors of the two π pulses involved in generating the corresponding 4LE. We denote a 4LE generated by the input signal at t_0 and the π pulses at t_m and t_n as $4LE_{0mn}$; for example, $4LE_{023}$ is generated by the input signal at t_0 and the π pulses at t_2 and t_3 .

In the forward-retrieval sequence shown in Fig. 4(b), three corresponding 4LE noise components, $4LE_{026}$, $4LE_{035}$, and $4LE_{056}$, can be generated and detected. According to the 4LE timing condition in Eq. (6), the SE and these noise components can be temporally separated by appropriately adjusting the pulse sequence. The first 4LE noise peak in Fig. 4(b) is formed by the overlap of $4LE_{026}$ and $4LE_{035}$. In contrast, the backward-retrieval configuration generates two backward-propagating 4LEs and one forward-propagating 4LE. The experimental results are shown in Fig. 4(c), where $4LE_{025}$ is not observed because it is emitted before the application of the Stark pulse. More details of the noise analysis are presented in Appendix C.

The total transfer efficiency of the SE involves all four π pulses, so its intensity scales as

$$I_{SE} \propto \eta_{control}^4. \quad (7)$$

By contrast, each 4LE noise component undergoes only two π pulses. Its intensity therefore scales as

$$I_{4LE} \propto (1 - \eta_{control})^2 \eta_{control}^2. \quad (8)$$

Details of the derivation are provided in Appendix B. Therefore, from the measured SE-to-4LE intensity ratio I_{SE}/I_{4LE} together with the corresponding decoherence factors, we infer a control-pulse efficiency of $\eta_{control} = 81.5\%$ for forward retrieval, in agreement with the value obtained from the preceding analysis. For backward retrieval, the same method yields $\eta_{control} = 85.3\%$. Using this inferred value together with the experimentally determined linewidths and decoherence rate in Eq. (4), we estimate the phase-matching efficiency to be $\eta_{pm} \approx 7.1\%$ for the backward-retrieval experiment. Increasing $\eta_{control}$ can effectively suppress the 4LE noise. In addition, a filter crystal can be used to eliminate the residual noise [29,44], enabling noiseless quantum storage.

C. Fidelity

The input time-bin coherent state was prepared in the superposition $|\psi\rangle = (|e\rangle + e^{i\Delta\alpha}|l\rangle)/\sqrt{2}$, where $|e\rangle$ and $|l\rangle$ denote the early and late temporal modes, respectively, and $\Delta\alpha$ is their relative phase. The two time bins were separated by a fixed temporal interval Δt , as shown in Fig. 5(a). The fidelities for the early and late basis states were calculated from the signal-to-noise ratio $F_{e,l} = (S + N)/(S + 2N)$, where S is the noise-subtracted signal and N is the noise [24]. The average fidelity for the computational-basis states is then given by $F_{poles} = (F_e + F_l)/2$.

