

Noise-enabled transparency in optically pumped atomic magnetometers

Xu-Xing Geng^{1,2}, Zhao-Yuan Liu,¹ Ming Xue,¹ Kai Jin,¹ Wang-Wang Tang,¹ Guoqing Yang,³ Shao-Ping Wu,¹ Guang-Ming Huang,^{1,*} and Gao-Xiang Li^{1,†}

¹Department of Physics, *Huazhong Normal University*, Wuhan 430079, China

²School of Mathematics and Statistics, *Huazhong Normal University*, Wuhan 430079, China

³College of Electronics and Information, *Hangzhou Dianzi University*, Hangzhou 310018, China



(Received 20 November 2024; revised 11 February 2026; accepted 2 March 2026; published 25 March 2026)

Non-Gaussian low-frequency noise, such as random telegraph noise (RTN), is widely present in quantum sensing systems and is usually regarded as the main perturbation limiting sensitivity and stability. However, whether it can be transformed into a usable spectroscopic resource still lacks systematic experimental verification. We demonstrate Ξ - and Λ -type noise-induced transparency (NIT) effects in an optically pumped atomic magnetometer coupled to non-Gaussian RTN reservoirs, showing that NIT arises from interference between stochastic transition channels in spins, both theoretically and experimentally. In the Ξ -type configuration, the spectral splitting driven by RTN exhibits a significant change in the transparency window, and the relationship between dynamic criticality and observable spectral structure is elucidated by the correlation between the Liouvillian exceptional point and the resolvable spectral-splitting threshold. In the Λ -type configuration, the interference effect of amplitude-dependent noise enables the transparency effect to be stably manifested under weak noise, providing a calibrable noise-amplitude read-out mechanism. These results extend magnetometer spectra from noise suppression to probing noise with noise, providing a new approach for non-Gaussian noise measurement and adaptive background-noise sensing.

DOI: [10.1103/pbwy-lfk](https://doi.org/10.1103/pbwy-lfk)

I. INTRODUCTION

In quantum sensing, unavoidable coupling to the environment can diminish quantum advantages [1,2]. Especially in the sensing of atoms and solid-state devices based on the magnetic resonance principle, sensing systems are restricted by intrinsic and external noise sources [3–6]. The competition between these noise sources and their internal coherent drivers leads to degradation of sensor capabilities. Therefore, understanding how noise affects quantum sensing is of central importance across a broad range of disciplines [7–13].

Optically pumped atomic magnetometers (OPAMs) [14–19], based on the principle of laser-radio-frequency double-resonance technology, are widely used for detecting various weak magnetic fields [20–23]. In unshielded,

real-world environments, OPAMs are limited not by laboratory settings but by ambient fields: the Earth's field is 30–60 μT [24], and hundreds-of-nT fluctuations render unshielded OPAMs unable to operate stably [25]. Therefore, elucidating and quantitatively characterizing how noise shapes the system response and measurement sensitivity is essential for achieving high-performance OPAMs. More broadly, similar fluctuations are common in solid-state spins [26–31]. Classical low-frequency noise is a key limiter of sensor performance [32–34].

Random telegraph noise (RTN) is a prototypical non-Gaussian fluctuation relevant to many spin-based sensors [35–38]. Moreover, investigating the impact of classical fluctuations on quantum sensors is important for understanding the environmental-response effect in sensing [39–42]. Furthermore, one possibility that has been explored is that quantum systems exposed to a noise reservoir could themselves be potential sources for technological applications [43–47]. This alternative approach transforms noise from a nuisance into a resource.

In this paper, we first propose and demonstrate an open quantum system with noise-induced transparency (NIT) effect in room-temperature cesium atomic vapor based on the OPAM. We reveal that the Ξ -type and Λ -type NIT effects originate from the interference effect of different

*Contact author: gmhuang@ccnu.edu.cn

†Contact author: gaox@ccnu.edu.cn

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI.

spin-moment transition channels enabled by RTN. The NIT effect leads to the appearance of a transparent window in the OPAM spectra, which is significantly affected by the RTN amplitude and switching frequency. We experimentally probe the relationship between the RTN-induced Liouvillian exceptional point (LEP) and the peak-splitting point (PSP). More importantly, this more noise-sensitive Λ -type NIT provides a candidate for noise-traceable measurement and sensor-accuracy calibration. We introduce a frequency-offset auxiliary noise amplitude, correlated and flip-synchronized with the target RTN, to stochastically modulate the detection field, thereby simulating unshielded ambient coupling in a controlled manner and realizing a noise-referenced lock-in response. Not only do we propose a new protocol for measuring noise with noise, but the NIT effect also opens up a new perspective for non-Gaussian open-quantum engineering [48].

The paper is organized as follows. Section II presents the general theoretical model and analytical processing framework for NIT. Sections III and IV report the spectral properties, eigenvalues, and peak-splitting characteristics of RTN-driven OPAM systems in Ξ - and Λ -type NIT, both theoretically and experimentally. Section V analyzes the sensitivity of Ξ - and Λ -type NIT for noise sensing. Section VI concludes and discusses the outlook.

II. GENERAL MODEL FOR NIT

We consider a situation in which dc/ac OPAMs are coupled to RTN. The general Hamiltonian with the NIT effect is

$$H(t) = [\Delta_z + \xi_z(t)]F_z + [\Omega_x + \xi_x(t)]F_x, \quad (1)$$

where the coefficients Δ_z and Ω_x are the detuning and the Rabi frequency of the coupling of the coherent magnetic fields to the angular momenta $F_{z,x}$, respectively. In Eq. (1), setting $\xi_x(t) = 0$ reduces the model to the Ξ -type NIT configuration, in which the transverse coupling is provided solely by the coherent term Ω_x . This configuration can be realized in single- and two-photon OPAM schemes [49,50]. By contrast, setting $\Omega_x = 0$ yields the Λ -type NIT configuration, where the transverse coupling is induced entirely by the noise term $\xi_x(t)$, as realized in the two-photon OPAM scheme [51]. Table I summarizes realistic field scales used in our simulation of unshielded operation; under these correlation conditions (synchronized stochastic modulation), the Λ - scheme yields clear transparency and splitting. For all measurements, the applied noise fields $\xi_{z,x}(t)$ follow an RTN process.

The form of RTN with amplitude α is $\mu(t) = \alpha(-1)^{n(t,0)}$, where $n(t,0)$ is the number of telegraph jumps from time 0 to time t [52]. RTN is a dichotomous Markovian process with a Poissonian distribution of $n(t,0)$, whose correlation function is $\langle \mu(t)\mu(0) \rangle = \alpha^2 \exp(-2\gamma t)$

TABLE I. Representative field magnitudes used in this work.

Quantity	Value (rep.)
Bias field B_0	≈ 290 mG
Larmor frequency $\omega_0/2\pi$	≈ 100 kHz
RTN amplitude $\mu_1(t)$	0.0–12 mG
RTN Rabi frequency $\xi_z/2\pi$	0.0–4.0 kHz
RTN Rabi frequency $\xi_x/2\pi$	≈ 67 Hz
rf Rabi frequency $\Omega_x/2\pi$	≈ 70 Hz
RTN flipping rate γ	0.1–5.0 kHz
Spin relaxation $\Gamma/2\pi$	≈ 46.4 Hz

with telegraph switching rate γ . Meanwhile, the evolution of the joint probability $p(\alpha, t|\alpha_0, t_0)$ satisfies the Chapman-Kolmogorov-Smoluchowski (CKS) equation

$$\dot{p}(\alpha, t|\alpha_0, t_0) = -\gamma p(\alpha, t|\alpha_0, t_0) + \gamma p(-\alpha, t|\alpha_0, t_0), \quad (2)$$

with the initial condition $p(\alpha, t_0|\alpha_0, t_0) = \delta_{\alpha, \alpha_0}$ [53,54]. The physical properties of $H(t)$ are governed by the master equation

$$\dot{\rho} = i[\rho, H(t)] + \mathcal{L}\rho, \quad (3)$$

where ρ is the density matrix and $\mathcal{L}\rho$ describes the relaxation process [55]. Ultimately, by tracing the atomic subsystem and introducing the expected value of the observation, $\mathcal{F}_j = \text{Tr}[F_j \rho]$ ($j = z, \pm$), and combining Eqs. (2) and (3), the dynamic equation for the spin magnetic moment containing the RTN effect can be derived. Solving this dynamic equation allows the complete reconstruction of the spectral response of the OPAM system [56]. A detailed derivation is provided in Appendix A.

To facilitate the correspondence between the above model and the experimental implementation, we briefly describe the physical implementation of each driver in the OPAM platform. The experiment employs the standard circularly polarized light pump-probe technique [14], using circularly polarized light to optically pump and probe atomic spins, with the magnetometer output read out from the transmitted light signal. A weak transverse ac magnetic field is generated by a calibration coil and used to establish deterministic coupling. The RTN field is achieved by injecting our homemade RTN source into the corresponding coil current and converting it into magnetic field units for independent calibration. Details on device setup, calibration, and data processing are provided in Appendix C.

III. THE Ξ -TYPE NIT IN A dc OPAM

A. Hamiltonian of Ξ -type NIT

In this section, we study the Ξ -type NIT configuration in a dc OPAM [49]. The system operates under a static bias

field and a weak transverse ac drive, while the longitudinal field component is stochastically modulated by an RTN field. The equivalent driving magnetic field is

$$\hat{B}(t) = [B_0 + \mu_1(t)]\hat{e}_z + B_2 \cos(\omega_2 t)\hat{e}_x, \quad (4)$$

where B_0 and B_2 are the amplitudes of the dc and ac components of the magnetic field, respectively. The parameter ω_2 is the frequency of the ac field, and $\mu_1(t)$ is an RTN. Based on the field-atomic interaction theory [57] and the angular momentum theory [58], and under the rotating-wave approximation, its effective Hamiltonian in a rotating frame with frequency ω_2 is

$$H_{\text{dc}} = [\Delta_z + \xi_z(t)]F_z + \Omega_x F_x, \quad (5)$$

where the detuning is defined as $\Delta_z = \Omega_0 - \omega_2$. The Rabi frequencies are $\Omega_0 = \mu_B g_F B_0$, $\Omega_x = \mu_B g_F B_2/2$, while the noise-induced coupling is $\xi_z(t) = \mu_B g_F \mu_1(t)$. In these expressions, μ_B denotes the Bohr magneton and g_F is the Landé g factor. This corresponds to the case where $\xi_x(t) = 0$ in Eq. (1). The model is characterized by using a coherent driving field to detect noise. The transverse coherent field Ω_x provides a stable probe that converts the longitudinal random modulation $\xi_z(t)$ into a readable spin response change. Without $\xi_z(t)$, the model is a standard OPAM based on laser-radio-frequency double-resonance technology [49].

B. Spin-moment evolution of Ξ -type NIT

To gain further insights, substituting the Hamiltonian (5) into Eqs. (2) and (3) yields the evolution of the system as

follows:

$$\dot{\mathcal{F}}_{z,S} = -\Gamma(\mathcal{F}_{z,S} - \mathcal{F}_{z,\text{eq}}) + i\Omega_x(\mathcal{F}_{+,S} - \mathcal{F}_{-,S}), \quad (6a)$$

$$\dot{\mathcal{F}}_{\pm,S} = -(\Gamma \pm i\Delta_z)\mathcal{F}_{\pm,S} \mp i\xi_z\mathcal{F}_{\pm,A} \pm i\Omega_x\mathcal{F}_{z,S}, \quad (6b)$$

$$\dot{\mathcal{F}}_{\pm,A} = -(\gamma' \pm i\Delta_z)\mathcal{F}_{\pm,A} \mp i\xi_z\mathcal{F}_{\pm,S} \pm i\Omega_x\mathcal{F}_{z,A}, \quad (6c)$$

with $\gamma' = 2\gamma + \Gamma$. The moments $\mathcal{F}_{j,S(A)} = \mathcal{F}_{j,\alpha} \pm \mathcal{F}_{j,-\alpha}$ are the symmetric and antisymmetric superpositions of spin moments $\mathcal{F}_{j,\pm\alpha}$ ($j = z, \pm$) tied to the joint probability $p(\pm\alpha, t|\alpha_0, t_0)$ [52]. The moment $\mathcal{F}_{z,\text{eq}}$ is the thermal-equilibrium spin moment [59]. The parameter Γ ($\Gamma \ll \gamma$) is the isotropic spin-moment relaxation rate. Observing $\mathcal{F}_{z,S}$, Eqs. (6b) and (6c) predict a Ξ -type EIT-like response [60]; see Fig. 1(a). Here ξ_z acts as the coupling drive (addressing $\mathcal{F}_{\pm,A}$ and $\mathcal{F}_{\pm,S}$), while Ω_x probes the link between $\mathcal{F}_{z,S}$ and $\mathcal{F}_{\pm,S}$. RTN and the rf field are coupled to the spin moments, leading to interference of two different quantumlike routes—directly via the $\mathcal{F}_{z,S} \leftrightarrow \mathcal{F}_{\pm,S}$ pathway, or indirectly via $\mathcal{F}_{z,S} \leftrightarrow \mathcal{F}_{\pm,S} \leftrightarrow \mathcal{F}_{\pm,A} \leftrightarrow \mathcal{F}_{z,A}$ pathway (or by higher-order variants) [61,62]. When the two paths match the resonance conditions, destructive interference occurs, inducing transparency. This is our so-called NIT effect. We provide a detailed comparison between Ξ -type NIT and Ξ -type EIT in Appendix A 1.

C. Analytical spectra of Ξ -type NIT

Experimentally, the spectrum (circularly polarized pump-probe) is proportional to $\mathcal{P}_z \propto \mathcal{F}_{z,S}$ [63–65]. In the

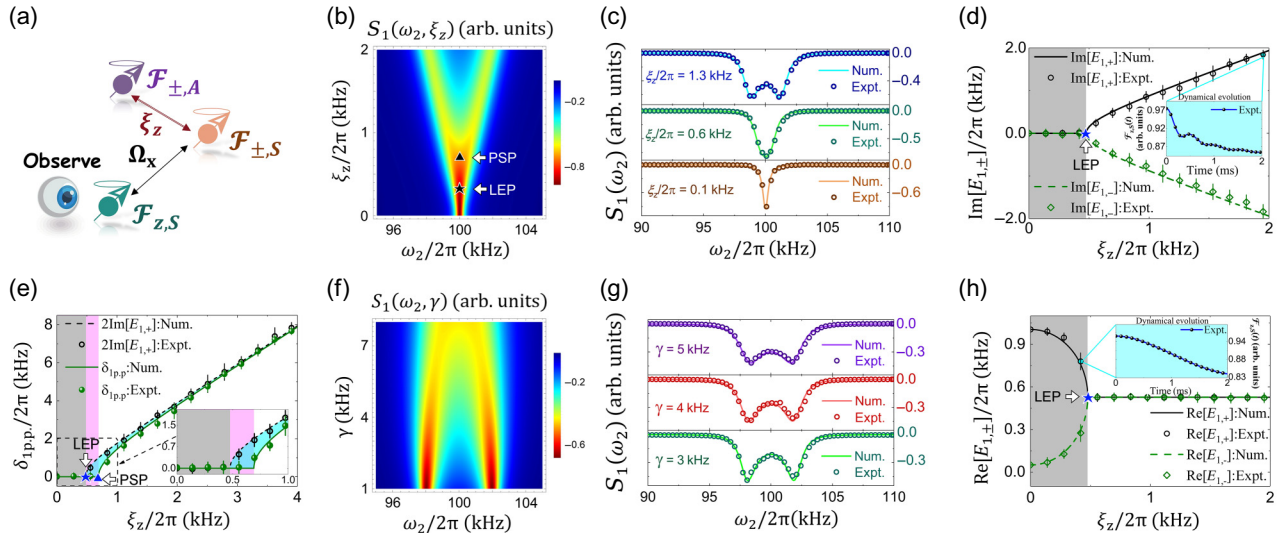


FIG. 1. (a) Ξ -type NIT spin-moment coupling channels for Eqs. (6b) and (6c). (b) Numerical simulation of the $S_1(\omega_2, \xi_z)$ map on the ω_2 - ξ_z plane at $\gamma = 3$ kHz. (c) Spectra (experiment and simulation) at fixed $\gamma = 3$ kHz for $\xi_z/2\pi = 1.3, 0.6$, and 0.1 kHz. (d),(h) Real and imaginary parts of the eigenvalues versus ξ_z ; insets: exact- and broken-phase dynamics. (e) The influence of ξ_z on $\delta_{1p,p}$ and $2\text{Im}[E_{1,+}]$ at $\gamma = 3$ kHz. (f) Numerical simulation of the $S_1(\omega_2, \gamma)$ map on the ω_2 - γ plane at $\xi_z/2\pi \approx 1.9$ kHz. (g) Spectra (experiment and simulation) at fixed $\xi_z/2\pi \approx 1.9$ kHz for $\gamma = 5, 4$, and 3 kHz. Other parameters are given in Table I.

steady state ($\dot{\mathcal{F}}_{j,S(A)} \equiv 0, j = z, \pm$), the OPAM spectrum is

$$\mathcal{S}_1(\Delta_z) = \mathcal{F}_{z,\text{eq}} \left(1 - \frac{2\Omega_x^2}{\Gamma Z_1} \left\{ \Delta_z^2 (\Gamma + \gamma') - \gamma' \left[\Delta_z^2 - \Gamma \gamma' - \xi_z^2 \left(1 + \frac{2\Omega_x^2 \Gamma}{\xi_z^2 \gamma'} \right) \right] \right\} \right) \quad (7)$$

with

$$Z_1 = [\Delta_z^2 + c_2^2 (2\gamma^2 + \Gamma \gamma')]^2 + 4(\Gamma + \gamma)^2 \times [(c_1 \xi_z)^2 - \gamma^2], \quad (8)$$

where the coefficients c_1 and c_2 are $c_1 = \sqrt{1 + \epsilon_1}$ and $c_2 = \sqrt{1 + \epsilon_2}$, with saturation parameters $\epsilon_1 = 2\Omega_x^2/(\Gamma \gamma')$ and $\epsilon_2 = (2\Omega_x^2 - \xi_z^2)/(2\gamma^2 + \Gamma \gamma')$. Although Eqs. (7) and (8) give a compact spectral form, the physical picture reflected by this form is not intuitive and does not clearly reveal the physical mechanism of NIT.

