

Collective quantum stochastic resonance in Rydberg atoms

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(Received 4 February 2024; accepted 27 August 2024; published 21 November 2024)

We study the collective response of a group of dissipative Rydberg atoms to a periodic modulation of the Rydberg excitation laser. Focusing on the emergent collective-jump dynamics, where the system stochastically switches between states with distinct Rydberg excitations, we show that the counting statistics of the state switching is qualitatively changed by the periodic drive. The impact is most prominent when the driving frequency is comparable to the emergent collective-jump rate, as the jumps tend to synchronize with the external drive, and their counting statistics exhibits a series of suppressed subharmonics of the driving frequency. These phenomena are manifestations of a novel type of stochastic resonance, where a cooperative collective state switching is facilitated by quantum fluctuations in a many-body open system. Such a collective quantum stochastic resonance further leads to an enhanced signal-to-noise ratio in the power spectrum of the Rydberg excitations, for which the synchronized collective jumps are viewed as the output signal. We confirm the many-body quantum nature of the resonance by devising a cluster model, under which the role of many-body correlations is analyzed by changing the size of the atom clusters.

DOI: [10.1103/PhysRevResearch.6.L042046](https://doi.org/10.1103/PhysRevResearch.6.L042046)

Introduction. Stochastic resonance describes the counter-intuitive signal amplification by adding an appropriate dose of noise [1,2]. Originally suggested to explain the recurring ice ages [3], stochastic resonance proves to be a universal phenomenon in nature, appearing also in biological processes [4,5], and across a variety of physical systems ranging from optics [6–11] to electronics [12,13]. The occurrence of stochastic resonance generally entails a noisy bistable system periodically modulated by a weak signal. When the frequency of the driving signal and the noise strength are appropriate, the response of the system exhibits resonancelike behaviors, including the synchronization of the bistable-state switching with the external drive, and an enhanced signal-to-noise ratio (SNR). While stochastic resonances are predominantly found in nonlinear systems with classical noise, they also emerge in quantum mechanical settings, assisted by quantum fluctuations [14–18]. These quantum stochastic resonances have only recently been observed in quantum-dot microstructures [19] and scanning tunneling microscopes [20], where the intrinsic fluctuations of single-electron tunneling play a key role.

But quantum fluctuations, particularly those in many-body open systems, also manifest in collective dynamics. Take a group of dissipative Rydberg atoms for instance. The atomic

ground states therein are laser coupled to high-lying energy levels with significant principle quantum number n , but subject to spontaneous decay back to the ground state. An outstanding feature of the Rydberg atoms is the strong long-range interactions, which are central to the recent advances in quantum computation and simulation with Rydberg atoms [21–29]. Previous studies have shown that the entanglement generated by the long-range interactions, combined with the repetitive quantum measurements (in the form of spontaneous emission), lead to collective quantum jumps between two metastable many-body states with high and low Rydberg excitations, respectively [30]. Since these emergent collective jumps are a manifestation of quantum fluctuations in a many-body open system, it is tempting to ask whether a quantum stochastic resonance can be engineered based on them.

In this work, we reveal such a quantum stochastic resonance in the counting statistics of the collective jumps in a group of Rydberg atoms, where the Rydberg excitation laser is subject to a periodic modulation. While rich resonance behaviors emerge in the stochastic state evolution of the many-body open system, many-body correlations are indispensable for the observed resonance, in contrast to previously studied stochastic resonances. We explicitly demonstrate this using a cluster model, where quantum coherence and entanglement are only retained within each atom cluster. Specifically, the cluster size and partitions dramatically impact the collective jumps and resonance condition.

