Decoherence Cancellation through Noise Interference

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(Received 11 June 2025; revised 16 September 2025; accepted 8 January 2026; published 27 January 2026)

We propose a novel, feedback-free method to cancel the effects of dephasing in the dynamics of open quantum systems. The protocol makes use of the coupling with an auxiliary system when they are both subjected to the same noisy dynamics in such a way that their interaction leads to cancellation of the noise on the system itself. This requires tuning the strength of the coupling between the main and auxiliary systems as well as the ability to prepare the auxiliary system in a Fock state, which solely depends on the coupling strength. We investigate the protocol's efficiency to protect NOON states against dephasing in setups such as tweezer arrays of cold atoms. We show that the protocol's efficiency is robust against fluctuations of the optimal parameters and, remarkably, that it is independent of the temporal features of the noise. Therefore, it could be applied to cancel both Markovian and non-Markovian dephasing noise, reaching regimes where error-correction protocols become inefficient.

DOI: 10.1103/251j-cnbp

Loss of quantum coherence (decoherence) arises from the unavoidable interactions of quantum systems with their environments [1,2] and is an outstanding challenge in the quest for next-generation quantum devices [3]. A severe limitation is pure dephasing, which suppresses phase coherence and can be described either microscopically by coupling to a quantum bath or by stochastic fluctuations of Hamiltonian parameters [4]. A range of approaches has been proposed to improve the stability and reliability of quantum dynamics, thereby bringing the state of the art closer to truly practical and scalable quantum technologies. These approaches include dynamical decoupling protocols, namely, open-loop quantum control techniques that are provably effective for Markovian processes [5–8]. Other approaches make use of weak and strong measurements, with success rates limited by the failure rates of the measurement processes involved [9,10]. Recently, protocols based on spectator qubits have been implemented that perform correction of phase errors during the circuit dynamics by measuring the states of qubits not involved in the quantum dynamics but subjected to noise correlated with that affecting the computational qubits [11]. By combining measurements, data processing, and feedforward operations, correlated errors were suppressed during the execution of the quantum circuit. Similar ideas are at the basis of erasure conversion protocols in tweezer arrays of Rydberg atoms [12,13]. The use of ancillary systems for robust quantum state preparation and computing is at the basis of a series of protocols, where entanglement between system and ancilla followed by measurements of the ancilla's state, or decay of the ancilla, can effectively pull the system into a fixed point of the composite dynamics [14–21]. An alternative ansatz is to prepare states that are immune from specific types of decoherence, "dark states" belonging to decoherence-free subspaces, which are symmetry-protected from the external environment [22–25]. This approach requires the design of appropriate coupling depending on the target state and on the environment [25,26], and is typically effective for memoryless, Markovian reservoirs. A different class of schemes crucially rely on measurement and postselection. Such schemes have been implemented for quantum state preparation; see, e.g., Refs. [27,28]. In quantum communication, it has been argued that repeated teleportation steps can probabilistically reverse or mitigate decoherence [29]. In [30] it has been shown that hybrid multipartite entangled resources can be tailored to increase the robustness of teleportation against pure decoherence.

In this Letter, we propose a protocol that achieves exact cancellation of a correlated pure-dephasing channel for arbitrary initial states of the system by means of quantum interference. The strategy is illustrated in Fig. 1. It requires controlling the coupling between computational and

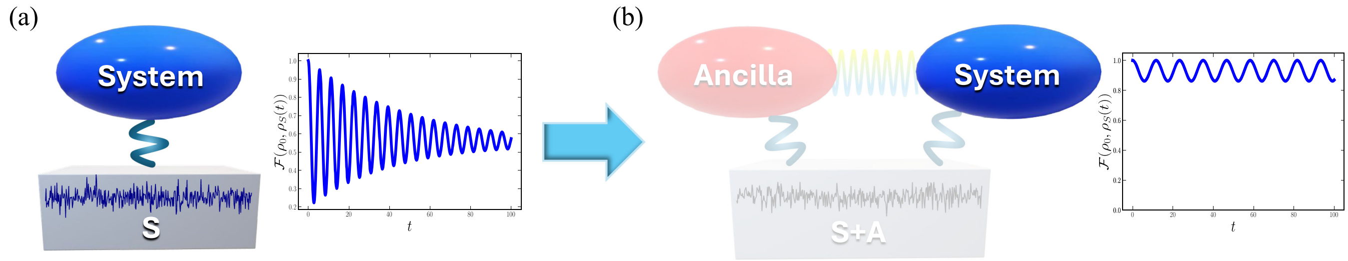

FIG. 1. Schematic illustration of the proposed decoherence cancellation strategy. (a) The target system S is initially subject to decoherence. (b) Upon enabling the interaction between S and A by bringing the ancilla into close proximity with the target system, both subsystems are exposed to a common correlated dephasing noise. Through appropriate preparation of the ancilla's state and precise tuning of the interaction, a dark state is formed in the combined $S+A$ system, thereby canceling decoherence in S and preserving only its unitary dynamics after tracing out the ancilla.

spectator qubits that we refer to as system S and ancilla A. By properly engineering the coupling between S and A, the individual decoherence processes affecting each subsystem no longer simply add; instead, they can interfere destructively, yielding noise cancellation. We argue that the scheme is efficient for arbitrary input states of the system within a class of correlated-noise models. The scheme does not require feedback operations, intermediate measurements, symmetry constraints, or repeated protocol rounds. Most importantly, the cancellation condition is independent of the temporal noise spectrum (Markovian or non-Markovian), as long as the noise of system and ancilla are correlated.