To obtain a more interpretable physical picture, a natural approach is to define the detection process as a linear response. In experiments, the probe field is usually much weaker than the pump and noise drive, and its role is mainly to read out rather than reshape the system dynamics. Therefore, we further consider the weak-detection limit $\Omega_x \rightarrow 0$. Under this limit, the detection field excites only the transverse coherent components with first-order perturbations. The observed $\mathcal{S}_1(\Delta_z)$ can be equivalently understood as the linear polarizability response of transverse coherence to detuning Δ_z . The transverse moments $\mathcal{F}_{+,S}$ and $\mathcal{F}_{+,A}$ in Eqs. (6a)–(6c) will be decoupled from the remaining spin moments and form a closed two-dimensional subspace

$$\begin{pmatrix} \dot{\mathcal{F}}_{+,S} \\ \dot{\mathcal{F}}_{+,A} \end{pmatrix} = \mathcal{L}_1 \cdot \begin{pmatrix} \mathcal{F}_{+,S} \\ \mathcal{F}_{+,A} \end{pmatrix}, \quad (9)$$

with

$$\mathcal{L}_1 = \underbrace{\begin{pmatrix} -\Gamma & -i\xi_z \\ -i\xi_z & -2\gamma - \Gamma \end{pmatrix}}_{\mathcal{L}_{1,s}} + \underbrace{\begin{pmatrix} -i\Delta_z & 0 \\ 0 & -i\Delta_z \end{pmatrix}}_{\mathcal{L}_{1,p}}, \quad (10)$$

where $\mathcal{L}_{1,s}$ is the Liouvillian superoperator of the system [66,67]. The Liouvillian superoperator $\mathcal{L}_{1,p}$ is related to the probe field. By solving the determinant $\det(\mathcal{L}_{1,s} - E_1 \mathcal{I}) = 0$, the eigenvalues of the system can be obtained, where $\mathcal{I}$ is the 2×2 identity matrix. The eigenvalues of $\mathcal{L}_{1,s}$ are

$$E_{1,\pm} = \gamma + \Gamma \pm \sqrt{\gamma^2 - \xi_z^2}. \quad (11)$$

As can be seen from the eigenmodes of this two-dimensional subspace, the modal properties are controlled by the discriminant $\gamma^2 - \xi_z^2$. For $\xi_z < \gamma$, the system exhibits overdamped relaxation (real eigenvalues). For $\xi_z > \gamma$, it exhibits underdamped oscillation (complex-conjugate eigenvalues). The critical point $\xi_z = \gamma$ corresponds to the LEP. Therefore, the partition boundary of the spectrum in Eq. (7) is naturally defined. Based on this boundary, we introduce partition parameters

$$\Pi = \begin{cases} \sqrt{\gamma^2 - \xi_z^2}, & \xi_z < \gamma, \\ 0, & \xi_z = \gamma, \\ \sqrt{\xi_z^2 - \gamma^2}, & \xi_z > \gamma. \end{cases} \quad (12)$$

Under this partition, the spectrum in Eq. (7) is decomposed into a more physical form as a superposition of multiple Lorentzian spectra

$$\mathcal{S}_1(\Delta_z) = \begin{cases} \mathcal{F}_{z,\text{eq}} \left\{ 1 - \frac{\Omega_x^2}{4\Gamma\Pi} \left[\frac{(\Gamma + 2\Pi)^2 - \gamma'^2}{\Delta_z^2 + (\gamma + \Gamma + \Pi)^2} - \frac{(\Gamma - 2\Pi)^2 - \gamma'^2}{\Delta_z^2 + (\gamma + \Gamma - \Pi)^2} \right] \right\}, & \xi_z < \gamma, \\ \mathcal{F}_{z,\text{eq}} \left\{ 1 - \frac{2\Omega_x^2}{\Gamma} \frac{\Delta_z^2 \Gamma + (\gamma + \Gamma)^2 (2\gamma + \Gamma)}{[\Delta_z^2 + (\gamma + \Gamma)^2]^2} \right\}, & \xi_z = \gamma, \\ \mathcal{F}_{z,\text{eq}} \left(1 - \frac{\Omega_x^2}{\Gamma} \left[\frac{\gamma'}{(\Delta_z - \Pi)^2 + (\gamma + \Gamma)^2} + \frac{\gamma'}{(\Delta_z + \Pi)^2 + (\gamma + \Gamma)^2} - \frac{4\Delta_z^2 \gamma}{[(\Delta_z - \Pi)^2 + (\gamma + \Gamma)^2][(\Delta_z + \Pi)^2 + (\gamma + \Gamma)^2]} \right] \right), & \xi_z > \gamma. \end{cases} \quad (13)$$

Compared with the conventional OPAM response, which is well approximated by a single Lorentzian line-shape [49,63], the Ξ -type NIT spectrum (13) exhibits richer and more tunable spectral characteristics, including transparent windows induced by modal competition and interference, as well as structural changes such as splitting and reconstruction. Based on the theoretical prediction of Eq. (13), we then experimentally verify the Ξ -type NIT response. The next section presents the measured spectra and provides a quantitative analysis of the associated transparency window, splitting, and reconstruction.

D. Experimental results and analysis of Ξ -type NIT

We next examine the Ξ -type NIT spectra under RTN and quantify their dependence on the noise amplitude and switching rate. To expose the physical origin of the observed features, We further investigate the relationship and differences between the spectral splitting and eigenvalues observed in the experiment.

1. Spectrum versus RTN amplitude and switching rate

For $\xi_z < \gamma$, Eq. (13) shows that the NIT spectrum is determined by the superposition of two Lorentzian components with opposite signs, both centered at $\Delta_z = 0$ but with different effective linewidths $\gamma + \Gamma \pm \Pi$. Physically, this subtraction corresponds to destructive interference near the center of the spectrum, thus naturally producing a tendency for transparency at $\Delta_z = 0$. This tendency is more readily apparent in two-dimensional parametric plots, as shown in Fig. 1(b). As ξ_z increases within a small range and gradually approaches γ , the decrease in $\Pi = \sqrt{\gamma^2 - \xi_z^2}$ enhances the difference effect, and the central suppression deepens accordingly, showing a trend of gradually excavating transparent features. The experimental results in Fig. 1(c) also demonstrate that the spectrum retains the quasi-single-Lorentzian form at $\xi_z/2\pi = 0.1$ kHz, without the appearance of a transparent window.

Conversely, for $\xi_z > \gamma$, Eq. (13) shows that the spectrum is dominated by two Lorentzian components that are symmetrically shifted to $\Delta_z = \pm\Pi$, where $\Pi = \sqrt{\xi_z^2 - \gamma^2}$ directly sets the splitting scale. Additional interference terms further determine the indentation between the two side structures and the spectral reconstruction of the central region. Physically, this means that after crossing $\xi_z = \gamma$, the spectral weights shift from central anchoring to symmetric splitting. As ξ_z increases, Π increases monotonically, and the splitting distance 2Π gradually widens. Therefore, as shown in Fig. 1(b), the evolutionary trajectory of the central structure branching to both sides and forming a symmetrical enhancement near $\Delta_z = \pm\Pi$ can be clearly traced. The curve at the top of Fig. 1(c) shows a clear central transparent window that is experimentally measured at $\xi_z/2\pi \approx 1.3$ kHz. This result clearly shows that, when $\xi_z \gg \gamma$, RTN is no longer merely a spectral

correction but instead induces a pronounced transparent window.

The switching rate γ of telegraph noise directly modulates the visibility of NIT. At smaller γ values, the transparent window is clearer, whereas as γ increases, the transparent window gradually becomes lighter and tends to close. Figure 1(f) presents a global view of $\mathcal{S}_1(\omega_2, \gamma)$ at a fixed $\xi_z/2\pi \approx 1.9$ kHz. In the small- γ region, a transparent window around $\omega_2 \approx 100$ kHz is visible, coexisting with the symmetrical response branches on both sides. As γ increases, the two branches converge toward the center and merge, while the transparent window is continuously filled at the same time, resulting in a significant reduction in contrast. The experimental curves in Fig. 1(g) further confirm that, as γ increases, the transparent window is gradually filled. The window becomes significantly shallower from 3 to 5 kHz, and the spectral lines correspondingly revert to a single-envelope morphology. Experimental and numerical results show consistent trends toward window closure, indicating that γ is the key parameter regulating the transition of NIT from a distinguishable window to an average, flattened state.

2. Eigenvalue signatures and spectral splitting

The origin of spectral reconstruction under ξ_z scanning can be more clearly understood from the perspective of system eigenvalues $E_{1,\pm}$. Figure 1(d) shows that $\text{Im}[E_{1,\pm}]$ is approximately zero below the threshold ($\xi_z = \gamma$), while the imaginary part splits into a pair of finite branches with opposite signs after crossing the threshold and expands as ξ_z increases. This provides a direct scale for the resolvable characteristic frequencies in the spectrum and corresponds to the trend of spectral weights expanding from center-locked to features on both sides. Complementing this, in Fig. 1(h), $\text{Re}[E_{1,\pm}]$ converges and merges near the threshold. This indicates that the two decay timescales merge into a common broadened background, thus determining the contrast and resolvability of the transparent window or split structure in the spectrum. Furthermore, this threshold also aligns with the statistical limit of noise. Rapid flipping below the threshold causes the noise to become Markovian and induces dynamic averaging (motional narrowing), thereby weakening the transparency window. The increased correlation time above the threshold and the enhanced effective non-Markovianity allow the underdamped modes to appear in the spectrum as transparent windows or split features [68–71].

The dynamical splitting scale $2\text{Im}[E_{1,+}]$ is compared side by side with the peak-splitting point $\delta_{\text{p.p.}}$ of the steady-state spectrum; see Fig. 1(e). Interestingly, the dynamical threshold (LEP) does not coincide with the spectroscopic splitting threshold (PSP). As ξ_z increases, $2\text{Im}[E_{1,+}]$ is almost zero before the LEP and immediately

becomes nonzero after crossing it, indicating that the Liou-villian eigenmode has generated finite oscillations. However, $\delta_{\text{lp.p.}}$ only jumps significantly from the near-zero state until a larger threshold PSP ($\xi_z \approx \sqrt{\gamma'^3/(2\gamma' + \Gamma)}$), indicating that, under finite linewidth and contrast constraints, the dynamical oscillation does not necessarily immediately manifest as a resolvable double peak in the steady-state spectrum. After entering the resolvable region, $\delta_{\text{lp.p.}}$ grows rapidly and shows an evolutionary trend consistent with $2\text{Im}[E_{1,+}]$, approaching $\delta_{\text{lp.p.}} \approx 2\xi_z$ in the large- ξ_z limit. In this view, the LEP provides a dynamical scale for noise-induced spectral reconstruction, while the PSP marks the threshold at which the noise-transparency window reaches a resolvable value in the steady-state spectrum.

Similar two-photon magnetometer models also exhibit the consistent Ξ -type NIT effect. The relevant results are summarized in Appendix D. To avoid repetition, they are not discussed further here. To achieve higher-contrast NIT, we further propose a Λ -type NIT via amplitude-correlated RTN.

IV. THE Λ -TYPE NIT IN AN ac OPAM

A. Hamiltonian of Λ -type NIT

In this section, we study the Λ -type NIT configuration in an ac OPAM [18,25,51]. The system is driven by a transverse ac magnetic field, while along the longitudinal axis it experiences a bias field and two amplitude-correlated RTN fields, one of which includes a frequency offset. The resulting effective driving field is

$$\hat{B}(t) = [B_0 + \mu_1(t) + \mu_2(t) \cos(\omega_1 t)]\hat{e}_z + B_2 \cos(\omega_2 t)\hat{e}_x, \quad (14)$$

where $\mu_k(t)$ ($k = 1, 2$) are the amplitude-correlated RTN sources. This can be achieved by modulating the noise amplitude. In the rotating frame, the Hamiltonian of the system is

$$H_{\text{ac}} = [\Omega_0 - \omega_2 + \xi_1(t) + \xi_2(t) \cos(\omega_1 t)]F_z + \frac{\Omega_2}{2}F_x, \quad (15)$$

where the coupling parameters $\xi_k(t)$ and Ω_j are defined as $\xi_k = \mu_B g_F \mu_k(t)$ ($k = 1, 2$) and $\Omega_j = \mu_B g_F B_j$ ($j = 0, 2$), respectively. The oscillation frequencies of the two ac fields satisfy the cascade two-photon resonance relationship $\omega_1 + \omega_2 \approx \Omega_0$. If the noise term $\xi_1(t)$ is ignored in Eq. (15) and $\xi_2(t)$ is taken as a constant driver, the model reduces to the standard cascaded two-photon OPAM configuration [50]. To provide an intuitive picture of how this noise impacts the OPAM, Hamiltonian (15) is incorporated into the eigenrepresentation of $H' = \Omega_0 F_0 + \Omega_2 F_x/2$. In

this representation, the Hamiltonian of the system is

$$H'_{\text{ac}} = [\Delta_z + \xi_z(t)]F_z + \xi_x(t)F_x. \quad (16)$$

In this case, the parameters for Eq. (16) are $\Delta_z = \lambda - \Omega_0 + \omega_2$, $\xi_z(t) = \omega_1 \xi_1(t)/\lambda$, and $\xi_x(t) = -\xi_2(t)\Omega_2/(2\lambda)$, with $\lambda = \sqrt{4\omega_1^2 + \Omega_2^2}/2$, where $\xi_x(t)$ is a noisy detection field that is amplitude correlated with the target RTN $\mu_1(t)$ and synchronized in its flipping dynamics. In the parameter regime considered here, the hierarchy $\Omega_0 > \omega_1, \omega_2 \gg |\xi_x(t)|$ holds. This model corresponds to the case where $\Omega_z = 0$ in Eq. (1). The key feature of this configuration is that the effective detection field $\xi_x(t)$ is random. Specifically, the effective detection $\xi_x(t)$ is not an externally added independent perturbation but rather the correlated noise associated with the target RTN contained in $\xi_z(t)$. It is this intrinsic correlation that transforms noise from interference into a resource, enabling noise-on-noise probing and allowing it to be directly tracked within the measured OPAM signal.