Emergent collective jumps in Rydberg atoms. We consider a group of N atoms uniformly subject to a Rydberg excitation laser. As illustrated in Fig. 1(a), the ground state of the j th atom $|g_j\rangle$ is continuously coupled to a Rydberg state $|r_j\rangle$. The Rydberg state has a finite lifetime, which can be further

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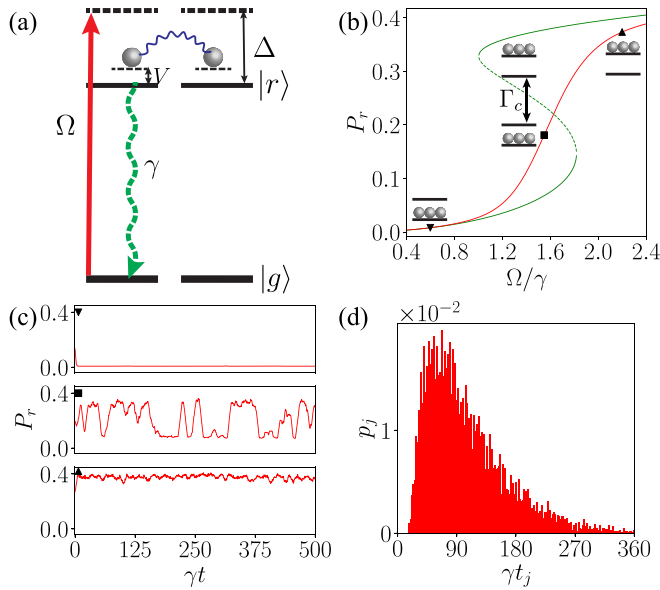


FIG. 1. (a) Schematic illustration of the level scheme. The Rydberg excitation laser couples the atomic ground state $|g\rangle$ and a Rydberg state $|r\rangle$. Atoms in the excited state interact with each other through the Rydberg long-range interaction, and can spontaneously decay back to the ground state. (b) Rydberg-state population P_r in the steady state of the master equation (red), and in the bistable states of the optical Bloch equations (green). Note that solid (dashed) sectors of the green curve correspond to stable (unstable) solutions. Collective jumps with rate Γ_c (defined in the main text) predominantly occur in the mean-field bistable region, but only visible through quantum trajectories. Here the mean-field bistable region exists in the range $\Omega/\gamma \in [1.00, 1.82]$. (c) Rydberg excitation along a single trajectory. From top to bottom: $\Omega/\gamma = 0.6, 1.55, 2.2$, consistent with the corresponding black markers in (b). (d) Normalized distribution p_j of interjump intervals, indicative of the counting statistics of the collective jumps in the middle panel of (c). We take $\gamma\delta t = 2$ for the definition of p_j and t_j (see main text for details). For numerical calculations, we take $\Delta/\gamma = 3.4$, $(N-1)V/\gamma = 10$, and $N = 8$ in all figures.

tuned through additional couplings to other excited states. Approximating the Rydberg atoms as an ensemble of two-level systems and neglecting the much slower external motion, the dissipative dynamics of the density matrix ρ is depicted by the Lindblad master equation

$$\dot{\rho} = -i[H, \rho] + \gamma \sum_j \left(-\frac{1}{2}\{L_j^\dagger L_j, \rho\} + L_j \rho L_j^\dagger \right), \quad (1)$$

with the Hamiltonian

$$H = \sum_j \left[-\frac{\Delta}{2}(\sigma_j^z + 1) + \frac{\Omega}{2}\sigma_j^x \right] + \frac{V}{4} \sum_{j < k} (\sigma_j^z + 1)(\sigma_k^z + 1). \quad (2)$$

The quantum jump operator $L_j = (\sigma_j^x - i\sigma_j^y)/2$ describes the spontaneous decay of the j th atom back to its ground state, with a decay rate γ . Here Ω is the Rabi frequency of the Rydberg excitation laser, Δ is its detuning, and $\sigma_j^{x,y,z}$ are the Pauli operators associated with the two-dimensional subspace spanned by $\{|g_j\rangle, |r_j\rangle\}$ of the j th atom, with

$\sigma_j^z = |r_j\rangle\langle r_j| - |g_j\rangle\langle g_j|$. For simplicity, we model the long-range Rydberg interactions using a constant all-to-all coupling V . Note that all the key results still hold when we consider the spatial dependence of the Rydberg interactions [31].