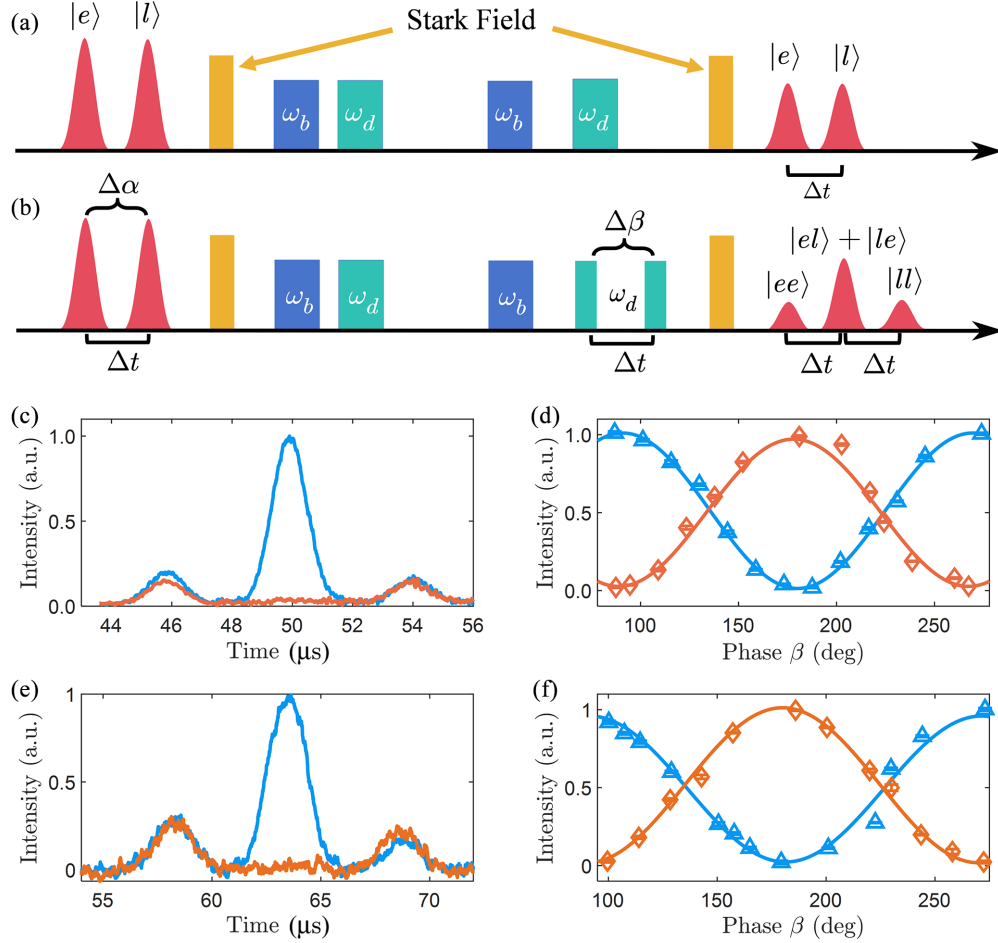


FIG. 5. Fidelity measurement. (a) illustrates the storage and retrieval process of the $|e\rangle$ and $|l\rangle$ time-bin states. (b) shows the interference measurement sequence used to evaluate the coherence between these two states. (c)–(f) correspond to the fidelity measurements of time-bin states for forward and backward retrieval, respectively. The time-bin delays are $\Delta t = 4 \mu\text{s}$ in (c) and (d) and $\Delta t = 5 \mu\text{s}$ in (e) and (f). (c) and (e) show the constructive and destructive interference fringes, while (d) and (f) show the corresponding interference curves for $\Delta\alpha = 0^\circ$ (blue) and $\Delta\alpha = 90^\circ$ (orange).

For superposition-state measurements, we employed an unbalanced Mach-Zehnder interferometer [24,30]. As shown in Fig. 5(b), the π pulse at t_6 was replaced by two $\pi/2$ pulses separated by a delay Δt with a relative phase $\Delta\beta$, enabling interference between the two time bins. From the measured interference fringes, we extracted the visibility V and determined the corresponding superposition-state fidelity as $F_{\text{equator}} = (1 + V)/2$. This procedure was performed for input states with $\Delta\alpha = 0^\circ$ and $\Delta\alpha = 90^\circ$, yielding two superposition-state fidelities. The corresponding interference results for forward and backward retrieval are shown in Figs. 5(c), 5(d), 5(e), and 5(f), respectively.

The overall conditional fidelity was obtained by combining the contributions from the basis-state and superposition-state measurements [24]:

$$F_T = \frac{1}{3}F_{\text{poles}} + \frac{2}{3}F_{\text{equator}}. \quad (9)$$

All results are summarized in Table I, yielding an overall conditional fidelity of $97.7\% \pm 0.8\%$ for forward retrieval and $97.5\% \pm 0.9\%$ for backward retrieval.

D. Discussion

The limited backward-retrieval efficiency can be primarily attributed to two experimental constraints. First, imperfect alignment and spatial mode overlap between the counter-propagating beams limit the phase-matching efficiency η_{pm} . Second, the available control-field power does not fully support complete population transfer, thereby reducing the overall control efficiency η_{control} . Both limitations can be addressed using integrated photonic structures. Waveguide confinement provides superior spatial mode overlap and significantly boosts the control-field intensities, enabling shorter

TABLE I. Fidelities of forward and backward retrieved states.

Fidelity	Forward	Backward
F_e	$98.4\% \pm 0.2\%$	$96.3\% \pm 0.4\%$
F_l	$96.7\% \pm 0.3\%$	$97.1\% \pm 0.4\%$
F_+	$97.0\% \pm 1.2\%$	$98.2\% \pm 1.0\%$
F_-	$98.6\% \pm 0.8\%$	$97.6\% \pm 1.3\%$
F_T	$97.7\% \pm 0.8\%$	$97.5\% \pm 0.9\%$

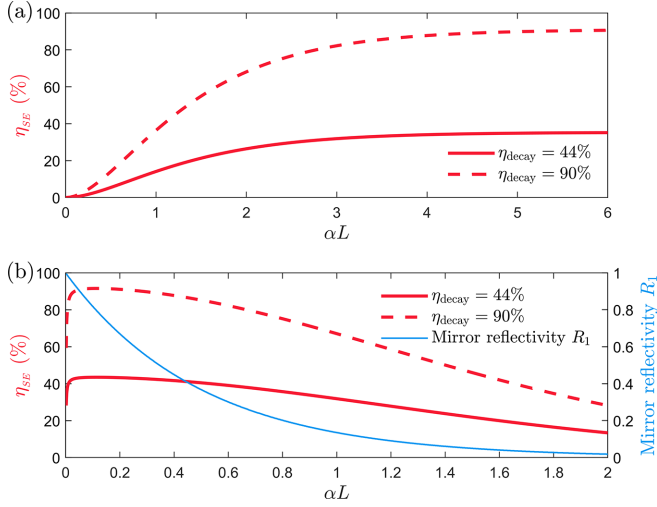


FIG. 6. Theoretical memory-efficiency estimation of (a) backward retrieval and (b) cavity enhancement, with different η_{decay} and under the assumption of $\eta_{control} = 1$ and $\eta_{pm} = 1$. The red curves in (b) represent the maximum achievable efficiency under different optical depths, while the blue curve denotes the first mirror reflectivity R_1 required to achieve this maximum efficiency. R_2 is set to a typical value of 0.999 [45].

control pulses and efficient adiabatic population-transfer techniques [46–48]. This approach can realistically push both η_{pm} and $\eta_{control}$ toward unity while further suppressing residual 4LE noise, as detailed in Appendix B.

Figure 6(a) shows the expected backward-retrieval efficiency η_{SE} , calculated using the experimentally determined $\eta_{decay} \approx 44\%$ at a storage time of 23 μs . Achieving high-efficiency backward retrieval primarily requires improving both $\eta_{retrieval}$ and η_{decay} [28,29]. Because backward emission eliminates reabsorption, $\eta_{retrieval}$ can be enhanced directly by increasing the optical depth. Recent advances in isotopically purified $^{151}\text{Eu}^{3+} : \text{Y}_2\text{SiO}_5$ indicated that substantially higher optical depths can be achieved without sacrificing coherence [28,29]. These developments suggest that stoichiometric rare-earth crystals may be a promising option for future implementations of the present scheme, owing to their potentially favorable combination of high optical depth and narrow linewidths [49,50].