To illustrate the procedure, consider a quantum system in a d-dimensional Hilbert space with Hamiltonian

$$H_S = \eta_S (J_S^z)^2 - \gamma_S J_S^x - \Delta_S J_S^z, \tag{1}$$

where η_S, γ_S, Δ_S are real scalars and $J_S^{x,y,z}$ satisfy angular-momentum commutation relations, $[J_S^x, J_S^y] = iJ_S^z$ and $(J_S^x)^2 + (J_S^y)^2 + (J_S^z)^2 = (d^2 - 1/4)$. Let $|\ell\rangle_S$ be eigenstates of J_S^z: $J_S^z|\ell\rangle_S = \ell|\ell\rangle_S$, $-(d-1/2) \le \ell \le (d-1/2)$. Random fluctuations in the parameter Δ_S give rise to decoherence effects that suppress off-diagonal terms with respect to the chosen basis. In the Markovian limit, after averaging over the noise, the dephasing effects are described by an irreversible, Markovian dynamics for the system density matrices $\rho_S(t)$ generated by a master equation of Gorini-Kossakowski-Sudarshan-Lindblad (GKSL) form [31],

$$\partial_t \rho_S(t) = -i\left[H_S, \rho_S(t)\right] + \lambda \mathbb{M}[\hat{J}_S^z]\rho_S(t), \tag{2}$$

where $\mathbb{M}[\hat{O}]\hat{\rho} = \hat{O}\,\hat{\rho}\,\hat{O}^\dagger - \{\hat{O}^\dagger \hat{O}, \hat{\rho}\}/2$, and λ is the dimensionless decoherence rate. In this Letter, however, we consider a generic type of noise, which includes Markovian noise as a specific white-noise limit. For our discussion, we thus refer to a stochastic Hamiltonian,

$$H'_S(t) = H_S + w_S(t) J_S^z, \tag{3}$$

with $w_S(t)$ a generic centered stochastic potential, with zero mean, $\langle\!\langle w_S(t) \rangle\!\rangle = 0$, and generic higher moments.

In order to mitigate decoherence, the system S is coupled to an ancilla A. Here, for simplicity, we assume that it is also a d-level quantum system with Hamiltonian

$$H_A = \eta_A (J_A^z)^2 - \Delta_A J_A^z. \tag{4}$$

In working experimental conditions, the ancilla is also affected by the same noise as the target system. Assume the SA interaction has the form $H_{SA}^I = \alpha J_S^z \otimes J_A^z$, where $\alpha \in \mathbb{R}$ scales the strength. Let us suppose that, by switching on the interaction, the noise affects the compound system $S+A$ in a correlated way as described by the stochastic Hamiltonian $H_{SA}(t) = H_0 + H_{\mathrm{stoch}}(t)$, where $H_0 \coloneqq H_S \otimes \mathbb{I}_A + \mathbb{I}_S \otimes H_A + H_{SA}^I$ is a deterministic contribution and the stochastic term is given by

$$H_{\mathrm{stoch}}(t) = w_S(t) J_S^z \otimes \mathbb{I}_A + w_A(t) \mathbb{I}_S \otimes J_A^z + w_{SA}(t) \alpha J_S^z \otimes J_A^z, \tag{5}$$

where α is a time-independent parameter.

In the case of a white noise [as shown in Sec. I of Supplemental Material (SM) [32]], one obtains a master equation of GKSL type. The same effective description can be derived by coupling S and the ancilla A to a truly quantum environment consisting of a Fermionic bath in equilibrium at high temperature as in the case in the so-called "singular coupling limit" [33,34]. More in general, that is, in the presence of generic quantum baths inducing dephasing, the core mechanism of our proposal can be described as follows. One extends the Hilbert space to include the bath degrees of freedom in a Hamiltonian of the form

$$H_{SAB} = J_S^z \otimes \mathbb{I}_A \otimes X_B^S + \beta \mathbb{I}_S \otimes J_A^z \otimes X_B^A + \alpha J_S^z \otimes J_A^z \otimes X_B^{SA}, \tag{6}$$