The driven-dissipative Rydberg system features a unique steady state (with density matrix ρ^s), which can be solved by setting $\dot{\rho}^s = 0$ in Eq. (1). In Fig. 1(b), we show the steady-state Rydberg excitation $P_r = \rho_{rr}^s$ for $N = 8$ atoms (red). For later discussions, it is instructive to compare the steady-state solution from the master equation with those under a mean-field approximation. In the latter case, the interatomic correlations are neglected, the density matrices of individual atoms become decoupled through $\rho = \bigotimes_j \bar{\rho}$, and the dynamics of $\bar{\rho}$ is governed by the optical Bloch equations [30,32]

$$\dot{\bar{\rho}}_{rr} = -\Omega \text{Im} \bar{\rho}_{rg} - \gamma \bar{\rho}_{rr}, \quad (3)$$

$$\dot{\bar{\rho}}_{rg} = i[\Delta - (N-1)V\bar{\rho}_{rr}]\bar{\rho}_{rg} - \frac{\gamma}{2}\bar{\rho}_{rg} + i\Omega\left(\bar{\rho}_{rr} - \frac{1}{2}\right), \quad (4)$$

where the Rydberg interactions contribute toward a density-dependent effective detuning. Steady-state solutions can be obtained by setting $\dot{\bar{\rho}} = 0$ in the equations above, giving rise to the green lines in Fig. 1(b). Notably, a bistable region emerges, where the two steady-state solutions feature high and low Rydberg excitations, respectively. The bistability derives from the nonlinear interaction terms under the mean-field approximation, and is well studied in thermal Rydberg gases [33–35]. While such a bistability is absent in the full-quantum description of Eq. (1), its quantum correspondence nevertheless manifests in the counting statistics of the open-system dynamics.

To see this, we unravel the density-matrix dynamics into an ensemble of quantum trajectories [36–38]. In a given time interval $[t, t + \Delta t]$ along each trajectory (with $\gamma\Delta t \ll 1$), the many-body wave function $|\psi(t)\rangle$ undergoes the quantum jump $|\psi(t)\rangle \rightarrow L_j|\psi(t)\rangle$ with a probability $\gamma\Delta t \langle \psi(t) | L_j^\dagger L_j | \psi(t) \rangle$, or evolves under a non-Hermitian effective Hamiltonian $H_{\text{eff}} = H - i\frac{\gamma}{2} \sum_j L_j^\dagger L_j$ with the probability $1 - \gamma\Delta t \sum_j \langle \psi(t) | L_j^\dagger L_j | \psi(t) \rangle$. The state is then normalized to repeat the process for the following time interval. Given our unraveling scheme above, each single trajectory simulates the many-body wave-function evolution in a single-run experiment, provided the spontaneously emitted photons are directly detected [31,39–41].

An important feature of an individual quantum trajectory of the system is the possible emergence of collective jumps between metastable many-body states with high and low Rydberg excitations [30]. The two states are characterized by the joint emissions or nonemissions of individual atoms, and the collective jumps between them originate from the interplay of Rydberg interaction and dissipation. In Fig. 1(c), we show the single-trajectory evolution of the Rydberg excitation $P_r = \frac{1}{2} \langle \psi(t) | \sum_j (\sigma_j^z + 1) | \psi(t) \rangle$ with different Ω . Apparently, when the parameters are within the mean-field bistable region [Fig. 1(c) middle panel], P_r switches between two plateaus, whose values correspond to the bistable steady-state solutions under the mean-field approximation. When the system is tuned away from the bistable region, the collective jumps dramatically decrease in