Cavity enhancement provides a complementary route to improve the memory efficiency. We consider an optical cavity formed by an input coupling mirror and an end mirror, with reflectivities R_1 and R_2 , respectively. Under impedance-matching conditions, the transmission of the input coupling mirror matches the absorption of the intracavity medium, such that the incident field is efficiently coupled to the cavity and interacts repeatedly with the ensemble [45,51]. Owing to the destructive interference between the promptly reflected field and the intracavity leakage field, nearly all of the incident light can be absorbed by the sample, thereby effectively enhancing the light-matter interaction. If spectral dispersion and inhomogeneous broadening are neglected [52,53], Eq. (4) becomes

$$\eta_{retrieval} = \frac{4d^2 e^{-2d} (1 - R_1)^2 R_2}{(1 - \sqrt{R_1 R_2} e^{-d})^4}. \quad (10)$$

As shown in Fig. 6(b), cavity enhancement enables near-unity efficiency even at moderate optical depths, making it particularly suitable for low-doping systems in which η_{decay} is already high and the optical depth remains the primary bottleneck.

V. CONCLUSION

In summary, we reported an experimental demonstration of backward retrieval in a spin-wave quantum memory while retaining the full optical depth of the ensemble and suppressing coherent noise. SE modulation enables deterministic phase control without population removal, thereby providing geometrical flexibility and full optical-depth utilization, which are not readily available in AFC- or NLPE-based protocols. These results demonstrate the feasibility of high-fidelity backward retrieval and suggest a practical route toward higher efficiency through improved backward retrieval and its combination with cavity enhancement. Overall, our work identifies the SE protocol as a promising platform for efficient backward retrieval and highlights its potential for scalable, long-lived quantum network nodes based on rare-earth-ion materials.

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Z.X. and M.G. contributed equally to this work.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX A: DETAILS OF THE EXPERIMENTAL SETUP

As shown in Fig. 7, the crystal used in the experiment was mounted inside a cryostat and cooled down to 1.8 K. Copper plate electrodes were attached to the crystal surface using silver paint. The 580-nm continuous-wave laser beam passed through an optical isolator and then a polarization beam splitter (PBS), which divided the beam into two paths. A half-wave plate (HWP) was placed before the PBS to adjust the relative intensity between the two beams. After the PBS, each beam

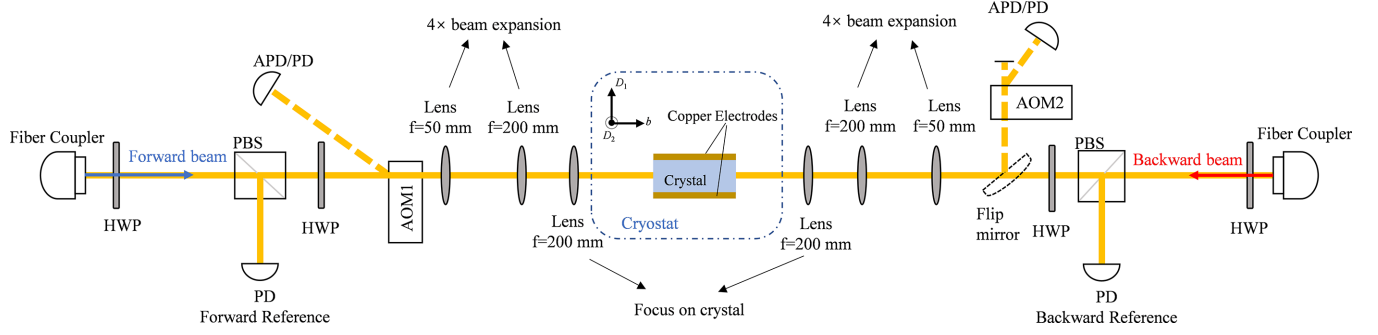


FIG. 7. Experimental setup. The yellow solid lines represent the main optical paths, while the dashed lines indicate the detection paths. The blue and red arrows denote the forward- and backward-propagating beams, respectively. AOM: acousto-optic modulator; HWP: half-wave plate; PBS: polarization beam splitter; APD: avalanche photodiode; PD: photodiode.

was sent into a double-pass AOM for pulse generation and frequency tuning. The modulated beams were then coupled to two optical fibers. At the output of each coupler, as shown in Fig. 7, three lenses with focal lengths of 50, 200, and 200 mm were placed sequentially before the beams entered the experimental crystal. The lenses expanded the beam by a factor of 4 and then focused it onto the crystal, resulting in a beam waist radius of 30 μm . Before the lens, a PBS and two HWPs were used to adjust laser power to reference and selected memory beam polarization. To achieve backward retrieval, after each beam passed through the experimental crystal, it was coupled to the coupler corresponding to the output of the other beam. The configuration ensured spatial overlap between the two beams. Because multiple lenses and other optical components were involved, the beam quality was degraded, leading to a coupling efficiency of only 25%. This efficiency is directly related to the degree of wave-vector matching. Two additional AOMs were placed in the optical path to gate the detectors and prevent potential damage caused by the strong control beams. For backward retrieval, AOM1 was used for gating and detection, while for forward retrieval, a flip mirror together with AOM2 was used to perform the gating and detection. Since the forward-retrieval configuration employs a copropagating and collinear optical setup, all pulses can be generated using a double-pass AOM (not shown in Fig. 7). The peak power of the π pulses entering the cryostat is 300 mW. In the backward-retrieval configuration, however, the optical path must be split into two beams, and two AOMs are used to generate the required pulses for the forward- and backward-propagating beams. Because the two AOMs have different frequency response characteristics, the optimized peak power of the π pulses in both directions is 170 mW.