B. Spin-moment evolution of Λ -type NIT

To reveal the physics of the Λ -type NIT, we write the evolution equations

$$\dot{\mathcal{F}}_{z,S} = -\Gamma(\mathcal{F}_{z,S} - \mathcal{F}_{z,\text{eq}}) + i\xi_x(\mathcal{F}_{+,A} - \mathcal{F}_{-,A}), \quad (17a)$$

$$\dot{\mathcal{F}}_{\pm,S} = -(\Gamma \pm i\Delta_z)\mathcal{F}_{\pm,S} \mp i\xi_z\mathcal{F}_{\pm,A} \pm i\xi_x\mathcal{F}_{z,A}, \quad (17b)$$

$$\dot{\mathcal{F}}_{\pm,A} = -(\gamma' \pm i\Delta_z)\mathcal{F}_{\pm,A} \mp i\xi_z\mathcal{F}_{\pm,S} \pm i\xi_x\mathcal{F}_{z,S}. \quad (17c)$$

It is worth noting that the above equations are almost completely isomorphic to the dynamic equations of the standard Λ -type EIT [72,73]. Moreover, the coupling channel, interference path, and key factors that determine the transparent window correspond one-to-one with those of traditional Λ -type EIT. The difference lies in the fact that both fields involved in the interference are realized by noise processes in this model. The two noise paths enter their respective channels through randomly flipped equivalent Zeeman modulation and still generate EIT-type destructive interference without a coherent phase reference, thereby opening the transparent window.

In this configuration, the detection field ξ_x serves as the probe that drives $\mathcal{F}_{\pm,A}$ and $\mathcal{F}_{z,S}$, while the correlated RTN ξ_z acts as the coupling field linking $\mathcal{F}_{\pm,A}$ and $\mathcal{F}_{\pm,S}$; see Fig. 2(a). Correspondingly, two amplitude-correlated RTN channels then generate interfering pathways: (i) $\mathcal{F}_{z,S} \leftrightarrow \mathcal{F}_{\pm,A}$ and (ii) $\mathcal{F}_{z,S} \leftrightarrow \mathcal{F}_{\pm,A} \leftrightarrow \mathcal{F}_{\pm,S} \leftrightarrow \mathcal{F}_{\pm,A}$ (or higher-order variants) [61,62]. Interference between these pathways opens a transparency window—the Λ -type NIT. We exploit this correlation as a noise-referenced control: a stochastic envelope synchronized with the RTN ξ_x . The two noise-modulated pathways in Eqs. (17b) and (17c)

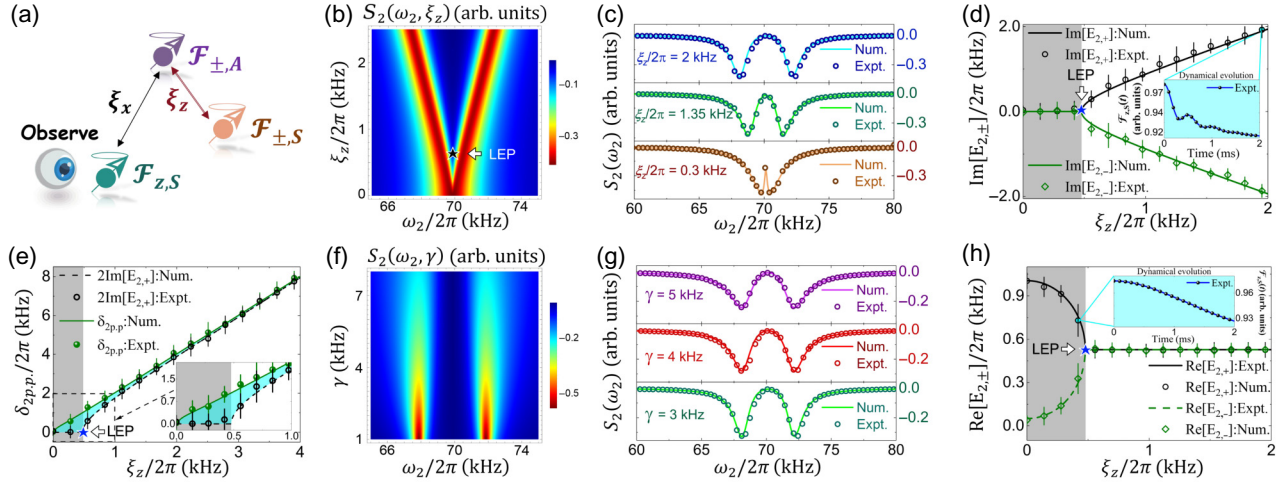


FIG. 2. (a) Λ -type NIT spin-moment coupling channels for Eqs. (17b) and (17c). (b) Numerical simulation of the $S_2(\omega_2, \xi_z)$ spectral map on the ω_2 - ξ_z plane at $\gamma = 3$ kHz. (c) Spectra (experiment and simulation) at fixed $\gamma = 3$ kHz for $\xi_z/2\pi = 2.0, 1.35$, and 0.3 kHz. (d),(h) Real and imaginary parts of the eigenvalues versus ξ_z ; insets: exact- and broken-phase dynamics. (e) Influence of ξ_z on $\delta_{2p,p}$ and $2\text{Im}[E_{2,+}]$ at $\gamma = 3$ kHz. (f) Numerical simulation of the $S_2(\omega_2, \gamma)$ map on the ω_2 - γ plane at $\xi_z/2\pi \approx 2.0$ kHz. (g) Spectra (experiment and simulation) at fixed $\xi_z/2\pi \approx 2.0$ kHz for $\gamma = 5, 4$, and 3 kHz. Other parameters are listed in Table I.

then interfere, opening a high-contrast transparency window (Fig. 2). We provide a detailed comparison between Λ -type NIT and Λ -type EIT in Appendix A 2.

C. Analytical spectra of Λ -type NIT

To gain further insight, using Eqs. (17a)–(17c), the steady-state spectrum is

$$S_2(\Delta_z) = F_{z,\text{eq}} \left(1 - \frac{2\xi_x^2}{\Gamma Z_2} \left\{ \Delta_z^2 (\Gamma + \gamma') - \Gamma \left[\Delta_z^2 - \Gamma \gamma' - \xi_z^2 \left(1 + \frac{2\xi_x^2}{\xi_z^2} \right) \right] \right\} \right), \quad (18)$$

with

$$Z_2 = [\Delta_z^2 + s_1^2 (2\gamma^2 + \Gamma\gamma') - \xi_z^2]^2 + 4s_1^2 (\Gamma + \gamma)^2 \times [\xi_z^2 - (\gamma s_1)^2] \quad (19)$$

and $s_1 = \sqrt{1 + \beta}$. The parameter $\beta = 2\xi_x^2/(\Gamma\gamma')$ represents the RTN saturation effect. In the weak-detection limit ($\xi_x \rightarrow 0, \beta = 0$), the spectrum (18) can also be decomposed into segments at the LEP ($\xi_z = \gamma$) to obtain a structure similar to that of the Ξ -type NIT spectrum (13). For related details, we refer the reader to Appendix B. Interestingly, Eq. (18) is almost isomorphic to the classical spectrum of Λ -type EIT in terms of analytical structure [72], so the spectroscopic laws it predicts should be consistent with the core physical conclusions of Λ -type EIT.

D. Experimental results and analysis of Λ -type NIT

Below, we present the experimental results of Λ -type NIT under the action of RTN and directly compare them with those of the aforementioned Ξ type. We employ consistent spectral characterization measures to ensure that the observed similarities and differences primarily stem from the configurational differences themselves. This further reveals the distinction between the Ξ - and Λ -type NIT effects.

1. Spectrum versus RTN amplitude and switching rate

Figure 2(b) shows the spectral line-shape changes of the Λ -type NIT under a ξ_z scan. Near $\omega_2 \approx 70$ kHz, the central transparent window remains stable throughout the scan range. The spectral center does not drift with ξ_z , while symmetrically enhanced response branches appear on both sides, forming a clear bifurcation profile. It is worth noting that, unlike the Ξ -type NIT spectrum, the transparency window of this Λ configuration can still be directly identified even in the region where $\xi_z < \gamma$. The key reason is that the two RTN fields in this Λ configuration are correlated in amplitude, thereby preserving a stable interference structure in the readout channel. In other words, although the drive comes from random flipping, the amplitude correlation causes the relative weights of the two coupled paths to be modulated synchronously, and the transparency condition is not easily smoothed out by random averaging. As a result, the window can maintain limited contrast even in the overdamped regime. Consistent with this behavior, Fig. 2(c) shows that, as $\xi_z/2\pi$ increases from 0.3 to 1.35 and 2.0 kHz, the central transparent window gradually deepens and widens, and the responses on both sides

increase synchronously. The experimental data agree well with the numerical curves. This feature provides a route to tunable, high-contrast noise sensing under low-noise-amplitude or overdamped operating conditions.

With a fixed $\xi_z/2\pi \approx 2.0$ kHz, Fig. 2(f) shows the variation of the Λ -type NIT spectrum with the switching rate γ . It can be seen that increasing γ mainly changes the width of the transparent window and the spectral intensity. As γ increases from 3 to 5 kHz, the transparent window around $\omega_2 \approx 70$ kHz shrinks significantly along the frequency axis, while the symmetrical branches on both sides gradually weaken, making the bifurcation profile of the spectrum more convergent. The experimental data in Fig. 2(g) provide consistent evidence that, as γ increases, the width of the central depression decreases, while the amplitudes of the two side peaks decrease synchronously and tend to become rounded. Physically, an increase in γ means that noise flips more frequently, and the system responds as if it were driven by a rapidly averaged field. Therefore, the transparent structure is maintained only within a narrower range near the center frequency, and the window becomes correspondingly narrower, while the peak values on both sides are thinned and reduced. Overall, γ provides a control knob for continuously adjusting the bandwidth and visibility of the transparent window at a nearly fixed center frequency.

2. Eigenvalue signatures and spectral splitting

It is worth emphasizing that the eigenvalue evolution shown in Figs. 2(d) and 2(h) has the same critical structure as the Ξ -type NIT configuration. Specifically, in Fig. 2(d), $\text{Im}[E_{2,\pm}]$ is approximately zero to the left of $\xi_z = \gamma$, and after crossing this point, it splits into a pair of branches with opposite signs and expands with ξ_z , corresponding to the appearance of the oscillation-frequency scale; see Fig. 2(d). Complementing this, Fig. 2(h) shows that $\text{Re}[E_{2,\pm}]$ merges and then separates near the critical point, indicating that the relaxation timescales merge at this point and recombine on both sides of the threshold. The simultaneous occurrence of the abovementioned splitting of the imaginary part and merging of the real part corresponds one-to-one with the conclusions in the Ξ -type diagram, indicating that the LEP at $\xi_z = \gamma$ is an eigenspectral marker shared by the two configurations.

To contrast an intrinsic dynamical scale with a directly observable spectroscopic metric, we also present the eigenvalue-based oscillation scale $2\text{Im}[E_{2,+}]$ for this configuration and the variation of the peak-to-peak splitting $\delta_{2p,p}$ directly read from the spectral peak position with the telegraph-noise amplitude ξ_z ; see Fig. 2(e). First, $\delta_{2p,p}$ remains resolvable almost from the weak-noise end in this configuration and grows approximately linearly with ξ_z , which means that the double structure induced by transparency can be directly used as a spectroscopic scale for

the noise amplitude. This behavior is distinctly different from that of Ξ -type NIT, whose spectral transparency window typically becomes clearly visible only after crossing the LEP, whereas the Λ -type configuration exhibits a traceable transparency trend in the interval $\xi_z < \gamma$.

In contrast, the threshold behavior of $2\text{Im}[E_{2,+}]$ is closer to the general eigenvalue reconstruction. The imaginary part is zero when $\xi_z < \gamma$, corresponding to the overdamped mode. When ξ_z exceeds the LEP, the imaginary part splits in pairs and grows with ξ_z , which is consistent with the eigenvalue evolution trend of the Ξ -type NIT configuration. Therefore, it is important to emphasize that there is no one-to-one correspondence between imaginary-eigenvalue splitting and spectral splitting. The former represents the intrinsic oscillation scale of the dynamic mode, while the latter also depends on the projection weight of the read-out channel, the broadening of the line width, and the reshaping of the line shape by interference.

In the Λ -type NIT configuration, it is the amplitude correlation of the two RTN sources that stabilizes the interference weights, so that even at small ξ_z , the transparency effect can appear in advance in the form of distinguishable spectral splitting. Its linear properties are highly isomorphic to the bimodal transparency structure of standard Λ -type EIT. Furthermore, Fig. 2(e) shows that $\delta_{2p,p}$ responds more directly to ξ_z and enters the linear region earlier, demonstrating higher noise sensitivity and a wider usable dynamic range. Under high noise, the two curves gradually tend to overlap and approach $2\xi_z$, indicating that, when noise dominates the splitting scale, the spectral peak splitting and the dynamical oscillation scale become locked to the same amplitude scale. These results enable traceable non-Gaussian-noise metrology [74] and background-noise self-adaptive sensors [75].

V. SENSITIVITY ANALYSIS

We compare and analyze the sensitivity performance of dc and ac magnetometers in background magnetic noise and RTN-field measurements; see Fig. 3. All spectra are obtained by continuous open-loop acquisition for 2 h to ensure sufficient resolution in the low-frequency band. Without external RTN sources, Figs. 3(a) and 3(b) show the background-noise spectra of the dc and ac magnetometers, respectively, with corresponding magnetic field sensitivity noise floors of approximately 28.7 and 16.3 pT/ $\sqrt{\text{Hz}}$. Both show an upward trend at low frequencies, reflecting the unavoidable $1/f$ drift and slow-varying disturbances in the system, such as laser power and frequency fluctuations, photoelectric detection noise, atomic-cell temperature drift, and residual environmental fluctuations in magnetic shielding. In contrast, ac magnetometers modulate the signal to shift it out of the low-frequency drift-dominant region, resulting in a lower noise floor across the entire frequency range.

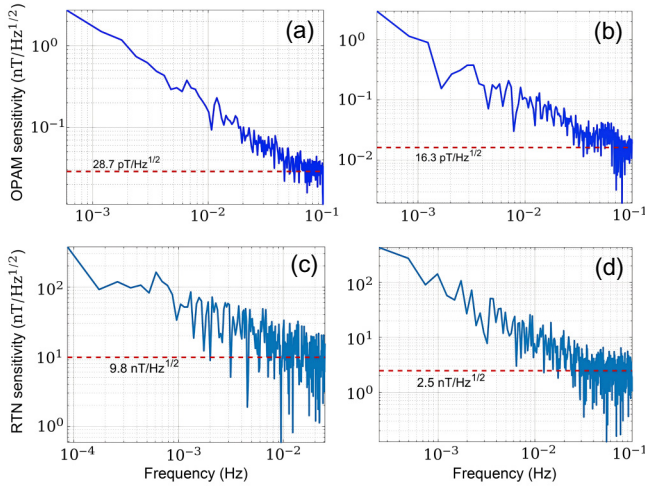


FIG. 3. Sensitivity of dc and ac magnetometers to measure magnetic and RTN fields. These sensitivities are obtained by collecting magnetometer signals for 2 h in each measurement. (a) The sensitivity of the dc magnetometer for measuring the magnetic field is about $28.7 \text{ pT}/\sqrt{\text{Hz}}$. (b) The sensitivity of the ac magnetometer for measuring an ac magnetic field is about $16.3 \text{ pT}/\sqrt{\text{Hz}}$. (c) The sensitivity of the dc magnetometer for measuring the RTN field is about $9.8 \text{ nT}/\sqrt{\text{Hz}}$. (d) The sensitivity of the ac magnetometer for measuring the RTN field is about $2.5 \text{ nT}/\sqrt{\text{Hz}}$. In both cases, the RTN sensitivity is measured at $\xi_z \approx 1 \text{ kHz}$. Both the vertical and horizontal axes are plotted on logarithmic scales.