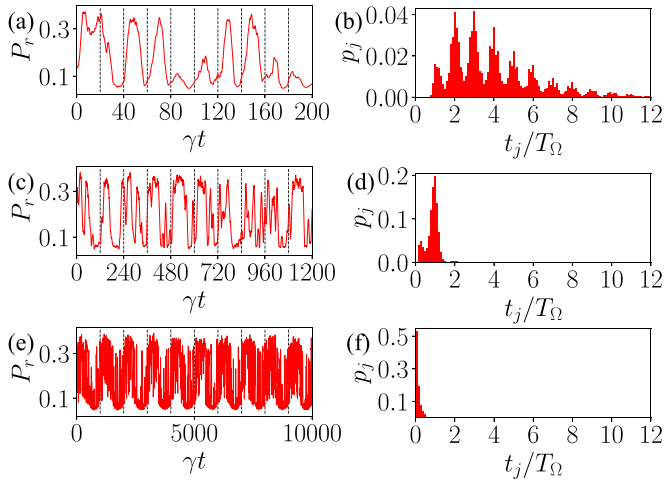


FIG. 2. (a), (c), (e) Single-trajectory dynamics of the Rydberg excitation P_r under a periodic drive. (b), (d), (f) Distribution of inter-jump intervals for a single trajectory within the evolution time $\gamma t \in [0, 10^6]$. The driving parameters are $\Omega_0/\gamma = 1.55$, $A_\Omega/\gamma = 0.2$, and (a), (b) $\gamma T_\Omega = 20$, (c), (d) $\gamma T_\Omega = 120$, (e)(f) $\gamma T_\Omega = 1000$. For all figures, we take $N = 8$ and $\delta t = 0.17T_\Omega$. Other parameters are the same as those in Fig. 1.

frequency, and eventually vanish in effect [Fig. 1(c) top and bottom panels]. This is the case when the steady-state solution of the master equation is close to one of the steady-state solutions of the optical Bloch equations.

To characterize the counting statistics of the collective jumps, we plot the distribution of interjump intervals in Fig. 1(d). We uniformly discretize the overall evolution time T into small segments of δt , and count the number of adjacent downward collective jumps n_j that occur within the time interval $[t_j - 0.5\delta t, t_j + 0.5\delta t]$, where $t_j = j\delta t$. Here the downward collective jumps are the state switching from high to low Rydberg excitations. The time of occurrence of these downward jumps along a quantum trajectory is determined numerically by setting a population-difference threshold $P_r^{\text{th}} = 0.25$ [31]. The normalized distribution for the interjump interval is defined as $p_j = n_j / \sum_j n_j$, which represents the probability for adjacent downward collective jumps to occur within the corresponding time interval. The distribution shows an exponential decay for large j , suggesting the stochastic nature of the jumps. We further define a collective-jump rate $\Gamma_c = \sum_j n_j / T$, which represents an emergent many-body energy scale for the driven-dissipative system.

Counting statistics and quantum stochastic resonance. We now consider the counting statistics of the collective jumps in response to a periodic modulation of the Rydberg excitation laser, with $\Omega(t) = \Omega_0 + A_\Omega \cos(2\pi t/T_\Omega)$, where T_Ω and A_Ω are, respectively, the driving period and amplitude, and Ω_0 is the time-independent Rabi frequency.

In Figs. 2(a), 2(c), 2(e), we show the time-evolved Rydberg excitations of single trajectories under different driving frequencies. The collective jumps appear to be regulated by the driving frequency, distinct from the more random pattern in Fig. 1(c). Importantly, at an appropriate intermediate driving frequency, the collective jumps fully synchronize with the