APPENDIX B: THEORETICAL MODELING OF NOISE

We perform theoretical modeling of the noise and coherent population transfer efficiency in the Stark-echo (SE) modulated spin-wave quantum memory scheme. Consider an atomic ensemble with four energy levels, denoted as $|1\rangle$, $|3\rangle$, $|\bar{3}\rangle$, and $|\bar{5}\rangle$, where $|1\rangle$ and $|3\rangle$ represent the ground states and $|\bar{3}\rangle$ and $|\bar{5}\rangle$ represent the excited states. For example, in the Eu^{3+} -ion system used in our experiment [Fig. 1(a)], the

relevant hyperfine levels can be defined as follows:

$$\begin{aligned} |1\rangle &= |\pm 1/2\rangle_g, |3\rangle = |\pm 3/2\rangle_g, \\ |\bar{3}\rangle &= |\pm 3/2\rangle_e, |\bar{5}\rangle = |\pm 5/2\rangle_e. \end{aligned} \quad (\text{B1})$$

We assume that all atoms are distributed uniformly along the same propagation direction in the ensemble. In the Heisenberg picture, the quantum state of the j th ion at time t is described by a set of atomic operators $\hat{\sigma}_{\mu\nu}^{(j)}$, defined as [54]

$$\hat{\sigma}_{\mu\nu}^{(j)} = |\mu\rangle_{jj}\langle\nu| e^{i\omega_{\mu\nu}t}, \quad (\text{B2})$$

where $\mu, \nu = 1, 3, \bar{3}, \bar{5}$ correspond to the specific energy levels. When $\mu = \nu$, $\hat{\sigma}_{\mu\mu}$ represents the population operator of level $|\mu\rangle$; when $\mu \neq \nu$, $\hat{\sigma}_{\mu\nu}$ describes the coherence between levels $|\mu\rangle$ and $|\nu\rangle$ and satisfies $\hat{\sigma}_{\mu\nu} = \hat{\sigma}_{\nu\mu}^*$. The quantity $\omega_{\mu\nu}$ denotes the angular frequency difference between $|\mu\rangle$ and $|\nu\rangle$. To study the coherence transfer and phase-matching conditions in the system, the population of each level and the coherences between different levels can be expressed in terms of the atomic density matrix. After the state-preparation process at $t = t_0$, all atoms remain in the ground state $|\pm 1/2\rangle_g$, and the corresponding initial density matrix can be written as

$$\sigma(t_0^-) = \begin{pmatrix} \sigma_{11} & \sigma_{13} & \sigma_{1\bar{3}} & \sigma_{1\bar{5}} \\ \sigma_{31} & \sigma_{33} & \sigma_{3\bar{3}} & \sigma_{3\bar{5}} \\ \sigma_{\bar{3}1} & \sigma_{\bar{3}3} & \sigma_{\bar{3}\bar{3}} & \sigma_{\bar{3}\bar{5}} \\ \sigma_{\bar{5}1} & \sigma_{\bar{5}3} & \sigma_{\bar{5}\bar{3}} & \sigma_{\bar{5}\bar{5}} \end{pmatrix} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}. \quad (\text{B3})$$

During the subsequent evolution, the atomic coherence experiences two distinct processes: interaction with the control field and free evolution. It is assumed that each driving field affects only its corresponding transition and has a negligible effect on other transitions. After we apply the rotating-wave approximation [55,56] and assume that the applied pulses are “hard”, all relevant field-induced evolutions can be represented by a unitary evolution operator [56,57]:

$$|\psi'\rangle = U|\psi\rangle, \quad \sigma' = U\sigma U^\dagger. \quad (\text{B4})$$

An ideal π pulse completely exchanges the population between two energy levels. Since each optical pulse carries a specific wave vector $\mathbf{k}$, it introduces a spatially dependent phase factor. In quantum mechanics, a plane-wave field with

wave vector $\mathbf{k}$ introduces a term $e^{i\mathbf{k}\cdot\mathbf{r}}$ into the interaction Hamiltonian. Consequently, after the application of a π pulse, the density-matrix elements accumulate additional phases associated with this spatial dependence. In the present analysis, to account for imperfections, we consider that each π pulse drives its transition with an effective population transfer efficiency $\eta_{\text{control}} < 1$.

In the memory sequence, two types of π pulses are used to drive transitions $|3\rangle \leftrightarrow |\bar{5}\rangle$ and $|1\rangle \leftrightarrow |\bar{3}\rangle$. Their corresponding unitary evolution operators can be expressed as

$$U_{\pi}^{3\bar{5}} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \sqrt{1-\eta_{\text{control}}} & 0 & -i\sqrt{\eta_{\text{control}}}e^{i\mathbf{k}_{24}\cdot\mathbf{r}} \\ 0 & 0 & 1 & 0 \\ 0 & -i\sqrt{\eta_{\text{control}}}e^{-i\mathbf{k}_{24}\cdot\mathbf{r}} & 0 & \sqrt{1-\eta_{\text{control}}} \end{pmatrix}, \quad (\text{B5})$$

$$U_{\pi}^{1\bar{3}} = \begin{pmatrix} \sqrt{1-\eta_{\text{control}}} & 0 & -i\sqrt{\eta_{\text{control}}}e^{i\mathbf{k}_{13}\cdot\mathbf{r}} & 0 \\ 0 & 1 & 0 & 0 \\ -i\sqrt{\eta_{\text{control}}}e^{-i\mathbf{k}_{13}\cdot\mathbf{r}} & 0 & \sqrt{1-\eta_{\text{control}}} & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}, \quad (\text{B6})$$

where $\mathbf{r}$ denotes the position of a single atom.