Furthermore, under the condition $\xi_z/2\pi \approx 1 \text{ kHz}$, Figs. 3(c) and 3(d) show the equivalent sensitivities of the Ξ -type and Λ -type NIT configurations to the RTN field, which reach $9.8 \text{ nT}/\sqrt{\text{Hz}}$ and $2.5 \text{ nT}/\sqrt{\text{Hz}}$, respectively. The significant advantage of Λ -type NIT arises not only from lower background noise but also from its noise-involved interferometric readout mechanism. Amplitude-dependent noise modulation enables the equivalent detection channel to be statistically synchronized with the RTN under test, thereby forming an enhanced response similar to that of phase-locked and frequency-locked circuits. Spectroscopically, this manifests as a narrower effective linewidth and a steeper response slope, meaning that even small changes in RTN parameters can cause significant changes in the window width and depth, thereby improving the sensitivity of RTN measurements. Therefore, Λ -type NIT is particularly suitable for highly sensitive characterization of low-frequency, non-Gaussian noise fields, such as the noise spectrum of quantum devices and low-frequency magnetic anomalies in the environment.

VI. CONCLUSIONS

We propose and experimentally demonstrate that RTN is not merely a performance burden but can be transformed into a controllable resource in the OPAM, opening

an NIT window in the spectral domain. Two complementary realizations are demonstrated on the same platform: a Ξ -type NIT driven by a single RTN channel and a Λ -type NIT enabled by two amplitude-correlated, flip-synchronized noise fields. For both configurations, we develop a unified analysis that quantifies the emergence of the transparency window and the evolution of peak splitting, and establishes how the LEP and its eigenvalue restructuring map onto experimentally resolvable spectral thresholds, thereby translating dynamical criticality into observable spectroscopic fingerprints.

Mechanistically, the Λ -type uses a noise-assisted detection field that is amplitude correlated with the target RTN and synchronized in flipping rate, which enhances sensitivity to weak, low-frequency fluctuations. This extends OPAMs to precise low-frequency noise characterization and marks a practical step toward optimizing sensors operating under non-Gaussian environmental noise. The protocol is generally applicable to NIT-style spectroscopy of defect reservoirs in solid-state devices [76,77] and to quantum state engineering in open environments [48].

ACKNOWLEDGMENTS

This work is supported by the National Natural Science Foundation of China (Grants No. 12374330, No. 12404408, and No. 12174139) and the Key Research and Development Special Project of Hubei Provincial Technology Innovation Program (Grant No. 2023BCB086).

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: THEORETICAL DERIVATION OF DYNAMICAL EVOLUTION EQUATIONS FOR OPAMS WITH RTN

The telegraph process, as a special class of Markov chain, can be described by the CKS equation [52,53]

$$\begin{aligned} \partial_t p(\alpha, t|\alpha_0, t_0) = & -c(\alpha, t)p(\alpha, t|\alpha_0, t_0) \\ & + \sum_{\beta} c(\alpha, t)f(\alpha|\beta, t) \\ & \times p(\beta, t|\alpha_0, t_0), \end{aligned} \quad (\text{A1})$$

where $p(\alpha, t|\alpha_0, t_0)$ is the joint probability distribution of the chain's possible states α and α_0 . The coefficient $c(\alpha, t)$ is the rate at which the telegraph signal in state α at time t changes its state during the time interval $(t, t + dt)$. The function $f(\alpha|\beta, t)$ is the conditional probability of a transition from β during the interval $(t, t + dt)$, given that the

transition ends in state α . For the two-valued telegraph process considered here, we have $c(\alpha, t) = \gamma$ and $f(\alpha|\beta; t) = \delta_{\alpha,-\beta}$, and the Eq. (A1) simplifies to

$$\begin{aligned} \partial_t p(\alpha, t|\alpha_0, t_0) = & -\gamma p(\alpha, t|\alpha_0, t_0) \\ & + \gamma p(-\alpha, t|\alpha_0, t_0), \end{aligned} \quad (\text{A2})$$

with the initial condition $p(\alpha, t_0|\alpha_0, t_0) = \delta_{\alpha,\alpha_0}$ [78]. The dynamical evolution of the system is governed by the master equation

$$\partial_t \rho = -i[H(t), \rho] + \mathcal{L}\rho, \quad (\text{A3})$$

where ρ and $\mathcal{L}\rho$ denote the density matrix and relaxation process of the system, respectively. We use the following relations between the density matrix and atomic multipole moments, the angular momentum and irreducible tensors, and the tensor commutation relations:

$$\rho = \sum_{k=0}^{2F} \sum_{q=-k}^k m_{k,q} T_{k,q}(F), \quad (\text{A4})$$

$$T_{1,q}(F) = \frac{\sqrt{3}}{\sqrt{F(F+1)(2F+1)}} F_q \quad (q = 0, \pm 1), \quad (\text{A5})$$

$$\begin{aligned} [T_{k_1,q_1}(F), T_{k_2,q_2}(F)] \\ = \sqrt{(2k_1+1)(2k_2+1)} \\ \times \sum_{k_3} (-1)^{2F+k_3} [1 - (-1)^{k_1+k_2+k_3}] \\ \times \left\{ \begin{matrix} k_1 & k_2 & k_3 \\ F & F & F \end{matrix} \right\} C_{k_1,q_1;k_2,q_2}^{k_3,q_3} T_{k_3,q_3}(F), \end{aligned} \quad (\text{A6})$$

where the physical quantities $m_{k,q}$ and $T_{k,q}(F)$ are the atomic multipole moment and the irreducible tensor of the q th component of order k , respectively. The symbols $C_{\dots}^{\dots}$ and $\{\dots\}$ are the Clebsch-Gordan coefficient and the Wigner $6j$ symbols of angular momentum, respectively. Next, we substitute Eqs. (1) and (A4)–(A6) into Eq. (A3). According to Eq. (A2), the evolution equation of the atomic multipole moments in the presence of telegraph noise is given by

$$\begin{aligned} \partial_t m_{k,q,\mu} = & -i(\mathcal{K}_{k,q,\mu,0} m_{k,q,\mu} + \mathcal{K}_{k,q,\mu,1} m_{k,q+1,\mu} \\ & - \mathcal{K}_{k,q,\mu,-1} m_{k,q-1,\mu}) \\ & - \gamma(m_{k,q,\mu} - m_{k,q,-\mu}) - \Gamma(m_{k,q,\mu} - \frac{1}{2} \delta_{q,0} \mathcal{F}_{z,\text{eq}}) \end{aligned} \quad (\text{A7})$$

for $k = 0, 1, \dots, 2F$, $q = -k, -k+1, \dots, k$ and $\mu = \pm 1$, with

$$\mathcal{K}_{k,q,\mu,0} := \sqrt{k(k+1)} C_{1,q;1,0}^{k,q} (\Delta_z + \xi_{z,\mu}), \quad (\text{A8a})$$

$$\mathcal{K}_{k,q,\mu,1} := \sqrt{k(k+1)} C_{1,q+1;1,-1}^{k,q} (\Omega_x + \xi_{x,\mu}), \quad (\text{A8b})$$

$$\mathcal{K}_{k,q,\mu,-1} := \sqrt{k(k+1)} C_{1,q-1;1,1}^{k,q} (\Omega_x + \xi_{x,\mu}), \quad (\text{A8c})$$

where the subscripts k and q represents the q component of the k th atomic multipole moment. The subscript μ takes the values 1 and -1 , representing the two states of the RTN.

1. Dynamical evolution equations of Ξ -type NIT for the dc/ac OPAM with longitudinal RTN

In this case, the index k is set to 1, and the symmetric component $m_{1,q,S} = m_{1,q,1} + m_{1,q,-1}$ and anti-symmetric component $m_{1,q,A} = m_{1,q,1} - m_{1,q,-1}$ generated by the telegraph-noise-induced system are introduced. In addition, the transformation relations $\{\mathcal{F}_{-j} = \sqrt{2}m_{1,-1,j}, \mathcal{F}_{+j} = -\sqrt{2}m_{1,1,j}, \mathcal{F}_{z,j} = m_{1,0,j}\}$ ($j = A, S$) between the atomic multipole moments and the spin moments are used. The evolution equations of the spin moments $[\mathcal{F}_{\mu,j} \ (\mu = \pm, z; j = A, S)]$ in the presence of telegraph noise are obtained as

$$\dot{\mathcal{F}}_{-,S} = -(\Gamma - i\Delta_z) \mathcal{F}_{-,S} + i\xi_z \mathcal{F}_{-,A} - i\Omega_x \mathcal{F}_{z,S}, \quad (\text{A9a})$$

$$\dot{\mathcal{F}}_{+,S} = -(\Gamma + i\Delta_z) \mathcal{F}_{+,S} - i\xi_z \mathcal{F}_{+,A} + i\Omega_x \mathcal{F}_{z,S}, \quad (\text{A9b})$$

$$\dot{\mathcal{F}}_{-,A} = -(\gamma' - i\Delta_z) \mathcal{F}_{-,A} + i\xi_z \mathcal{F}_{-,S} - i\Omega_x \mathcal{F}_{z,A}, \quad (\text{A9c})$$

$$\dot{\mathcal{F}}_{+,A} = -(\gamma' + i\Delta_z) \mathcal{F}_{+,A} - i\xi_z \mathcal{F}_{+,S} + i\Omega_x \mathcal{F}_{z,A}, \quad (\text{A9d})$$

$$\dot{\mathcal{F}}_{z,S} = -\Gamma (\mathcal{F}_{z,S} - \mathcal{F}_{z,\text{eq}}) + i\Omega_x (\mathcal{F}_{+,S} - \mathcal{F}_{-,S}), \quad (\text{A9e})$$

$$\dot{\mathcal{F}}_{z,A} = -\gamma' \mathcal{F}_{z,A} + i\Omega_x (\mathcal{F}_{+,A} - \mathcal{F}_{-,A}), \quad (\text{A9f})$$

with $\gamma' = 2\gamma + \Gamma$. From Eqs. (A9a)–(A9f), it can be seen that the evolution equations of the dc magnetic resonance system with longitudinal telegraph noise are similar to those of the Ξ -type EIT system [60]. To illustrate this relationship more intuitively, we reproduce two key equations of the Ξ -type EIT given in Ref. [60]:

$$\begin{aligned} \dot{\rho}_{12} = & -(\gamma_{12} - i\Delta_p) \rho_{12} - i\Omega_c \rho_{13} \\ & + i\Omega_p (\rho_{22} - \rho_{12}), \end{aligned} \quad (\text{A10a})$$

$$\begin{aligned} \dot{\rho}_{13} = & -[\gamma_{13} - i(\Delta_p + \Delta_c)] \rho_{13} \\ & + i\Omega_p \rho_{23} - i\Omega_c \rho_{12}. \end{aligned} \quad (\text{A10b})$$

By comparing and analyzing our Eqs. (A9b) and (A9d) with the Ξ -type EIT Eqs. (A10a) and (A10b) respectively, we find that the two sets of equations are almost completely consistent in form, with only slight differences in detail. To systematically reveal the differences and connections between our Ξ -type NIT configuration and the Ξ -type EIT system, we refer to the physical quantities listed in Table II and compare and analyze the various physical quantities of the two systems one by one.

Driving fields. In the Ξ -type EIT system, the pump field ω_c and the probe field Ω_p couple the levels $|2\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |1\rangle$, respectively, forming a typical cascade structure. In our Ξ -type NIT system, the corresponding driving

TABLE II. Comparison between our Ξ -type NIT system and the Ξ -type EIT system in Ref. [60]. The probe and pump Rabi frequencies, Ω_p and Ω_c , couple the transitions $|2\rangle \leftrightarrow |1\rangle$ and $|2\rangle \leftrightarrow |3\rangle$, respectively. The detunings Δ_p and Δ_c denote the frequency offsets of the probe and pump fields from the corresponding transitions.

System	Pump field	Probe field	Density-matrix element/spin moment				Detuning		Decay rate	
Ref. [60]	Ω_c	Ω_p	ρ_{12}	ρ_{13}	$\rho_{22} - \rho_{11}$	ρ_{23}	Δ_p	$\Delta_p + \Delta_c$	γ_{12}	γ_{13}
Ours	ξ_z	Ω_x	$\mathcal{F}_{+,S}$	$\mathcal{F}_{+,A}$	$\mathcal{F}_{z,S}$	$\mathcal{F}_{z,A}$	Δ_z	Δ_z	Γ	γ'

quantities are the longitudinal RTN ξ_z and the transverse rf drive ω_x . Among them, ξ_z provides an indirect driving effect between the coupling channels, similar to the role of the pump field, while Ω_x directly acts on the transverse component of the spin, corresponding to the probe field.

Density matrix elements and spin moments. In the Ξ -type EIT system, ρ_{12} and ρ_{13} are two key density-matrix coherence terms, describing the quantum coherence of $|1\rangle \leftrightarrow |2\rangle$ and $|1\rangle \leftrightarrow |3\rangle$, respectively. In our system, $\mathcal{F}_{+,S}$ and $\mathcal{F}_{+,A}$ are analogous to these two variables, representing the coherence of the transverse spin components in the symmetric and antisymmetric modes, respectively, and their evolution structurally corresponds to the evolution of ρ_{12} and ρ_{13} . The population-difference term $\rho_{22} - \rho_{11}$ in the Ξ -type EIT system is the modulation factor of the detection-field coupling strength and directly affects the formation of the EIT window. In our system, $\mathcal{F}_{z,S}$ plays a completely equivalent role as a longitudinal symmetric component.

Detunings. In the Ξ -type EIT system, the two driving fields have their own detunings Δ_p and Δ_c , and the frequency-matching condition in the coupling term appears as their sum. In our configuration, although the driving signal is in the form of noise, it still possesses frequency characteristics. These are manifested as a unified detuning term Δ_z , which participates in the interference path in the same way in different spin channels, corresponding to the total-detuning structure in the Ξ -type EIT system.

Relaxation coefficients. The main difference is that, for the Ξ -type EIT system, the decay rate γ_{13} of the density-matrix element ρ_{13} is much smaller than the decay rate γ_{12} of the density-matrix element ρ_{12} . However, for our system, the decay rate γ' of the spin moment $\mathcal{F}_{+,A}$ is much greater than the decay rate Γ of the spin moment $\mathcal{F}_{+,S}$. This is also the main reason why full transparency cannot be achieved in this configuration.