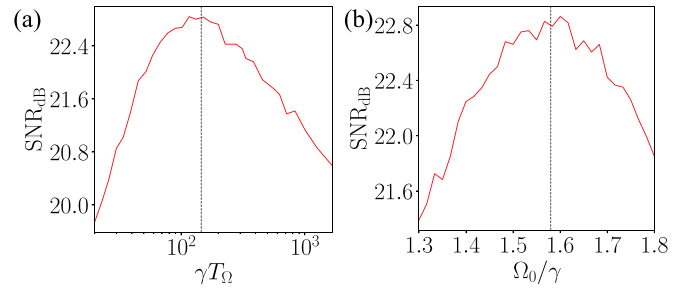


FIG. 3. (a) Trajectory-averaged SNR, with a fixed $\Omega_0/\gamma = 1.55$, $A_\Omega/\gamma = 0.2$, and varying T_Ω . (b) SNR as a function of Ω_0 with fixed $T_\Omega = 144.2$ and $A_\Omega/\gamma = 0.2$. The vertical dashed lines in both subplots indicate the theoretically predicted resonance condition $T_c \tilde{\Gamma}_c = 1$, which gives $\gamma T_c \approx 145.4$ in (a), and $\Omega_0/\gamma \approx 1.58$ in (b). Here $\tilde{\Gamma}_c$ is the downward collective-jump rate that is numerically averaged along the driving path [31]. For both calculations here, we take a single-trajectory evolution time of $\gamma t = 10000$, and average over 100 trajectories. Other parameters are the same as those in Fig. 1.

drive [see Fig. 2(c)]. Such a pattern is more visible from the counting statistics of the jumps, shown in Figs. 2(b), 2(d), 2(f). When the driving period is small, the response of the system, in terms of the downward collective-jump intervals, features a series of subharmonics of the driving frequency [see Fig. 2(b)]. On increasing the driving period, the subharmonics become suppressed, and the distribution is peaked at $t_j/T_\Omega = 1$, indicating full synchronization [see Fig. 2(d)]. Under a very long driving period, more than one jump can occur within one driving period, and p_j is peaked at $t_j/T_\Omega = 0$ [see Fig. 2(f)]. In all cases, the distribution of p_j exhibits discrete peaks centered at integer t_j/T_Ω , in sharp contrast to the pattern in Fig. 1(d).

The synchronization of the system's response with the periodic drive is reminiscent of the stochastic resonance, where the resonance condition is qualitatively understood as the matching between the external driving frequency and the rate of the emergent collective jumps [1, 19, 42]. To confirm such an observation, we calculate the SNR of the collective jumps, which is an outstanding quantifier for the stochastic resonance. We view the period modulation as a weak input signal, and the collective jumps the output. The SNR can be extracted from the power spectrum $P(f)$ of the collective jumps in a single trajectory, where $P(f)$ is the square of the Fourier transform of $P_r(t)$. The SNR is defined as [2]

$$\text{SNR}_{\text{dB}} = 10 \log_{10} \left[\frac{P(f = \frac{1}{T_\Omega})}{\frac{1}{f_{\text{max}} - f_{\text{min}}} \int_{f_{\text{min}}}^{f_{\text{max}}} P(f) df} \right], \quad (5)$$

where the frequency range $[f_{\text{min}}, f_{\text{max}}]$ of the noise-power integral is typically taken close to but not over the driving frequency. To compensate for the arbitrariness in the choice of $f_{\text{min}, \text{max}}$, at any given point, we take three different sets of ranges: $[f_{\text{min}}, f_{\text{max}}] \in (1/T_\Omega)\{[0.25, 0.75], [1.25, 1.75], [2.25, 2.75]\}$, and retain the one with the largest integrated noise power.

We plot the numerically evaluated SNR in Fig. 3(a), as a function of the driving period, where the SNR is further

averaged over many trajectories for better convergence within the numerically feasible evolution time. The resonance point is identified as the peak location in the SNR, which is close to the theoretically predicted resonance point $T_c \tilde{\Gamma}_c = 1$ [1,42] (vertical dashed line). Here $\tilde{\Gamma}_c$ is estimated using the averaged collective-jump rate along the driving path $\Omega(t)$ [31]. The general behavior of the SNR can be understood in view of the distribution of Figs. 2(b), 2(d), 2(f). Specifically, as the driving period increases, the subharmonics in the distribution become suppressed, and the output signal (the peak near $t_j/T_\Omega = 1$) increases in weight in the response. The SNR peak is reached when the collective jumps are fully synchronized with the drive, as a dominant peak emerges near $t_j/T_\Omega = 1$. When the driving frequency is tuned past the resonance, multiple collective jumps can occur within a single driving period, leading to a deterioration of the signal and the SNR. Similar SNR behavior can be observed by fixing the driving period and tuning Ω_0 [see Fig. 3(b)]. This is equivalent to changing the intrinsic collective-jump rate to meet the resonance condition (indicated by the vertical dashed line). As such, emergent collective jumps assist the periodic drive to enhance the SNR in an open system, which is but the hallmark of the stochastic resonance.