During free evolution, each off-diagonal element $\sigma_{\mu\nu}$ accumulates a phase determined by the frequency detuning δ_i of energy level $|i\rangle$. This process can be represented by a diagonal

unitary operator describing the evolution over a delay τ :

$$U_{\tau} = \text{diag}(e^{-i\delta_1\tau}, e^{-i\delta_3\tau}, e^{-i\delta_5\tau}, e^{-i\delta_{\bar{5}}\tau}). \quad (\text{B7})$$

We specify that $\delta_{\mu\nu} = -\delta_{\nu\mu} = \delta_{\mu} - \delta_{\nu}$, and the spin transition detunings δ_{13} and $\delta_{3\bar{5}}$ are so small that they can be neglected in the following derivation.

We first consider the interaction between the signal field and the ions. After the input pulse, a coherence is established between energy levels $|1\rangle$ and $|\bar{5}\rangle$. If we consider a transfer efficiency of η_0 , the unitary evolution operator can be written as

$$U_0 = \begin{pmatrix} \sqrt{1-\eta_0} & 0 & 0 & -i\sqrt{\eta_0}e^{i\mathbf{k}_0\cdot\mathbf{r}} \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ -i\sqrt{\eta_0}e^{-i\mathbf{k}_0\cdot\mathbf{r}} & 0 & 0 & \sqrt{1-\eta_0} \end{pmatrix}. \quad (\text{B8})$$

After the input pulse, the density matrix becomes

$$\sigma(t_0^+) = U_0\sigma(t_0^-)U_0^\dagger = \begin{pmatrix} 1-\eta_0 & 0 & 0 & \sigma_0e^{i\mathbf{k}_0\cdot\mathbf{r}} \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ \sigma_0^*e^{-i\mathbf{k}_0\cdot\mathbf{r}} & 0 & 0 & \eta_0 \end{pmatrix}, \quad (\text{B9})$$

where $\sigma_0 = i\sqrt{\eta_0(1-\eta_0)}$ and σ_0^* denotes the complex conjugate of σ_0 .

Then, the complete evolution sequence of the density matrix at time T can be written as

$$\sigma(T) = U_{T-t_6}U_{t_6-t_5}U_{t_5-t_4}U_{t_4-t_3}U_{t_3-t_2}U_{t_2-t_1}U_{t_1-t_0}\sigma(t_0^+)U_{t_2-t_0}^\dagger U_{t_3-t_2}^\dagger U_{t_4-t_3}^\dagger U_{t_5-t_4}^\dagger U_{t_6-t_5}^\dagger U_{T-t_6}^\dagger. \quad (\text{B10})$$

Based on the above description, we can explicitly write out the key coherence transfers and phase evolutions in the density matrix at each step. Since only four coherence terms, σ_{13} , $\sigma_{1\bar{5}}$, $\sigma_{3\bar{3}}$, and $\sigma_{3\bar{5}}$, are mainly involved in the evolution process, we show only the derivation of these four matrix elements for clarity. The following notations are introduced:

$$\alpha = \sqrt{\eta_{\text{control}}}, \quad \beta = \sqrt{1-\eta_{\text{control}}}. \quad (\text{B11})$$

After a free-evolution period of $t_2 - t_0$, the π pulse $U_{\pi}^{3\bar{5}}$ is applied, transferring the coherence in $\sigma_{1\bar{5}}$ to σ_{13} . Due to the imperfect π pulse, a portion of the coherence remains in $\sigma_{1\bar{5}}$. After another free evolution during $t_3 - t_2$ and the application of the π pulse $U_{\pi}^{1\bar{3}}$ and a subsequent free evolution, the system evolves:

$$\begin{aligned} \sigma_{13}(t_4) &= \alpha\beta\sigma_0e^{-i\phi_{13}(t_4)}, & \phi_{13}(t_4) &= \delta_{1\bar{5}}(t_2 - t_0) + \delta_{1\bar{3}}(t_4 - t_2) - (\mathbf{k}_0 - \mathbf{k}_2) \cdot \mathbf{r}, \\ \sigma_{1\bar{5}}(t_4) &= \beta^2\sigma_0e^{-i\phi_{1\bar{5}}(t_4)}, & \phi_{1\bar{5}}(t_4) &= \delta_{1\bar{5}}(t_4 - t_0) - \mathbf{k}_0 \cdot \mathbf{r}, \\ \sigma_{3\bar{3}}(t_4) &= \alpha^2\sigma_0e^{-i\phi_{3\bar{3}}(t_4)}, & \phi_{3\bar{3}}(t_4) &= \delta_{1\bar{5}}(t_2 - t_0) + \delta_{3\bar{1}}(t_3 - t_2) + \delta_{3\bar{3}}(t_4 - t_3) - (\mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_3) \cdot \mathbf{r}, \\ \sigma_{3\bar{5}}(t_4) &= \alpha\beta\sigma_0e^{-i\phi_{3\bar{5}}(t_4)}, & \phi_{3\bar{5}}(t_4) &= \delta_{1\bar{5}}(t_3 - t_0) + \delta_{3\bar{5}}(t_4 - t_3) - (\mathbf{k}_0 - \mathbf{k}_3) \cdot \mathbf{r}. \end{aligned} \quad (\text{B12})$$

At this moment, the phase $\phi_{3\bar{3}}(t_4)$ can be rewritten as

$$\phi_{3\bar{3}}(t_4) = \delta(t_2 - t_0 + t_3 - t_4) + (\mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_3) \cdot \mathbf{r}.$$

When $t_2 - t_0 + t_3 - t_4 = 0$, phase rephasing occurs, and the echo-emission condition is satisfied, with the corresponding wave-vector matching condition

$$\mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_3 + \mathbf{k}_4 = 0,$$

which corresponds to the 4LE emission condition [32]. However, the Stark field applied at t_1 destroys this phase-matching condition, resulting in a silenced echo, while the corresponding coherence remains and continues to evolve in the subsequent pulse sequence.