2. Dynamical evolution equations of Λ -type NIT for the ac OPAM with amplitude-correlated RTN

Similar to the processing method in Appendix A 1, taking $k = 1$, the evolution equations of the spin moments with amplitude-correlated telegraph noise are obtained as

$$\begin{aligned} \dot{\mathcal{F}}_{-,S} = & -(\Gamma - i\Delta_z) \mathcal{F}_{-,S} + i\xi_z \mathcal{F}_{-,A} \\ & - i\xi_x \mathcal{F}_{z,A}, \end{aligned} \quad (\text{A11a})$$

$$\begin{aligned} \dot{\mathcal{F}}_{+,S} = & -(\Gamma + i\Delta_z) \mathcal{F}_{+,S} - i\xi_z \mathcal{F}_{+,A} \\ & + i\xi_x \mathcal{F}_{z,A}, \end{aligned} \quad (\text{A11b})$$

$$\begin{aligned} \dot{\mathcal{F}}_{-,A} = & -(\gamma' - i\Delta_z) \mathcal{F}_{-,A} + i\xi_z \mathcal{F}_{-,S} \\ & - i\xi_x \mathcal{F}_{z,S}, \end{aligned} \quad (\text{A11c})$$

$$\begin{aligned} \dot{\mathcal{F}}_{+,A} = & -(\gamma' + i\Delta_z) \mathcal{F}_{+,A} - i\xi_z \mathcal{F}_{+,S} \\ & + i\xi_x \mathcal{F}_{z,S}, \end{aligned} \quad (\text{A11d})$$

$$\begin{aligned} \dot{\mathcal{F}}_{z,S} = & -\Gamma (\mathcal{F}_{z,S} - \mathcal{F}_{z,\text{eq}}) \\ & + i\xi_x (\mathcal{F}_{+,A} - \mathcal{F}_{-,A}), \end{aligned} \quad (\text{A11e})$$

$$\dot{\mathcal{F}}_{z,A} = -\gamma' \mathcal{F}_{z,A} + i\xi_x (\mathcal{F}_{+,S} - \mathcal{F}_{-,S}). \quad (\text{A11f})$$

From Eqs. (A11a)–(A11f), it can be seen that the evolution equations of the ac magnetic resonance system with amplitude-correlated RTN are similar to those of the Λ -type EIT system [72]. To illustrate this relationship more intuitively, we reproduce two key equations of the Λ -type EIT given in Ref. [72]:

$$\begin{aligned} \dot{\rho}_{ab} = & -(\gamma_1 + i\Delta_v) \rho_{ab} - i\Omega_v (\rho_{aa} - \rho_{bb}) \\ & + i\Omega_\mu \rho_{cb}, \end{aligned} \quad (\text{A12a})$$

$$\begin{aligned} \dot{\rho}_{cb} = & -[\gamma_3 + i(\Delta_v + \Delta_\mu)] \rho_{cb} - i\Omega_v \rho_{ca} \\ & + i\Omega_\mu \rho_{ab}. \end{aligned} \quad (\text{A12b})$$

By comparing and analyzing our Eqs. (A11b) and (A11d) with the Λ -type EIT equations (A12a) and (A12b), respectively, we can find that the two sets of equations are almost completely consistent in form. In Table III, the evolution equations of our system are compared in detail with those of the Λ -type EIT system.

Analogy of pump and probe fields. In the Λ -type EIT system, the pump field Ω_μ and the probe field Ω_v are two classical light fields that excite the $|a\rangle \leftrightarrow |c\rangle$ and $|a\rangle \leftrightarrow |b\rangle$ transitions, respectively. In our system, the same roles are played by the transverse and longitudinal RTN sources ξ_x and ξ_z , which drive the spin-moment changes in different directions, thereby producing a coupling effect analogous to the corresponding transitions.

Analogy between density matrix and spin moments. The density-matrix elements ρ_{ab} , ρ_{cb} , ρ_{ca} , and $(\rho_{aa} - \rho_{cc})$ in the Λ -type EIT system correspond to $\mathcal{F}_{+,S}$, $\mathcal{F}_{+,A}$, $\mathcal{F}_{z,S}$, and $\mathcal{F}_{z,A}$ in our system, respectively. These quantities reflect the one-to-one correspondence between the spin-moment

TABLE III. Comparison between our Λ -type NIT system and the Λ -type EIT system in Ref. [72]. The probe and pump Rabi frequencies, Ω_v and Ω_μ , couple the transitions $|a\rangle \leftrightarrow |b\rangle$ and $|a\rangle \leftrightarrow |c\rangle$, respectively. The detunings Δ_v and Δ_μ denote the frequency offsets of the probe and pump fields from their corresponding transitions.

System	Pump field	Probe field	Density-matrix element/spin moment				Detuning		Decay rate	
Ref. [72]	Ω_μ	Ω_v	ρ_{ab}	ρ_{cb}	$\rho_{aa} - \rho_{bb}$	ρ_{ca}	Δ_p	$\Delta_v + \Delta_\mu$	γ_1	γ_3
Ours	ξ_z	ξ_x	$\mathcal{F}_{+,S}$	$\mathcal{F}_{+,A}$	$\mathcal{F}_{z,A}$	$\mathcal{F}_{z,S}$	Δ_z	Δ_z	γ'	Γ

components of our system and the density-matrix elements of the Λ -type EIT system.

Analogy between detuning and relaxation. In the Λ -type EIT system, the detunings Δ_μ and Δ_v between the light fields and the atomic transition determines the resonance conditions of the system. In our system, the corresponding detuning is Δ_z , which also determines the resonance condition of our system. Similarly, the decay rates γ_1 and γ_3 correspond to γ' and Γ in our system, respectively, reflecting the relaxation process of the system. Moreover, the condition $\gamma' \gg \Gamma$ satisfied by the relaxation coefficient is completely consistent with the condition $\gamma_1 \gg \gamma_3$ satisfied by the relaxation coefficients in the Λ -type EIT system, so that the transparency window induced by the amplitude-correlated telegraph noise can approach complete transparency.

In this configuration, the coupling between the two noise fields and the spin moments is

$$\mathcal{F}_{z,S} \xleftrightarrow[\text{probe}]{\xi_x} \mathcal{F}_{\pm,A} \xleftrightarrow[\text{pump}]{\xi_z} \mathcal{F}_{\pm,S},$$

where ξ_x acts as the probe field and ξ_z acts as the pump field. Therefore, we refer to the transparency phenomenon

induced by the two noise fields as “ Λ -type NIT.” However, for the scenario of a dc OPAM with longitudinal RTN, the coupling form of the noise field ξ_z and the detection field Ω_x with the spin moments is

$$\mathcal{F}_{z,S} \xleftrightarrow[\text{probe}]{\Omega_x} \mathcal{F}_{\pm,S} \xleftrightarrow[\text{pump}]{\xi_z} \mathcal{F}_{\pm,A}.$$

By analogy with the EIT coupling mode, this structure is similar to that of the Ξ -type EIT. Therefore, we refer to this noise-induced transparency effect as “ Ξ -type NIT”. This analogy provides an intuitive and physically reasonable explanation for the transparent window and coherence-enhancement effect observed in our system.

APPENDIX B: STEADY-STATE SPECTRUM AND LEP OF Λ -TYPE NIT FOR AN ac OPAM

To further gain insight into the physical picture of the RTN-induced transparency effect, the steady-state analytical spectrum is obtained from the spin-magnetic-moment evolution equations as

$$\mathcal{S}_2(\Delta_z) = F_{z,\text{eq}} \left(1 - \frac{2\xi_x^2}{\Gamma Z_2} \left\{ \Delta_z^2 (\Gamma + \gamma') - \Gamma \left[\Delta_z^2 - \Gamma\gamma' - \xi_z^2 \left(1 + \frac{2\xi_x^2}{\xi_z^2} \right) \right] \right\} \right), \quad (\text{B1})$$

with

$$Z_2 = [\Delta_z^2 + s_1^2(2\gamma'^2 + \Gamma\gamma') - \xi_z^2]^2 + 4s_1^2(\Gamma + \gamma')^2[\xi_z^2 - (\gamma s_1)^2], \quad (\text{B2})$$

and $s_1 = \sqrt{1 + \beta}$. The dimensionless parameter $\beta = 2\xi_x^2/(\Gamma\gamma')$ represents the RTN saturation effect. Equation (B1) is a typical Λ -type EIT spectrum. Especially under the condition of weak ξ_x ($\xi_x \rightarrow 0$), ignoring the saturation effect of the transverse RTN field ($\beta = 0$), spectrum (B1) becomes

$$\mathcal{S}_2(\Delta_z) = \mathcal{F}_{z,\text{eq}} \left[1 - \frac{2\xi_x^2}{\Gamma} \frac{\Delta_z^2\gamma' + \Gamma(\Gamma\gamma' + \xi_z^2)}{\Delta_z^4 + \Delta_z^2(\Gamma^2 + \gamma'^2 - 2\xi_z^2) + (\Gamma\gamma' + \xi_z^2)^2} \right]. \quad (\text{B3})$$

The form of the spectrum (B3) is very similar to the Λ -type EIT spectrum in Ref. [72]. Next, the spectrum (B3) can be written in a more explicit and physically transparent form, i.e.,

$$\mathcal{S}_2(\Delta_z) = \begin{cases} \mathcal{F}_{z,\text{eq}} \left\{ 1 - \frac{\xi_x^2}{4\Gamma\Pi} \left[\frac{(\gamma' + 2\Pi)^2 - \Gamma^2}{\Delta_z^2 + (\gamma + \Gamma + \Pi)^2} - \frac{(\gamma' - 2\Pi)^2 - \Gamma^2}{\Delta_z^2 + (\gamma + \Gamma - \Pi)^2} \right] \right\}, & \xi_z < \gamma, \\ \mathcal{F}_{z,\text{eq}} \left\{ 1 - \frac{2\xi_x^2}{\Gamma} \frac{\Delta_z^2 \gamma' + \Gamma(\gamma + \Gamma)^2}{[\Delta_z^2 + (\gamma + \Gamma)^2]^2} \right\}, & \xi_z = \gamma, \\ \mathcal{F}_{z,\text{eq}} \left(1 - \frac{\xi_x^2}{\Gamma} \left\{ \frac{\Gamma}{(\Delta_z - \Pi)^2 + (\gamma + \Gamma)^2} + \frac{\Gamma}{(\Delta_z + \Pi)^2 + (\gamma + \Gamma)^2} + \frac{4\Delta_z^2 \gamma}{[(\Delta_z - \Pi)^2 + (\gamma + \Gamma)^2][(\Delta_z + \Pi)^2 + (\gamma + \Gamma)^2]} \right\} \right), & \xi_z > \gamma, \end{cases} \quad (\text{B4})$$

with

$$\Pi = \begin{cases} \sqrt{\gamma^2 - \xi_z^2}, & \xi_z < \gamma, \\ 0, & \xi_z = \gamma, \\ \sqrt{\xi_z^2 - \gamma^2}, & \xi_z > \gamma. \end{cases} \quad (\text{B5})$$

Under weak detection conditions ($\xi_x \rightarrow 0$), the two spin-moment modes $\mathcal{F}_{+,S}$ and $\mathcal{F}_{+,A}$ in the evolution equations (A11a)–(A11f) for amplitude-correlated RTN are decoupled from other modes. The evolution of these two modes, written in terms of the Liouvillian superoperator, is

$$\begin{pmatrix} \dot{\mathcal{F}}_{+,S} \\ \dot{\mathcal{F}}_{+,A} \end{pmatrix} = \mathcal{L}_2 \cdot \begin{pmatrix} \mathcal{F}_{+,S} \\ \mathcal{F}_{+,A} \end{pmatrix}, \quad (\text{B6})$$

with

$$\mathcal{L}_2 = \underbrace{\begin{pmatrix} -\Gamma & -i\xi_z \\ -i\xi_z & -2\gamma - \Gamma \end{pmatrix}}_{\mathcal{L}_{2,s}} + \underbrace{\begin{pmatrix} -i\Delta_z & 0 \\ 0 & -i\Delta_z \end{pmatrix}}_{\mathcal{L}_{2,p}}. \quad (\text{B7})$$

Here $\mathcal{L}_{2,s}$ is the Liouvillian superoperator associated with the LEPs of the system. The superoperator $\mathcal{L}_{2,p}$ is related to the probe field ξ_x . By solving the determinant equation $\det(\mathcal{L}_{2,s} - E_2 \mathcal{I}) = 0$, the eigenvalues of the system can be obtained, where $\mathcal{I}$ is the 2×2 identity matrix. The eigenvalues for the amplitude-correlated RTN case are $E_{2,+} = \gamma + \Gamma + \sqrt{\gamma^2 - \xi_z^2}$ and $E_{2,-} = \gamma + \Gamma - \sqrt{\gamma^2 - \xi_z^2}$. The energy degeneracy point (LEP) also occurs at $\gamma = \xi_z$. This point is also the phase-transition point, which divides the

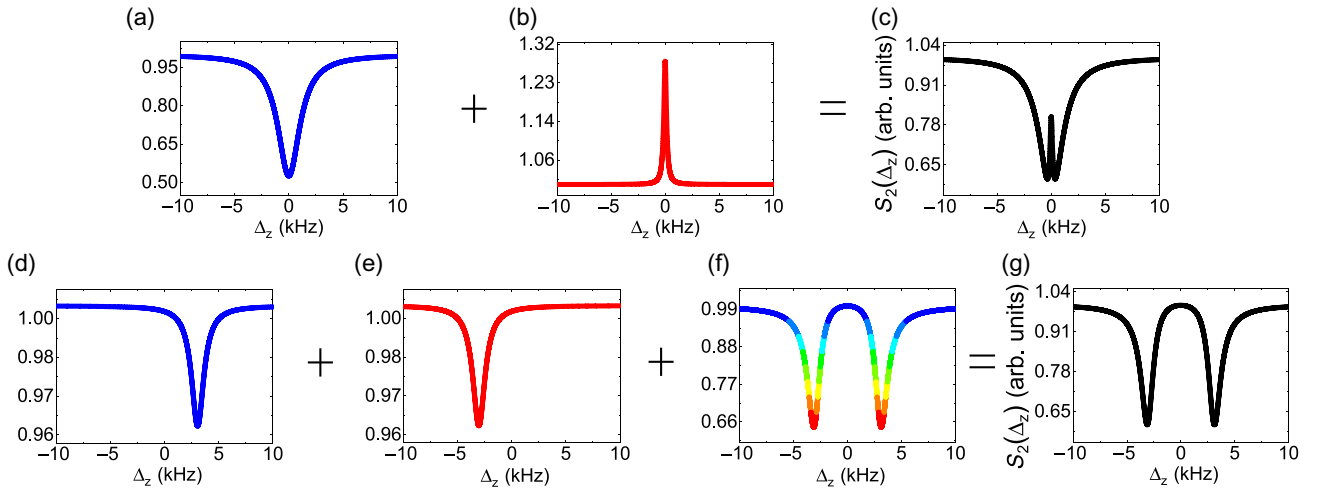


FIG. 4. The superposition-synthesis process of the spectrum $\mathcal{S}_2(\Delta_z)$ of the RTN amplitude ξ_z before and after the LEP. (a) The first Lorentzian peak in parentheses in spectrum (B4) for the RTN amplitude ξ_z before the LEP. (b) The second Lorentzian peak in parentheses in spectrum (B4) for the RTN amplitude ξ_z before the LEP. (c) The entire spectrum (B4) for the RTN amplitude ξ_z before the LEP. (d) The first Lorentzian peak in curly brackets in spectrum (B4) for the RTN amplitude ξ_z after the LEP. (e) The second Lorentzian peak in curly brackets in spectrum (B4) for the RTN amplitude ξ_z after the LEP. (f) The third term in curly brackets in the spectrum (B4) is the interference term between the two Lorentzian peaks for the RTN amplitude ξ_z after the LEP. (g) The entire spectrum (B4) for the RTN amplitude ξ_z after the LEP.

dynamical phase of the system into two parts: the broken phase ($\gamma > \xi_z$, weak coupling, nonoscillatory dynamics) and the exact phase ($\gamma < \xi_z$, strong coupling, oscillatory dynamics). Furthermore, the LEP can be understood as analogous to the critical-damping point of a damped harmonic oscillator, occurring at the phase transition between the oscillatory (exact phase) and nonoscillatory (broken phase) regimes of the dynamics. The two dynamical phases are again confirmed in the spectra (B4). Spectroscopically, Eq. (B4) shows that the RTN amplitude ξ_z below the LEP is composed of two Lorentzian peaks with the same center ($\Delta_z = 0$) but different linewidths ($E_{2,\pm}$), different amplitudes $[(\gamma' \pm 2\Pi)^2 - \Gamma^2]\xi_x^2/(4\Gamma\Pi)$ and opposite directions; see Figs. 4(a)–4(c). Since the amplitudes of these two Lorentzian peaks are of the same order of magnitude, an obvious transparency effect can occur even if the RNT amplitude ξ_z is very small. On the contrary, when the RTN amplitude ξ_z is above the LEP, the spectrum is composed of three parts: two Lorentzian peaks with different centers ($\pm\sqrt{\xi_z - \gamma^2}$) but the same linewidth ($\gamma + \Gamma$), amplitude (ξ_x^2), and direction, and the interference term between these two Lorentzian peaks; see Figs. 4(d)–4(g).