Impact of many-body correlations. To understand the many-body nature of the collective quantum stochastic resonance, we consider a cluster model where interactions are treated exactly within the finite-size clusters, while mean-field approximations are adopted for intercluster interactions. For simplicity, we assume all clusters to have the same number of atoms M (dubbed the cluster size).

The overall density matrix can then be formally written as a direct product $\rho = \bigotimes_n \rho_n$, where ρ_n is the density matrix of the n th cluster. Apparently, quantum coherence and entanglement only exist within the same cluster. While the master equation is formally the same as Eq. (1), the Hamiltonian of the n th cluster is

$$H_n = \sum_{j=1}^M \left[-\frac{\Delta}{2} (\sigma_{n,j}^z + 1) + \frac{\Omega}{2} \sigma_{n,j}^x \right] + \frac{V}{4} \sum_{j < k} (\sigma_{n,j}^z + 1) (\sigma_{n,k}^z + 1) + \frac{V}{4} \left(\sum_{m \neq n} P_r^m \right) \sum_{j=1}^M (\sigma_{n,j}^z + 1), \quad (6)$$

where $\sigma_{n,j}^{x,y,z}$ are the Pauli operators for the j th atom in the n th cluster, and the Rydberg excitation of the m th cluster is define as $P_r^m = \frac{1}{2} (\sum_{j=1}^M (\sigma_{m,j}^z + 1))$. Such a construction allows us to examine the impact of many-body correlations by decreasing the cluster size M , as the system is sequentially broken down into local clusters.

While the cluster model recovers the full-quantum description of Eq. (1) for $M = N$, the collective jumps are dramatically modified when $M < N$. Particularly, in the absence of the periodic modulation, the intrinsic collective-jump rate Γ_c significantly decreases for smaller M , as illustrated in Fig. 4(a). We also show the single-trajectory P_r for different cluster sizes in Figs. 4(b), 4(c), 4(d), where the slowdown of