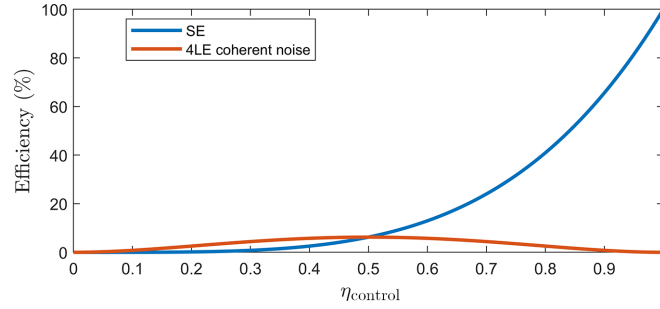


FIG. 8. Theoretical analysis of the 4LE noise level and SE level for different coherent transfer efficiencies of control π pulses.

After the π pulses U_{π}^{35} and U_{π}^{13} are applied in sequence and associated free evolution occurs and all nonrephasing terms are eliminated, the final density matrix at time T ($T > t_7$) can be expressed as

$$\sigma(T) = \begin{pmatrix} (1 - \eta_0)^2(1 - 2\alpha^2)^2 & 0 & 0 & \alpha^4\sigma_0A \\ 0 & 4\eta_0^2\alpha^2\beta^2 & \alpha^2\beta^2\sigma_0^*B^* & 0 \\ 0 & \alpha^2\beta^2\sigma_0B & 4(1 - \eta_0)^2\alpha^2\beta^2 & 0 \\ \alpha^4\sigma_0^*A^* & 0 & 0 & \eta_0^2(1 - 2\alpha^2)^2 \end{pmatrix}, \quad (\text{B13})$$

where $A = e^{i\phi_{15}(T)}$ and $B = e^{i\phi_{33}^{(1)}(T)} + e^{i\phi_{33}^{(2)}(T)} + e^{i\phi_{33}^{(3)}(T)}$. The phase shift induced by the Stark pulse applied at t_1 is canceled by the Stark pulse applied at t_7 . If we neglect the spin-detuning terms, the corresponding phase can be written as

$$\begin{aligned} \phi_{15}(T) &= \delta_{15}(t_2 - t_0) + \delta_{33}(t_5 - t_3) + \delta_{15}(T - t_6) - (\mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_3 + \mathbf{k}_5 + \mathbf{k}_6) \cdot \mathbf{r}, \\ \phi_{33}^{(1)}(T) &= \delta_{15}(t_3 - t_0) + \delta_{33}(T - t_5) - (\mathbf{k}_0 - \mathbf{k}_3 - \mathbf{k}_5) \cdot \mathbf{r}, \\ \phi_{33}^{(2)}(T) &= \delta_{15}(t_2 - t_0) + \delta_{33}(T - t_6) - (\mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_6) \cdot \mathbf{r}, \\ \phi_{33}^{(3)}(T) &= \delta_{15}(t_5 - t_0) + \delta_{33}(T - t_6) - (\mathbf{k}_0 - \mathbf{k}_5 - \mathbf{k}_6) \cdot \mathbf{r}, \end{aligned} \quad (\text{B14})$$

where $\phi_{15}(T = t_8)$ satisfies the echo emission and $t_8 = t_0 - t_2 - t_3 + t_5 + t_6$, with the corresponding wave vector

$$\mathbf{k}_{\text{echo}} = \mathbf{k}_0 - \mathbf{k}_2 - \mathbf{k}_3 + \mathbf{k}_5 + \mathbf{k}_6.$$

This emitted echo corresponds to the SE, whose coherence amplitude scales as $\alpha^4 = \eta_{\text{control}}^2$, leading to an echo intensity proportional to η_{control}^4 .

A 4LE generated by the input at t_0 and π pulses at t_m and t_n is denoted as 4LE_{0nm} (for example, 4LE₀₂₃ is generated by the input at t_0 and the π pulses at t_2 and t_3). When the condition $T = t_5 + t_3 - t_0$ is satisfied, rephasing occurs in $\phi_{33}^{(1)}(T)$, resulting in the emission of the 4LE₀₃₅ coherent noise. Similarly, when $T = t_6 + t_2 - t_0$, rephasing occurs in $\phi_{33}^{(2)}(T)$, and the 4LE₀₂₆ coherent noise is generated. When $T = t_6 + t_5 - t_0$, rephasing takes place in $\phi_{33}^{(3)}(T)$, leading to the emission of the 4LE₀₅₆ coherent noise. The intensity-efficiency coefficients of these three noise echoes are all given by $(\alpha^2\beta^2)^2 = \eta_{\text{control}}^2(1 - \eta_{\text{control}})^2$.