APPENDIX C: EXPERIMENT

1. Experimental setup

The experiments are performed using room-temperature cesium vapor confined in vacuum within a paraffin-coated cylindrical glass cell ($28 \times 40 \text{ mm}^2$) from Precision Glass-blowing Company. Experimental measurements show that the half width at half height of the atomic cell is about 46.4 Hz. The experimental optical setup is shown in Fig. 5. The cesium vapor cell is isolated from the surrounding magnetic field by a six-layer mu-metal shield,

where the residual magnetic field fluctuation of the magnetic shield here is less than 10 nT. Inside the shield, a nonmoment solenoid coil driven by a precision current source (Keysight B2962B) generates a static magnetic field B_0 along the light-propagation direction. Simultaneously, a high-precision multimeter (Agilent 3458A 8.5-digit multimeter) is connected to monitor the current drift of the precision current source. All rf fields used in the experiments are generated by driving three-dimensional Helmholtz coils wrapped with solenoids using an arbitrary waveform generator (RIGOL DG4162). The random telegraph fields are generated by driving the Helmholtz coil with our homemade two-channel random-telegraph source. The laser beam used to pump and probe the atomic vapor is generated by a Toptica tunable single-frequency diode laser ($\approx 894 \text{ nm}$) and is laser locked to the $6S_{1/2}, F_g = 4 \leftrightarrow 6P_{1/2}, F_e = 3$ hyperfine-structure component by saturated-absorption spectroscopy (SAS). Before the laser enters the magnetic shield, a quarter-wave plate followed by a polarization beam splitter produces circularly polarized light. After passing through the quarter-wave plate, the remaining noncircular polarization degree is less than 1%. This laser (power about $30 \mu\text{W}$) prepares an orientation-polarized state in the ground-state hyperfine structure $F_g = 4$. The light power transmitted through the cell is detected by an adjustable optical receiver (Newport 2051-FS-M), and its photocurrent is amplified by a low-noise transimpedance amplifier. The resulting signal is acquired by an oscilloscope (Keysight InfiniVision DSOX6004A).

a. Experimental design of Ξ -type NIT

A precision current source drives the solenoid along the z direction to generate a constant magnetic field B_0

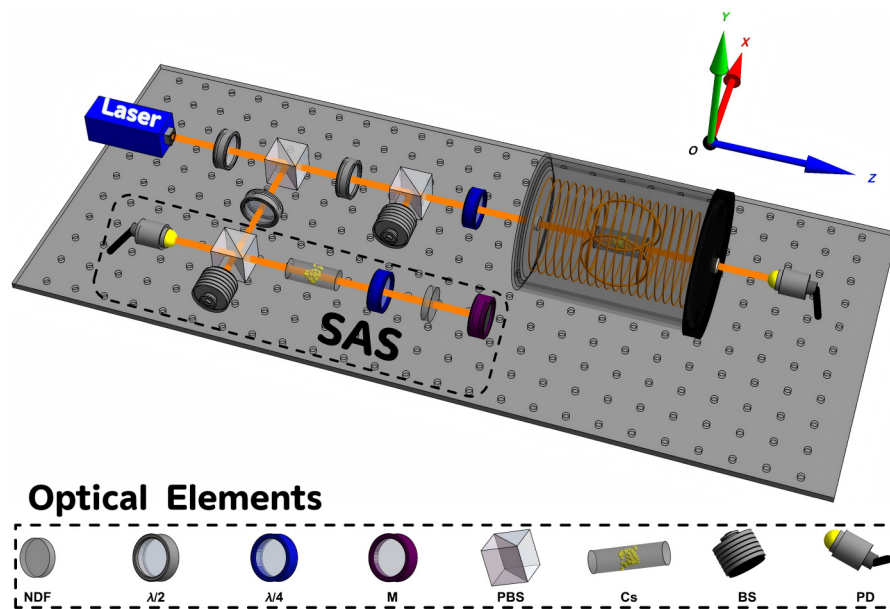


FIG. 5. Schematic of the experimental setup. NDF, neutral-density filter; $\lambda/2$, half-wave plate; $\lambda/4$, quarter-wave plate; M, mirror; PBS, polarization beam splitter; Cs, cesium atomic cell; BS, beam stopper; PD, photodetector; SAS, saturated-absorption spectroscopy. The electronic equipment used in the experiments is not shown; see the text for more details.

that uniformly splits the magnetic sublevel $F_g = 4$, and its Larmor-precession frequency is $\Omega_0/2\pi = 100$ kHz. The RTN source drives the Helmholtz coil in the z direction, and the resulting telegraph-noise magnetic field $\mu_1(t)$ couples to the sublevels to modulate the Larmor-precession frequency Ω_0 . This sets up a noise-dressed large-spin system. The weak rf field $B_2(t) = B_2 \cos \omega_2 t$ is introduced to drive the Helmholtz coil in the x direction to generate a transverse detection field, and the system properties are observed by slowly scanning the frequency of $B_2(t)$. The slow-scanning detection method is equivalent to performing a Fourier transform on the system, converting the time-domain information of the steady-state system into frequency-domain information. The center frequency of the detection field $B_2(t)$ is 100 kHz, the frequency span is 20 kHz, and the scanning time is 5 s. When the frequency ω_2 of the detection field $B_2(t)$ scans near the Larmor-precession frequency, population transfer between the sublevels occurs, thereby destroying the orientation polarization established by the pump-probe laser, which in turn causes the laser to continuously pump the atoms to establish a new equilibrium. In this process, the system information detected by the rf field $B_2(t)$ is recorded in the laser transmission spectrum. This detection technology is the so-called optical-radio-frequency double-resonance detection [79].

b. Experimental design of Λ -type NIT

Similar to Appendix C 1 a, a precision current source drives the solenoid in the z direction to produce a Larmor-precession frequency $\Omega_0/2\pi$ of 100 kHz. The two channels of the RTN source output noise $\mu_1(t)$ without spectrum shifting and noise $\mu_2(t) \cos(\omega_1 t)$ with spectrum shifting, respectively, which are added to the Helmholtz coil in the z direction through an adder. The flip frequency γ of both noises is 3 kHz. The noise-offset frequency $\omega_1/2\pi$ is 30 kHz. Note that these two noises are from the same source and have amplitude-correlation characteristics. These two noise fields are equivalent to an amplitude-modulated noise field. This field configuration constructs a large-spin system dressed by amplitude-correlated noise. The detection method still uses the optical-radio-frequency double-resonance detection protocol. The arbitrary waveform generator outputs an rf signal in linear-scanning mode to drive the Helmholtz coil in the x direction to generate an rf magnetic field $B_2(t) = B_2 \cos \omega_2 t$. The center frequency of the detection field $B_2(t)$ is 70 kHz, the frequency span is 20 kHz, and the scanning time is 5 s. By scanning the frequency of $B_2(t)$ to perform Fourier decomposition of the system, the system spectrum information is extracted. The system information is ultimately obtained from the oscilloscope display of the laser-transmission spectrum received by the optical receiver.

2. Random telegraph noise source

The RTN used in this paper is generated by our homemade two-channel frequency-offset-adjustable RTN source. The overall design scheme of the random-telegraph-noise source system is shown in Fig. 6, which is mainly composed of a human-machine interaction module, an exponential random-number generation module, a direct digital synthesis (DDS) module, and an output module. The exponential random-number generation module and the DDS module are implemented in a field-programmable gate array (FPGA). A critical pseudorandom-number generator module is designed, which utilizes the Mersenne Twister 19937 (MT19937) algorithm [80] to generate uniformly distributed pseudorandom numbers. Exponentially distributed pseudorandom numbers are calculated according to the inverse-function method, thereby reducing the occupation of hardware resources. The pseudorandom-number capacity generated by the algorithm is as high as $2^{19937} - 1$. When phase switching occurs at a frequency of 3 kHz, the maximum working time of the random-telegraph-noise source can reach 1.66×10^{5993} days, so that the designed pseudorandom telegraph-noise source can be approximately regarded as a true random telegraph-noise source. The inverse-function method is then used to calculate the exponential random numbers, which are used as the flip time intervals to generate the phase flips. The DDS module uses a phase accumulator and a waveform data table to combine the offset frequency f_0 with the flipped phase to generate a digital random-telegraph-noise signal. Through the digital-to-analog (D/A) converter in the output module, the random-telegraph-noise digital signal is converted into an analog signal. A low-pass filter is used to filter out high-frequency noise. Finally, the amplitude of the signal is amplified by an amplifier circuit, and the random-telegraph-noise analog signal is output. As for the human-machine interaction module, its core is a microprocessor. The adjustment knob and keys are connected to the microprocessor to adjust the parameters related to random telegraph noise, and the LCD connected to the microprocessor displays the random-telegraph-noise parameters. The microprocessor communicates with the FPGA to transmit the offset frequency f_0 and the phase-flip frequencies γ_0 and γ_1 . The microprocessor controls the reference voltage of the D/A converter through a slow-control DAC to adjust the amplitude of the random telegraph noise. Our RTN source is designed as a two-channel system, in which both channels share the same random switching sequence. That is, the state transitions—i.e., switching from one level to another—occur simultaneously in both channels, driven by a common random process. The synchronous flipping introduces a strong correlation between the two outputs.

The test results of our RTN source is shown in Fig. 7. Its time-domain signal and spectrum characteristics are consistent with the theoretical simulations and fully satisfy the

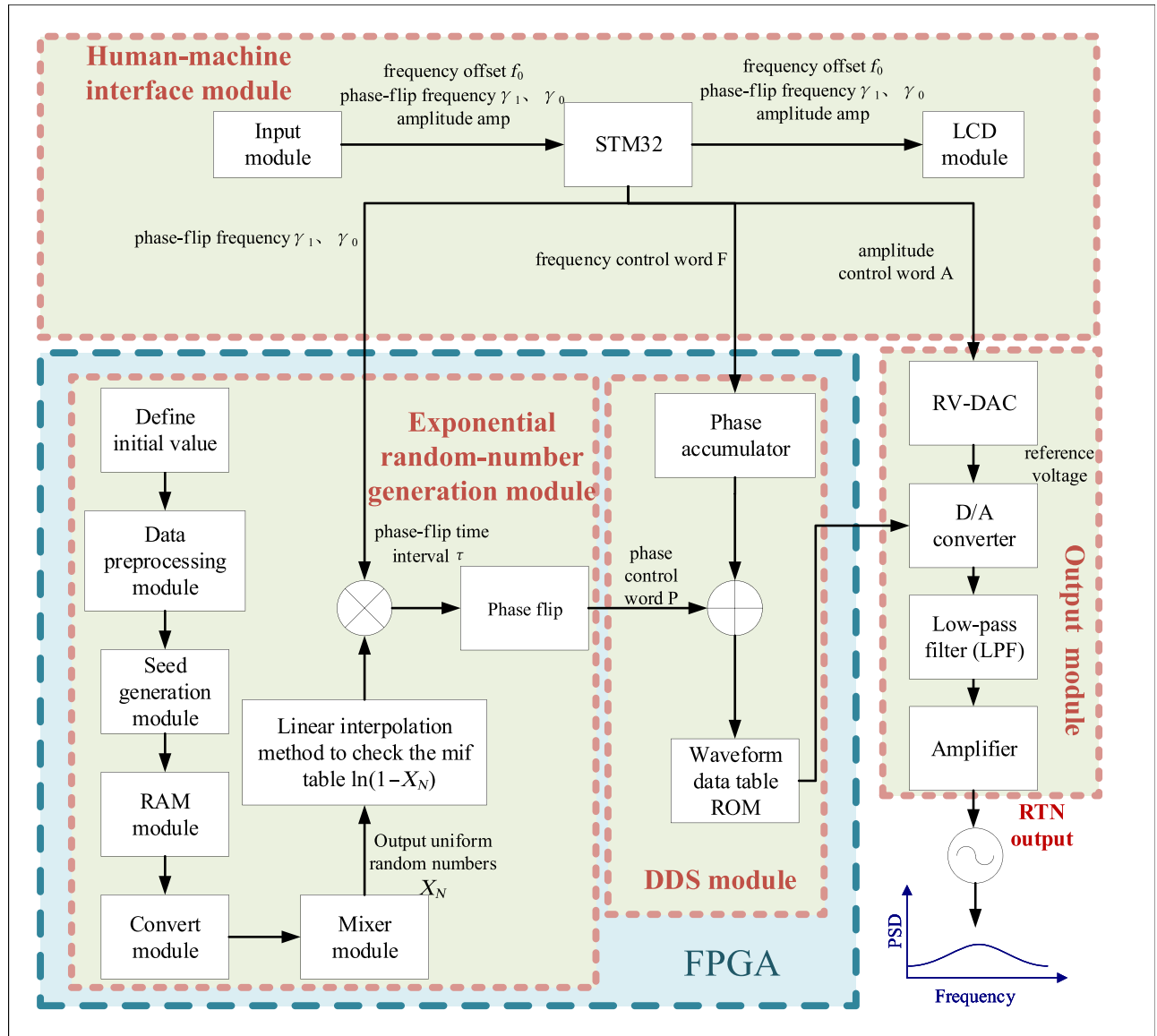


FIG. 6. System block diagram of a homemade multiparameter adjustable RTN source. The system consists of four main parts: the human-machine interface module, the exponential random-number generation module, the DDS module, and the output module. Users input parameters such as the frequency offset f_0 , phase-flip frequencies γ_1 and γ_0 , and amplitude (amp) via the input module. These are processed by an STM32 controller to generate the frequency control word F and amplitude control word A , which are also displayed on the liquid-crystal-display (LCD) module. Within the FPGA, the exponential random-number generation module produces phase-flip time intervals τ based on seed generation, linear interpolation, and lookup-table operations. The DDS module then uses a phase accumulator and waveform ROM to generate sine waveforms. These signals are converted to analogue via a D/A converter, filtered through a low-pass filter (LPF), amplified, and finally output as the RTN signal.

requirements of our experiment. When the parameters of the RTN source are set to $\gamma_1 = \gamma_0 = \gamma$ and $f_0 = 0$, the RTN source outputs telegraph noise $\mu_1(t)$ without a frequency offset. Applying this RTN field directly along the z direction of the standard dc OPAM system constitutes our first case—the Ξ -type noise-induced transparency. When the parameters of the first channel of the RTN source are set to $\gamma_1 = \gamma_0 = \gamma$ and $f_0 = 0$, and the parameters of the second channel are set to $\gamma_1 = \gamma_0 = \gamma$ and $f_0 = \omega_1/2\pi$, the

two channels of the RTN source output telegraph noises $\mu_1(t)$ and $\mu_2(t) \cos(\omega_1 t)$ without a frequency offset and with a frequency offset, respectively. Adding these two RTN fields $(\mu_1(t) + \mu_2(t) \cos(\omega_1 t))$ directly along the z direction of the ac OPAM system constitutes our other case—the Δ -type NIT. Moreover, these two noises can be written in standard amplitude-modulation form, i.e., $\mu_{AM}(t) = \mu_1(t)[1 + m(t) \cos(\omega_1 t)]$, where $\mu_1(t)$ is equivalent to the envelope amplitude varying with time. The

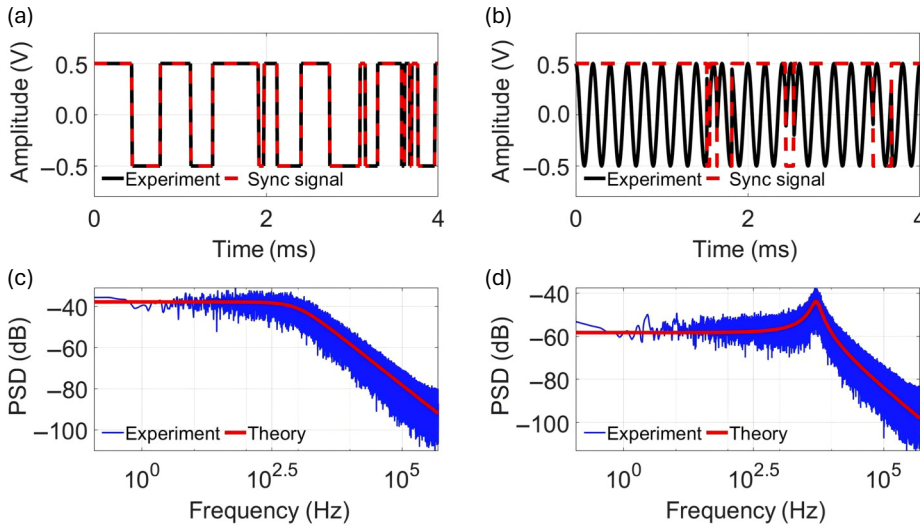


FIG. 7. Time- and frequency-domain signals of RTN with a telegraph-switching frequency of 3 kHz. Subplots (a) and (c) are the time-varying signal and the power spectral density (PSD) of RTN without a frequency offset, respectively. Subplots (b) and (d) show the time-varying signal and the PSD of RTN with a frequency offset of 5 kHz, respectively.

amplitude ratio $m(t) = \mu_2(t)/\mu_1(t)$ between the noises from the same source is the normalized modulation index.