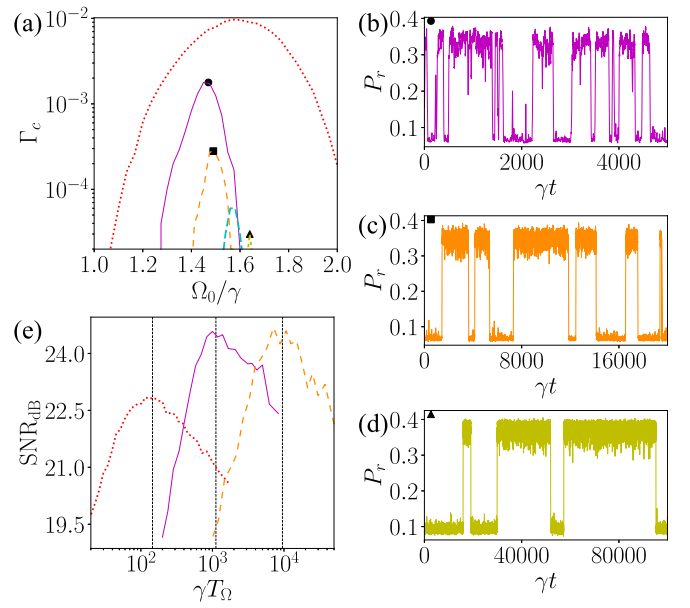


FIG. 4. (a) Intrinsic downward collective-jump rates as functions of Ω_0 under the cluster model in the absence of the periodic drive. The calculations are for a system of $N = 8$ atoms, with $M = 8$ (red), $M = 4$ (purple), $M = 2$ (blue), and $M = 1$ (yellow), respectively. The orange dashed curve corresponds to a partition of three clusters, with 3, 3, and 2 atoms, respectively [dubbed the (3,3,2) configuration]. (b), (c), (d) Single-trajectory dynamics for (b) $M = 4$ at $\Omega_0/\gamma = 1.47$, (c) the (3,3,2) configuration at $\Omega_0/\gamma = 1.49$, and (d) $M = 1$ at $\Omega_0/\gamma = 1.64$. The parameter Ω_0 is chosen at the position with the maximum Γ_c under any given M . (e) SNRs under the cluster model of $M = 8$ (red dotted), $M = 4$ (purple solid) and the (3,3,2) configuration (orange dashed). For $M = 4$ [the (3,3,2) configuration], we fix $\Omega_0/\gamma = 1.47$ ($\Omega_0/\gamma = 1.49$), $A_\Omega/\gamma = 0.1$ ($A_\Omega/\gamma = 0.08$), and average the SNR over 35 (12) trajectories, where the evolution time for a single trajectory is $\gamma t = 50000$ ($\gamma t = 300000$). The vertical dashed lines indicates the theoretical resonance points, with $\gamma T_c \approx 1118.4$ and $\gamma T_c \approx 9525.2$ for the two cases, respectively. Other parameters are the same as those in Fig. 1.

the collective-jump rate with decreasing M is directly visible. This suggests that the resonance point should also shift toward lower frequencies at smaller M . We confirm this using cluster models of smaller cluster partitions in Fig. 4(e). While the SNR peaks are consistent with theoretical predictions, they appear at larger T_Ω than the single-cluster case with $M = 8$.

Finally, we note that, in principle, the stochastic resonance persists even in the $M = 1$ case, albeit at a very small resonance frequency due to the scarcity of collective jumps. However, the quantum stochastic resonance does not exist under the optical Bloch equations of Eqs. (3) and (4), due to the complete lack of individual quantum jumps therein. Specifically, a key difference between the $M = 1$ cluster model and the full mean-field description is that, in the latter case, all the atoms are in the same state and described by the same density matrix. This renders sequential atomic emissions impossible, which is crucial for the buildup of a collective jump [30].

Discussion. The collective quantum stochastic resonance is readily observable in arrays of Rydberg atoms confined to

optical tweezers, or in Rydberg gases at low temperatures. Take the Rydberg state $|70S_{1/2}\rangle$ of ^{87}Rb in a tweezer as an example. Under typical experimental parameters [24,25], the parameters of Fig. 3(a) correspond to $\gamma = 2\pi \times 1$ MHz, $\Delta \sim 2\pi \times 3.4$ MHz. Note that the decay rate γ is larger than the typical spontaneous decay of Rydberg states, but can be achieved by laser coupling the Rydberg state to intermediate low-lying states. The resonance then occurs at $T_\Omega \sim 22.95$ μs for $N = 8$ atoms with an interatomic distance ~ 9.19 μm (where $V \sim 2\pi \times 1.43$ MHz). As the temperature rises, the collective-jump rate decreases along with the size of local coherent clusters, and it would become more difficult to observe the resonance behavior. For thermal Rydberg gases whose dynamics are well captured by the nonlinear optical Bloch equations [33–35], neither the collective jump nor the

quantum stochastic resonance persists. Nevertheless, based on the intrinsic non-linearity-induced bistability, stochastic resonances can still emerge when some form of classical noise is introduced [31]. The subsequent resonance behaviors would be similar in nature to those of conventional stochastic resonances, or related phenomena in nonlinear non-Hermitian systems [43], but are fundamentally different from what we study here. Finally, while we focus on a finite-size system, further analysis suggests that the resonance behavior becomes easier to detect with larger N [31].

Acknowledgments. We thank Zongkai Liu for helpful discussions. This work is supported by the National Natural Science Foundation of China (Grant No. 12374479) and the Innovation Program for Quantum Science and Technology (Grants No. 2021ZD0301904 and No. 2021ZD0301200).

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