Based on the results presented in Eq. (B13) and Fig. 8, the efficiency factors for the SE intensity and the 4LE noise intensity can be written as η_{control}^4 and $\eta_{\text{control}}^2(1 - \eta_{\text{control}})^2$, respectively. Following the efficiency expression used in the noiseless photon echo [30], the overall efficiency of the SE process in $\text{Eu}^{3+} : \text{Y}_2\text{SiO}_5$, neglecting the imperfection of electric field, can be expressed as

$$\eta_{\text{SE}} = \eta_{\text{retrieval}} \eta_{\text{pm}} \eta_{\text{control}}^4 \exp\left[-\frac{\Gamma_{13}^2(t_5 - t_2)^2}{2 \ln 2 / \pi^2}\right] \exp\left[-\frac{\Gamma_{35}^2(t_6 - t_3)^2}{2 \ln 2 / \pi^2} - 2\gamma(t_6 - t_3)\right]. \quad (\text{B15})$$

Here, $\eta_{\text{retrieval}}$ is the retrieval efficiency, which takes the form $\eta_{\text{retrieval}} = d^2 e^{-d}$ for forward retrieval and $\eta_{\text{retrieval}} = (1 - e^{-d})^2$ for backward retrieval, where $d = \alpha L$ is the optical depth, with α and L denoting the absorption coefficient and the crystal length, respectively. The term η_{pm} represents the phase-matching efficiency. Γ_{13} and Γ_{35} are the inhomogeneous broadenings of spin transitions $|1\rangle \leftrightarrow |3\rangle$ and $|3\rangle \leftrightarrow |\bar{5}\rangle$ [58], respectively. $\gamma = \gamma_{13} + \gamma_{35}$, where γ_{33} and γ_{15} represent the effective decoherence rates of optical transitions $|1\rangle \leftrightarrow |\bar{3}\rangle$ and $|3\rangle \leftrightarrow |\bar{5}\rangle$, respectively [30]. The decoherence of the 4LE process is mainly determined by γ . According to the phase evolution in Eq. (B14), the efficiency of the coherent noise can be expressed as

$$\eta_{\text{noise}} = \eta_{\text{retrieval}} \eta_{\text{pm}} \eta_{\text{control}}^2 (1 - \eta_{\text{control}})^2 \exp[-\gamma(T_{\text{noise}} - t_{\pi_2} + t_{\pi_1} + t_0)], \quad (\text{B16})$$

where t_{π_1} and t_{π_2} denote the times of the two π pulses in the 4LE subsequence and T_{noise} is the time when the 4LE

noise occurs. Here, the spin detuning and spin dephasing are neglected, so Eq. (B16) should be understood as a simplified

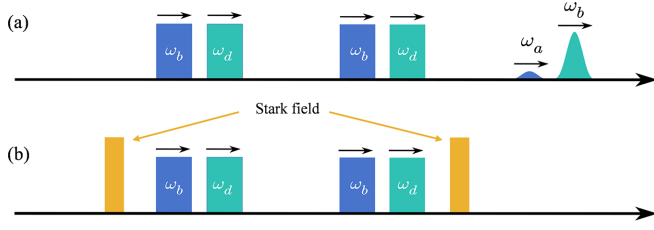


FIG. 9. Noise analysis without an input pulse. (a) The sequence illustrating the two-pulse photon-echo (2PE) noise, where each color corresponds to the noise generated by pulses of the same color. (b) Schematic diagram showing the suppression of noise by the applied Stark field.

effective expression for the coherent 4LE noise rather than a universal decay formula. If we combine Eqs. (B15) and (B16), η_{control} can be estimated from the experimentally measured SE-to-4LE ratio.

In summary, three coherent 4LE noise components appear in this storage scheme. Without considering decoherence, their intensities relative to the SE follow the theoretical dependence shown in Fig. 8. Therefore, the experimentally measured noise-to-readout intensity ratio can be used to determine the value of η_{control} . As illustrated in Fig. 8, a higher η_{control} leads to lower noise levels and hence a better SE quality.

APPENDIX C: NOISE CHARACTERIZATION

As shown in Fig. 9, due to pulse imperfections, 2PEs can be generated by the ω_b - ω_b and ω_d - ω_d pulse pairs. However,

since the population in the ω_b transition is depleted during the preparation process, the corresponding echo intensity is negligible. Theoretically, these coherent noise components are expected to be suppressed by the electric-field pulse applied at t_7 .

In the 4LE sequence [32], an input pulse and two π pulses are involved. The phase-matching conditions for the 4LE are given in Eq. (6). The corresponding 4LE emission time is determined by the arrival times t_{π_1} and t_{π_2} of two π pulses and is given by $t_{4\text{LE}} = t_{\pi_1} + t_{\pi_2} - t_0$ [32]. The notation for the 4LE noise components was already introduced in Appendix B.

There are four possible subsequences in the experimental sequence that can generate 4LEs, with optical frequencies different from that of the SE. Among them, 4LE_{023} generated at t_4 is suppressed by the electric field applied at t_1 and subsequently reappears as an echo revival at t_8 . It therefore does not contribute to the coherent noise in our experiment.

The remaining three 4LEs originate from imperfect π pulses, are rephased by the electric field applied at t_7 , and are emitted as noise. In the forward-retrieval configuration, all three of these 4LEs are emitted in the same direction as the SE. As shown in Fig. 10(a), the corresponding 4LEs are 4LE_{026} , 4LE_{035} , and 4LE_{056} . In contrast, in Fig. 10(b), where the π pulses at t_5 and t_6 are swapped, the three 4LEs become 4LE_{025} , 4LE_{036} , and 4LE_{056} .

In the backward-retrieval configuration, only two 4LEs are emitted in the same direction as the SE. As shown in Figs. 10(c) and 10(d), different π -pulse configurations lead to different generation times of the 4LE noise.