3. Calibration of Helmholtz coils driven by rf and RTN sources

To achieve precise control and quantitative evaluation of the transverse rf field and the RTN field, we systematically calibrated the electrical characteristics of the Helmholtz coils in the x and z directions. Since the transverse rf field and the longitudinal RTN field are different from the longitudinal constant magnetic field and cannot be directly measured from the Larmor frequency, they must be indirectly measured in combination with the electrical response of the coil.

According to the principle of magnetic resonance, during resonant driving the Rabi frequency Ω_{rf} and the driving rf amplitude B_{rf} satisfy the relationship $\Omega_{rf} = \gamma_L B_{rf}$, where γ_L is the gyromagnetic ratio. At the same time, the magnetic field B_{rf} generated by the Helmholtz coil is linearly related to its driving current I_{rf} , that is, $B_{rf} = k I_{rf}$, where k is the proportional constant determined by the geometry of the coil. In addition, the output voltage U_{rf} of the signal source is linearly related to the coil response current I_{rf} , that is, $I_{rf} = Y U_{rf}$, where Y is the equivalent admittance. Therefore, experimentally, the magnetic field amplitude B_{rf} can be inferred by measuring the voltage U_{rf} , current I_{rf} , or Rabi frequency Ω_{rf} .

Figure 8(a) shows the current-voltage (I - V) characteristics of the Helmholtz coil at an excitation frequency of 100 kHz. The relative error of the slope of the fitting result is about 0.006%, indicating that the coil response is highly linear and that the measurement error is extremely small. The inset of Fig. 8(a) further shows the relationship between the Rabi frequency and the change in the coil current. The relative error of the intercept of the fitting result is about 0.2%. Ideally, this relationship should

be strictly proportional; that is, the Rabi frequency should be zero when the current is zero. However, a small nonzero intercept (0.0006 ± 0.0006) appeared in the experimental fitting. Although it is not statistically significant (its confidence interval covers zero), it may still reflect the weak background field in the system, system drift in the frequency-fitting process, or an instrument measurement offset under low signal. Given that its error range is comparable to the fitted value, the fitting result can be considered consistent with the above prediction and does not affect the accuracy of the magnetic field estimation.

Figure 8(b) is used to examine the stability of the coil at different frequencies, showing the trend of the current changing with the excitation frequency under three driving voltages (0.2, 0.4, and 0.6 V). The relative errors of the three slopes are 0.207%, 0.128%, and 0.024%, respectively, indicating that the coil response has good linearity and stability. Although the fitting results show that the current decreases slightly as the frequency increases, the change is very small (less than 0.025 mA per kilohertz). This can be regarded as the frequency-response error being controlled within an acceptable range, which further verifies the applicability of the calibration curve within the working frequency band.

Figures 8(c) and 8(d) show the I - V fitting results of our homemade RTN signal source without frequency shift and with frequency shift, respectively. The relative errors of the slopes are 0.188% and 0.131%, respectively. Although the fitting curve should theoretically pass through the origin, the intercept of the fitting curve is extremely small and the uncertainty covers near zero, indicating that the offset is not statistically significant. The RTN source output shows good stability and linear response.

In summary, we have systematically evaluated the error in the response characteristics of the Helmholtz coil under continuous sinusoidal and RTN-noise excitation. The relative errors of all key fitting slopes are less than 0.2%, the

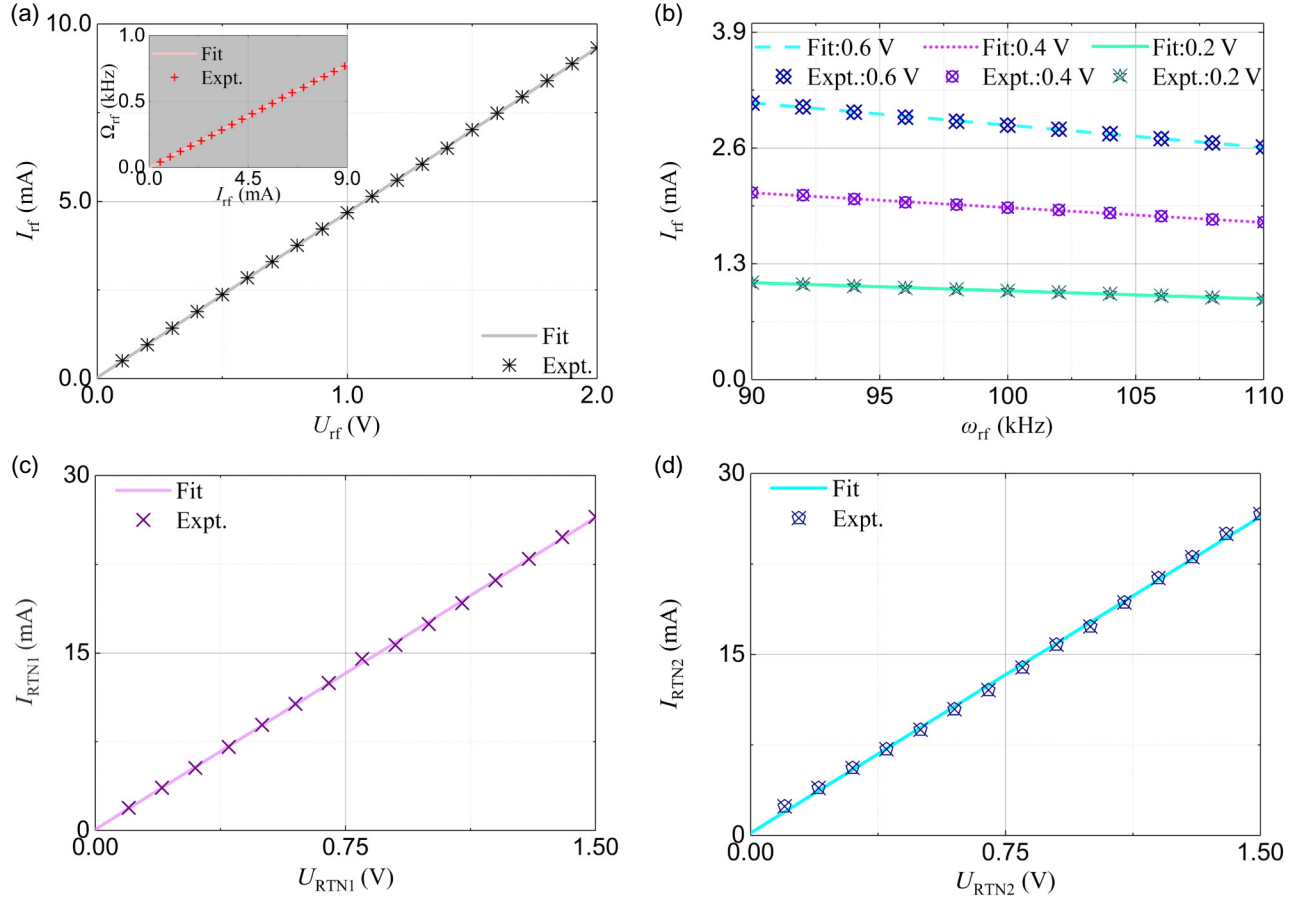


FIG. 8. Calibration of the Helmholtz coils in the x and z directions. (a) The current-voltage (I - V) characteristic curve of the Helmholtz coil with an rf frequency of 100 kHz in the x direction. The fit result is $I_{\text{rf}} = (4.648\,65_{\pm 0.000\,28})U_{\text{rf}} + (0.000\,09_{\pm 0.000\,22})$. The inset is the corresponding current-Rabi-frequency curve with the fit result $\Omega_{\text{rf}} = (0.086\,77_{\pm 0.000\,18})I_{\text{rf}} + (0.0006_{\pm 0.0006})$. (b) Dependence of current I_{rf} on the rf frequency ω_{rf} under varying signal-source output voltages. The fitting results for signal-source output voltages of 0.2, 0.4, and 0.6 V are $I_{\text{rf}} = (-0.009\,181_{\pm 0.000\,019})\omega_{\text{rf}} + (1.9148_{\pm 0.0019})$, $I_{\text{rf}} = (-0.016\,450_{\pm 0.000\,021})\omega_{\text{rf}} + (3.5778_{\pm 0.0020})$, and $I_{\text{rf}} = (-0.024\,654_{\pm 0.000\,006})\omega_{\text{rf}} + (5.3214_{\pm 0.0006})$, respectively. (c) The I - V curve of the RTN source without frequency shift. The fit result is $I_{\text{RTNI}} = (17.562_{\pm 0.033})U_{\text{RTNI}} + (0.04_{\pm 0.04})$. (d) The I - V curve of the RTN source with frequency shift. The fit result is $I_{\text{RTNI}} = (17.484_{\pm 0.023})U_{\text{RTNI}} + (0.028_{\pm 0.011})$. All fit functions are consistent in units with the axes of the respective plots.

intercepts are not statistically significant, and the coil frequency response is stable, which provides a guarantee for the reliability and repeatability of our magnetic resonance experimental results.

APPENDIX D: THE Ξ -TYPE NIT IN AN ac OPAM

1. Derivation of the effective Hamiltonian

In this configuration, the longitudinal spin is driven by a static bias field B_0 , a coherent ac field $B_1(t) = B_1 \cos \omega_1 t$, and an RTN field $\mu_1(t)$. The transverse spin is driven by a second ac field $B_2(t) = B_2 \cos \omega_2 t$. The two ac magnetic fields are tuned to satisfy the two-photon-resonance condition $\omega_1 + \omega_2 \approx \Omega_0$. In the rotating frame, the Hamiltonian of the system reads

$$H_2 = [\Omega_0 - \omega_2 + \xi_1(t) + \Omega_1 \cos(\omega_1 t)]F_z + \frac{\Omega_2}{2}F_x, \quad (\text{D1})$$

where the Rabi frequencies are defined as $\Omega_j = \mu_B g_F B_j$ ($j = 0, 1, 2$), and $\xi_1(t) = \mu_B g_F \mu_1(t)$. Next, the Hamiltonian $H'' = \omega_1 F_z + \Omega_2 F_x/2$ is diagonalized to bring the entire system into a dressed-state representation. The diagonalized eigenvalue is $\lambda_1 = \sqrt{4\omega_1^2 + \Omega_2^2}/2$. The effective Hamiltonian is

$$H'_2 = [\Delta_z + \xi_z(t)]F_z + \Omega_x F_x, \quad (\text{D2})$$

where the detuning Δ_z is defined as $\Delta_z = \lambda_1 - \Omega_0 + \omega_2$. The RTN coupling strength $\xi_z(t)$ is $\xi_z(t) = (\omega_1/\lambda_1)\xi_1(t)$. The effective detection Rabi frequency is $\Omega_x = -(\Omega_2/2\lambda_1)\Omega_1$. It is worth noting that Eqs. (D2) and (5) are highly consistent in structure and both correspond to the case $\xi_x(t) = 0$ in Eq. (1). This formal isomorphism suggests that, although their physical realization paths differ, their dominant dynamic mechanisms are the same. Therefore,

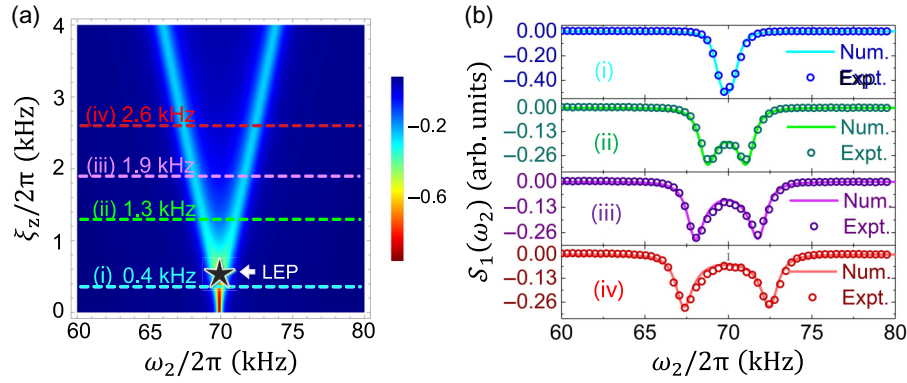


FIG. 9. The longitudinal RTN-induced transparency effect in an ac OPAM. (a) Numerical simulation of the OPAM spectrum $S_1(\omega_2, \xi_z)$ plotted in the ω_2 - ξ_z plane. (b) Theoretical and experimental fits of the OPAM spectra with different ξ_z . Other plot parameters are $\gamma = 3$ kHz, $\Gamma/2\pi \approx 46.4$ Hz, $\Omega_1/2\pi \approx 1.75$ kHz, $\Omega_2/2\pi \approx 1.75$ kHz, $\Omega_0/2\pi = 100$ kHz, and $\omega_1/2\pi = 30$ kHz.

the NIT characteristics and related parameter dependencies that appear in the system corresponding to Eq. (D2) are physically similar to the Ξ -type NIT described by Eq. (5).

2. Experimental results and analysis

Because this configuration is isomorphic to the system described by spectrum (7) at the level of the effective Hamiltonian, the spectral expression retains the same analytical form. Accordingly, the NIT spectrum for the present configuration is obtained by replacing the definitions of Δ_z and Ω_x in spectrum (7) with those in Eq. (D2).

Figure 9 shows the NIT effect triggered by the longitudinal RTN in a two-photon ac OPAM and its spectral variation with the noise amplitude ξ_z . When the RTN amplitude ξ_z is below the LEP, the spectrum shows a single-peak structure and the transparency effect disappears; see Fig. 9(b)(i). However, when the RTN amplitude ξ_z is above the LEP, the spectrum changes from a single-peak structure to a double-peak structure, and the transparent window deepens and broadens with increasing ξ_z ; see Fig. 9(b)(ii)–(iv). This result is consistent with the Ξ -type NIT behavior observed in the dc OPAM, further indicating a unified physical mechanism underlying NIT responses across different configurations.

averaging in a superconducting qubit, *Nat. Commun.* **4**, 1420 (2013).

- [5] J. Shao and P. Hänggi, Decoherent dynamics of a two-level system coupled to a sea of spins, *Phys. Rev. Lett.* **81**, 5710 (1998).
- [6] C.-Y. Hsieh and J. Cao, A unified stochastic formulation of dissipative quantum dynamics. II. Beyond linear response of spin baths, *J. Chem. Phys.* **148**, 014104 (2018).
- [7] J. Ye and P. Zoller, Essay: Quantum sensing with atomic, molecular, and optical platforms for fundamental physics, *Phys. Rev. Lett.* **132**, 190001 (2024).
- [8] Y. Yang, M. Xu, A. Liang, Y. Yin, X. Ma, Y. Gao, and X. Ning, A new wearable multichannel magnetocardiogram system with a SERF atomic magnetometer array, *Sci. Rep.* **11**, 5564 (2021).
- [9] X. Li, Z. Guo, R. Yang, and Y. Feng, Single-beam double-pass miniaturized atomic magnetometer for biomagnetic imaging systems, *IEEE Sens. J.* **23**, 12433 (2023).
- [10] K. Kim, S. Begus, H. Xia, S.-K. Lee, V. Jazbinsek, Z. Trontelj, and M. V. Romalis, Multi-channel atomic magnetometer for magnetoencephalography: A configuration study, *NeuroImage* **89**, 143 (2014).
- [11] W. Tang, G. Huang, G. Li, G. Yang, and X.-X. Geng, Measurement performance improvement method for optically pumped magnetometer based on multilayer perceptron, *IEEE Sensors J.* **24**, 38851 (2024).
- [12] F. Hu, G. Angelatos, S. A. Khan, M. Vives, E. Türeci, L. Bello, G. E. Rowlands, G. J. Ribeill, and H. E. Türeci, Tackling sampling noise in physical systems for machine learning applications: Fundamental limits and eigentasks, *Phys. Rev. X* **13**, 041020 (2023).
- [13] X.-X. Geng, K. Jin, L. Zhou, W.-W. Tang, G. Yang, S. Liang, S.-P. Wu, G.-M. Huang, and G.-X. Li, Spectral observation of symmetry-protected selection rules for dynamical high-dimensional parity in alignment magnetic resonance, *Sci. China Phys. Mech. Astron.* **67**, 100311 (2024).
- [14] D. Budker and D. F. Jackson Kimball, *Optical Magnetometry* (Cambridge University Press, Cambridge, 2013).
- [15] D. Budker and M. Romalis, Optical magnetometry, *Nat. Phys.* **3**, 227 (2007).

- [1] K. Funo and A. Ishizaki, Dynamics of a quantum system interacting with white non-Gaussian baths: Poisson noise master equation, *Phys. Rev. Lett.* **132**, 170402 (2024).
- [2] V. Link, K. Luoma, and W. T. Strunz, Non-Markovian quantum state diffusion for spin environments, *New J. Phys.* **25**, 093006 (2023).
- [3] C. L. Degen, F. Reinhard, and P. Cappellaro, Quantum sensing, *Rev. Mod. Phys.* **89**, 035002 (2017).
- [4] J. Li, M. P. Silveri, K. S. Kumar, J.-M. Pirkkalainen, A. Vepsäläinen, W. C. Chien, J. Tuorila, M. A. Sillanpää, P. J. Hakonen, E. V. Thuneberg, and G. S. Paraoanu, Motional

- [16] H. Yao, B. Maddox, and F. Renzoni, High-sensitivity operation of an unshielded single cell radio-frequency atomic magnetometer, *Opt. Express* **30**, 42015 (2022).
- [17] J. C. Allred, R. N. Lyman, T. W. Kornack, and M. V. Romalis, High-sensitivity atomic magnetometer unaffected by spin-exchange relaxation, *Phys. Rev. Lett.* **89**, 130801 (2002).
- [18] F. Renzoni, Electromagnetic induction imaging with atomic magnetometers: Coming of age, *Quantum Rev. Lett.* **1**, 22 (2025).
- [19] X.-X. Geng, G. Yang, K. Jin, W. Tang, S.-P. Wu, G.-M. Huang, and G.-X. Li, Observation of the influence of generalized parity on the alignment spectra in the periodically modulated warm atomic ensemble, *Laser Photonics Rev.* **16**, 2100614 (2022).
- [20] R. Zhang, W. Xiao, Y. Ding, Y. Feng, X. Peng, L. Shen, C. Sun, T. Wu, Y. Wu, Y. Yang, Z. Zheng, X. Zhang, J. Chen, and H. Guo, Recording brain activities in unshielded Earth's field with optically pumped atomic magnetometers, *Sci. Adv.* **6**, eaba8792 (2020).
- [21] Y. Lu, T. Zhao, W. Zhu, L. Liu, X. Zhuang, G. Fang, and X. Zhang, Recent progress of atomic magnetometers for geomagnetic applications, *Sensors* **23**, 5318 (2023).
- [22] M. Jiang, H. Su, Z. Wu, X. Peng, and D. Budker, Floquet maser, *Sci. Adv.* **7**, eabe0719 (2021).
- [23] H. Su, M. Jiang, Y. Wang, Y. Huang, X. Kang, W. Ji, X. Peng, and D. Budker, New constraints on axion-mediated spin interactions using magnetic amplification, *Phys. Rev. Lett.* **133**, 191801 (2024).
- [24] L. M. Rushton, T. Pyragius, A. Meraki, L. Elson, and K. Jensen, Unshielded portable optically pumped magnetometer for the remote detection of conductive objects using eddy current measurements, *Rev. Sci. Instrum.* **93**, 125103 (2022).
- [25] L. M. Rushton, L. M. Ellis, J. D. Zipfel, P. Bevington, and W. Chalupczak, Performance of a radio-frequency two-photon atomic magnetometer in different magnetic induction measurement geometries, *Sensors* **24**, 6657 (2024).
- [26] J. Picó-Cortés and G. Platero, Dynamical second-order noise sweetspots in resonantly driven spin qubits, *Quantum* **5**, 607 (2021).
- [27] J. H. Prechtel, A. V. Kuhlmann, J. Houel, L. Greuter, A. Ludwig, D. Reuter, A. D. Wieck, and R. J. Warburton, Frequency-stabilized source of single photons from a solid-state qubit, *Phys. Rev. X* **3**, 041006 (2013).
- [28] A. V. Kuhlmann, J. Houel, A. Ludwig, L. Greuter, D. Reuter, A. D. Wieck, M. Poggio, and R. J. Warburton, Charge noise and spin noise in a semiconductor quantum device, *Nat. Phys.* **9**, 570 (2013).
- [29] Y. Masuyama, K. Suzuki, A. Hekizono, M. Iwanami, M. Hatano, T. Iwasaki, and T. Ohshima, Gradiometer using separated diamond quantum magnetometers, *Sensors* **21**, 977 (2021).
- [30] J. Michl, J. Steiner, A. Denisenko, A. Bülau, A. Zimmermann, K. Nakamura, H. Sumiya, S. Onoda, P. Neumann, J. Isoya, and J. Wrachtrup, Robust and accurate electric field sensing with solid state spin ensembles, *Nano Lett.* **19**, 4904 (2019).
- [31] H. Malissa, M. Kavand, D. P. Waters, K. J. van Schooten, P. L. Burn, Z. V. Vardeny, B. Saam, J. M. Lupton, and C. Boehme, Room-temperature coupling between electrical current and nuclear spins in oleds, *Science* **345**, 1487 (2014).
- [32] P. Dutta and P. M. Horn, Low-frequency fluctuations in solids: $1/f$ noise, *Rev. Mod. Phys.* **53**, 497 (1981).
- [33] M. B. Weissman, $1/f$ noise and other slow, nonexponential kinetics in condensed matter, *Rev. Mod. Phys.* **60**, 537 (1988).
- [34] E. Paladino, Y. M. Galperin, G. Falci, and B. L. Altshuler, $1/f$ noise: Implications for solid-state quantum information, *Rev. Mod. Phys.* **86**, 361 (2014).
- [35] K. Hayakawa, S. Kanai, T. Funatsu, J. Igarashi, B. Jinnai, W. A. Borders, H. Ohno, and S. Fukami, Nanosecond random telegraph noise in in-plane magnetic tunnel junctions, *Phys. Rev. Lett.* **126**, 117202 (2021).
- [36] E. Kleinherbers, P. Stegmann, A. Kurzman, M. Geller, A. Lorke, and J. König, Pushing the limits in real-time measurements of quantum dynamics, *Phys. Rev. Lett.* **128**, 087701 (2022).
- [37] B. N. Costanzi and E. D. Dahlberg, Emergent $1/f$ noise in ensembles of random telegraph noise oscillators, *Phys. Rev. Lett.* **119**, 097201 (2017).
- [38] L. Gammaitoni, P. Hänggi, P. Jung, and F. Marchesoni, Stochastic resonance, *Rev. Mod. Phys.* **70**, 223 (1998).
- [39] A. W. Chin, S. F. Huelga, and M. B. Plenio, Coherence and decoherence in biological systems: Principles of noise-assisted transport and the origin of long-lived coherences, *Philos. Trans. R. Soc. A: Math. Phys. Eng. Sci.* **370**, 3638 (2012).
- [40] C. L. Kane, Telegraph noise and fractional statistics in the quantum Hall effect, *Phys. Rev. Lett.* **90**, 226802 (2003).
- [41] E. J. Davis, B. Ye, F. Machado, S. A. Meynell, W. Wu, T. Mittiga, W. Schenken, M. Joos, B. Kobrin, Y. Lyu, Z. Wang, D. Bluvstein, S. Choi, C. Zu, A. C. B. Jayich, and N. Y. Yao, Probing many-body dynamics in a two-dimensional dipolar spin ensemble, *Nat. Phys.* **19**, 836 (2023).
- [42] D. Niepce, J. J. Burnett, M. Kudra, J. H. Cole, and J. Bylander, Stability of superconducting resonators: Motional narrowing and the role of Landau-Zener driving of two-level defects, *Sci. Adv.* **7**, eabh0462 (2021).
- [43] V. V. Mkhitarian, C. Boehme, J. M. Lupton, and M. E. Raikh, Two-photon absorption in a two-level system enabled by noise, *Phys. Rev. B* **100**, 214205 (2019).
- [44] G. Guarnieri, M. Kolář, and R. Filip, Steady-state coherences by composite system-bath interactions, *Phys. Rev. Lett.* **121**, 070401 (2018).
- [45] A. Purkayastha, G. Guarnieri, M. T. Mitchison, R. Filip, and J. Goold, Tunable phonon-induced steady-state coherence in a double-quantum-dot charge qubit, *npj Quantum Inf.* **6**, 27 (2020).
- [46] C. Cao and X. Wang, Noise-assisted quantum autoencoder, *Phys. Rev. Appl.* **15**, 054012 (2021).
- [47] Z.-X. Luo and M. E. Raikh, Landau-Zener transition driven by slow noise, *Phys. Rev. B* **95**, 064305 (2017).
- [48] F. Verstraete, M. M. Wolf, and J. Ignacio Cirac, Quantum computation and quantum-state engineering driven by dissipation, *Nat. Phys.* **5**, 633 (2009).
- [49] G. Yang, H. bin Zhang, X. xing Geng, S. qing Liang, Y. Zhu, J. tao Mao, G. ming Huang, and G. xiang Li, Collision

- induced magnetic resonance in cesium atoms, *Opt. Express* **26**, 30313 (2018).
- [50] X.-X. Geng, G. Yang, P.-L. Qi, W. Tang, S. Liang, G.-X. Li, and G.-M. Huang, Laser-detected magnetic resonance induced by radio-frequency two-photon processes, *Phys. Rev. A* **103**, 053112 (2021).
- [51] B. Maddox and F. Renzoni, Two-photon electromagnetic induction imaging with an atomic magnetometer, *Appl. Phys. Lett.* **122**, 144001 (2023).
- [52] K. Wódkiewicz, B. W. Shore, and J. H. Eberly, Noise in strong laser-atom interactions: Frequency fluctuations and nonexponential correlations, *Phys. Rev. A* **30**, 2390 (1984).
- [53] J. H. Eberly, K. Wódkiewicz, and B. W. Shore, Noise in strong laser-atom interactions: Phase telegraph noise, *Phys. Rev. A* **30**, 2381 (1984).
- [54] P. C. Bressloff, Stochastic Liouville equation for particles driven by dichotomous environmental noise, *Phys. Rev. E* **95**, 012124 (2017).
- [55] F. Campaioli, J. H. Cole, and H. Hapuarachchi, Quantum master equations: Tips and tricks for quantum optics, quantum computing, and beyond, *PRX Quantum* **5**, 020202 (2024).
- [56] W. Xiao, T. Wu, X. Peng, and H. Guo, Atomic spin-exchange collisions in magnetic fields, *Phys. Rev. A* **103**, 043116 (2021).
- [57] M. Auzinsh, D. Budker, and S. Rochester, *Optically Polarized Atoms: Understanding Light-Atom Interactions* (Oxford University Press, New York, 2010).
- [58] D. A. Varshalovich, A. N. Moskalev, and V. K. Khersonskii, *Quantum Theory of Angular Momentum* (World Scientific, Singapore, 1988).
- [59] G. Bevilacqua, E. Breschi, and A. Weis, Steady-state solutions for atomic multipole moments in an arbitrarily oriented static magnetic field, *Phys. Rev. A* **89**, 033406 (2014).
- [60] X. Yang, J. Sheng, and M. Xiao, Electromagnetically induced absorption via incoherent collisions, *Phys. Rev. A* **84**, 043837 (2011).
- [61] S. Wielandy and A. L. Gaeta, Investigation of electromagnetically induced transparency in the strong probe regime, *Phys. Rev. A* **58**, 2500 (1998).
- [62] M. Fleischhauer, A. Imamoglu, and J. P. Marangos, Electromagnetically induced transparency: Optics in coherent media, *Rev. Mod. Phys.* **77**, 633 (2005).
- [63] N. Castagna and A. Weis, Measurement of longitudinal and transverse spin relaxation rates using the ground-state Hanle effect, *Phys. Rev. A* **84**, 053421 (2011).
- [64] W. Happer, Optical pumping, *Rev. Mod. Phys.* **44**, 169 (1972).
- [65] D. Budker, W. Gawlik, D. F. Kimball, S. M. Rochester, V. V. Yashchuk, and A. Weis, Resonant nonlinear magneto-optical effects in atoms, *Rev. Mod. Phys.* **74**, 1153 (2002).
- [66] F. Minganti, A. Miranowicz, R. W. Chhajlany, and F. Nori, Quantum exceptional points of non-Hermitian Hamiltonians and Liouvillians: The effects of quantum jumps, *Phys. Rev. A* **100**, 062131 (2019).
- [67] A. Jebraeili and M. R. Geller, Quantum simulation of a qubit with a non-Hermitian Hamiltonian, *Phys. Rev. A* **111**, 032211 (2025).
- [68] J.-W. Zhang, J.-Q. Zhang, G.-Y. Ding, J.-C. Li, J.-T. Bu, B. Wang, L.-L. Yan, S.-L. Su, L. Chen, F. Nori, S. K. Özdemir, F. Zhou, H. Jing, and M. Feng, Dynamical control of quantum heat engines using exceptional points, *Nat. Commun.* **13**, 6225 (2022).
- [69] W. Horsthemke and R. Lefever, *Noise-Induced Transitions in Physics, Chemistry, and Biology* (Springer, Berlin, 1984).
- [70] P. W. Anderson, A mathematical model for the narrowing of spectral lines by exchange or motion, *J. Phys. Soc. Jpn.* **9**, 316 (1954).
- [71] N. G. Van Kampen, *Stochastic Processes in Physics and Chemistry* (Elsevier Science, New York, 1992), Vol. 1.
- [72] M. O. Scully and M. S. Zubairy, *Quantum Optics* (Cambridge University Press, Cambridge, 1997).
- [73] K.-J. Boller, A. Imamoglu, and S. E. Harris, Observation of electromagnetically induced transparency, *Phys. Rev. Lett.* **66**, 2593 (1991).
- [74] M. Noaman, D. W. Booth, and J. P. Shaffer, Rydberg-atom sensors in bichromatic radio-frequency fields, *Phys. Rev. Appl.* **20**, 024068 (2023).
- [75] Z. Chen, B. Liu, and G. Guo, Adaptive single photon detection under fluctuating background noise, *Opt. Express* **28**, 30199 (2020).
- [76] C. Müller, J. Lisenfeld, A. Shnirman, and S. Poletto, Interacting two-level defects as sources of fluctuating high-frequency noise in superconducting circuits, *Phys. Rev. B* **92**, 035442 (2015).
- [77] C. Müller, J. H. Cole, and J. Lisenfeld, Towards understanding two-level-systems in amorphous solids: Insights from quantum circuits, *Rep. Progr. Phys.* **82**, 124501 (2019).
- [78] N. Ratanov and A. D. Kolesnik, *Telegraph Processes and Option Pricing* (Springer, Berlin, 2022).
- [79] A. Weis, G. Bison, and A. S. Pazgalev, Theory of double resonance magnetometers based on atomic alignment, *Phys. Rev. A* **74**, 033401 (2006).
- [80] M. Matsumoto and T. Nishimura, Mersenne twister: A 623-dimensionally equidistributed uniform pseudo-random number generator, *ACM Trans. Model. Comput. Simul.* **8**, 3 (1998).