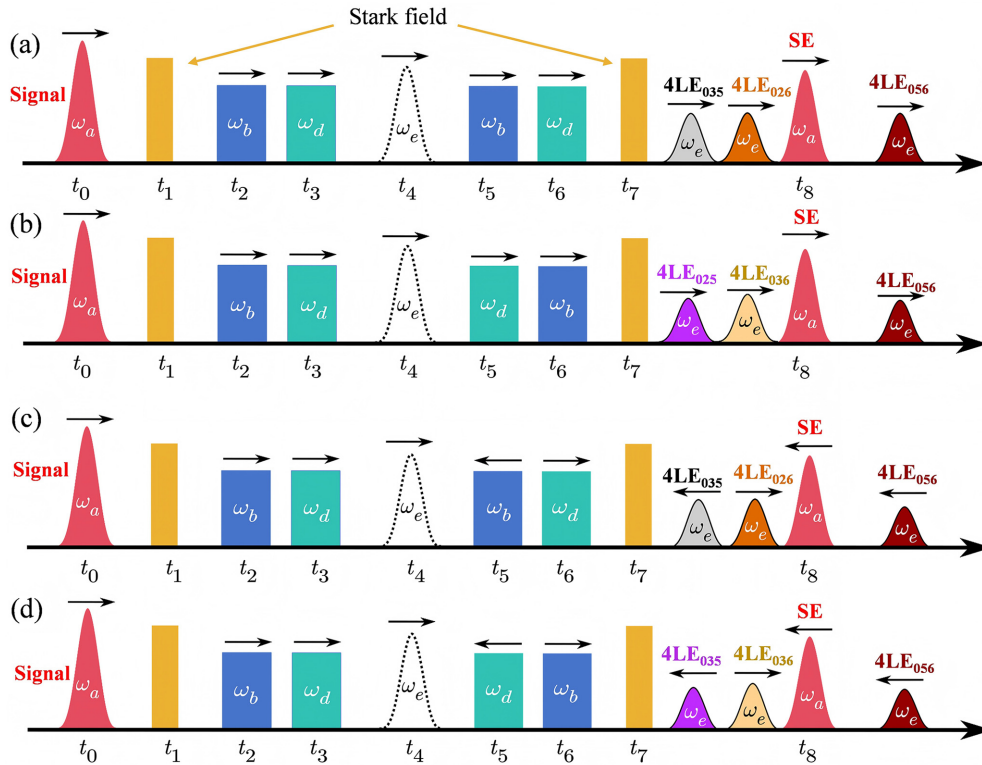


FIG. 10. Noise analysis of the four-level echo noise. (a) and (b) illustrate the noise analysis sequences for forward retrieval, while (c) and (d) illustrate those for backward retrieval.

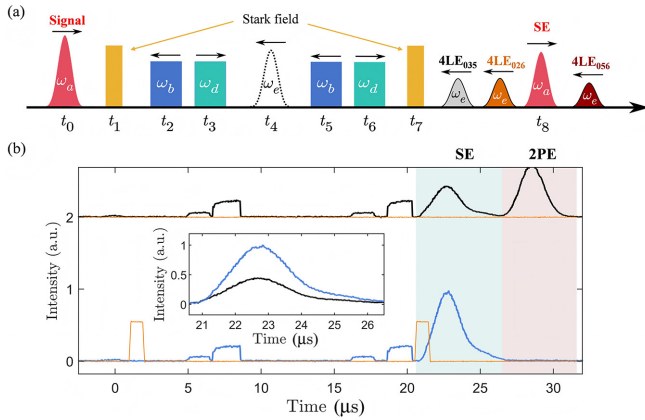


FIG. 11. Complete noiseless sequence. (a) Sequence illustration. (b) Experimental demonstration of the sequence in (a). The blue and black traces represent the APD-detected signals with and without the Stark field (orange traces), respectively. APD was gated on at 20.6 μs by an AOM. The signal within the light green window corresponds to the retrieved echo, and a comparison of the two echo intensities is magnified and shown in the inset. The signal within the light pink window represents the photon-echo (2PE) noise. The upper part of the 2PE curve is not fully displayed due to APD saturation. Except for these two highlighted regions, the signal observed before 20.6 μs originates from the transmission and reflection of the control pulses entering the APD.

APPENDIX D: VALIDATION OF THE PROTOCOL

The quantum memory sequence can be configured as a noise-free scheme, as illustrated in Fig. 11(a). In this

configuration, the π pulses applied at t_2 and t_3 form a pair of counterpropagating pulses; similarly, the π pulses at t_5 and t_6 constitute another counterpropagating pair, with pulses of the same frequency propagating in the same direction. Under these conditions, the SE propagates forward, while all 4LE noise signals propagate backward. Therefore, the detected SE is completely free from 4LE coherent noise.

The experimental result, shown in Fig. 11(b), confirms that the echo signal (blue curve) is perfectly clean, while the 2PE generated by the π -pulse pairs can be effectively suppressed by the applied electric-field pulses (black curve). However, in the free-space configuration, this approach suffers from low η_{pm} due to wave-vector mismatch and the imperfect overlap of counterpropagating beams, which prevents us from achieving high SE efficiency. Hence, this configuration was not used for efficiency calibration in our experiment. In addition, in our forward-propagating configuration without any Stark field applied, the rephasing is theoretically expected to occur at t_4 , corresponding to the 4LE emission. Interestingly, we observed a strong echo emission at t_8 even without the Stark field (not shown). We attribute this echo at t_8 not only to the imperfection of the π pulses, but rather to the reabsorption of the 4LE emitted at t_4 , which subsequently forms a new 4LE sequence with the π pulses at t_5 and t_6 . To verify the validity of our scheme, we configured the pulse sequence as shown in Fig. 11(a), so that the 4LE at t_4 is emitted backward to eliminate reabsorption. We then compared the echo intensity at t_8 with and without the Stark field, as presented in Fig. 11(b). The application of the Stark field significantly enhances the t_8 echo, confirming its crucial role in enabling echo revival, as further illustrated in the inset of Fig. 11(b).

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