where X_B^j, $j = S, A, SA$, are bath operators and the parameters β, α are assumed to be, at least partially, controllable, as we specify below. These operators can describe, for instance, the magnetic distortion of the vibrations of a bulk in which the spins of system and ancilla are embedded [35,36]. This example encompasses also Rydberg atoms, where the vibrations can be tweezer vibrations and the trapping potential is spin-dependent [37]. The coupling with the ancillary defect can be controlled by tuning the transition frequency by means of external fields. In other settings, terms with this structure can occur in cavity QED setups [38,39], where S is the atomic motion, A the cavity field, and the noise arises from spontaneous decay [40,41]. The coupling parameters are controlled by means of a laser driving the atoms. Let us assume one can prepare the auxiliary system in a right eigenstate of J_A^z, say, $J_A^z|\ell\rangle_A = \ell|\ell\rangle_A$. Consider then any state $|\chi\rangle = |\Psi\rangle_{SA} \otimes |\mathcal{B}\rangle$, where $|\Psi\rangle_{SA} = |\psi\rangle_S \otimes |\ell\rangle_A$ and $|\mathcal{B}\rangle$ is an arbitrary state of the bath. Under the action of the Hamiltonian in (6), this class of states satisfies the relation

$$H_{SA\mathcal{B}}|\chi\rangle = J_S^z \otimes \mathbb{I}_A \otimes (X_{\mathcal{B}}^S + \alpha\ell X_{\mathcal{B}}^{SA})|\chi\rangle + |\chi\rangle_A, \quad (7)$$

where $|\chi\rangle_A = \ell\beta\mathbb{I}_S \otimes \mathbb{I}_A \otimes X_{\mathcal{B}}^A|\chi\rangle$. Consider the bath correlation functions with respect to a suitable bath equilibrium state $\rho_{\mathcal{B}}$,

$$\langle X_{\mathcal{B}}^i(t) X_{\mathcal{B}}^j(s)\rangle_{\mathcal{B}} \coloneqq \mathrm{Tr}_{\mathcal{B}}\left[\rho_{\mathcal{B}} X_{\mathcal{B}}^i(t) X_{\mathcal{B}}^j(s)\right], \quad (8)$$

where the time dependence is due to the bath Hamiltonian. Suppose they are spatially coherent across the composite system, namely, $\langle X_{\mathcal{B}}^S(t) X_{\mathcal{B}}^{SA}(s)\rangle_{\mathcal{B}} = f(t-s)$ for some scalar function f. Perfect spatial correlations at the location of system and ancilla would imply that the operators $X_{\mathcal{B}}^j$ are identical except for a proportionality factor. Let $X_{\mathcal{B}}^{AS} = c_{SA} X_{\mathcal{B}}^S$, with $c_{SA} \in \mathbb{R}$. Then, $(X_{\mathcal{B}}^S + \alpha\ell X_{\mathcal{B}}^{SA}) = \gamma X_{\mathcal{B}}^S$ and $\gamma = 1 + c_{SA}\alpha\ell$ vanishes when $\alpha = -1/(\ell c_{SA})$. For this value, the system S is protected against the detrimental effects of the environment. When instead there are only partial correlations, the effect of the noise can be reduced but not entirely canceled. In this formalism, this case corresponds to a fundamental limitation to the possibility of tuning the parameter α.

We note that cancellation is achieved independently of the time correlations of the bath, as long as the noise is spatially correlated between system and ancilla. This demonstrates the ability of our proposal to cancel decoherence also for colored noise manifesting memory effects, including for $1/f$ noise, which present error-correction protocols cannot address (see Sec. VI of SM [32] for the details and a formal proof).

As in Secs. III and IV of SM [32] we show how it can be implemented in a Bose-Hubbard setting, we now return to the stochastic Hamiltonian (5) and concretely derive noise cancellation in this context.

Assume the correlations of the white noises $w_S(t)$, $w_A(t)$, and $w_{SA}(t)$ in the Hamiltonian (5) satisfy $\langle\langle w_i(t) w_j(s)\rangle\rangle = \Lambda_{ij}\delta(t-s)$ with correlation matrix

$$\Lambda = [\Lambda_{ij}] = \lambda \begin{pmatrix} 1 & 1 & \alpha \\ 1 & 1 & \alpha \\ \alpha & \alpha & \alpha^2 \end{pmatrix}, \quad (9)$$

which is positive semidefinite for $\lambda > 0$ and $\alpha \in \mathbb{R}$; see Refs. [42,43] and Sec. I of SM [32]. Averaging over the noises, these conditions yield a Lindblad master equation $\partial_t \rho_{SA}(t) = -i[H_0, \rho_{SA}] + \mathbb{L}\rho_{SA}(t)$ for the density matrix of S and A. Its GKSL generator is of the form $\mathbb{L} = \sum_{i,j=1}^3 \mathbb{L}_{ij}$, where

$$\mathbb{L}_{ij}\rho_{SA} = \Lambda_{ij}\left(A_i \rho_{SA} A_j - \frac{1}{2}\{A_j A_i, \rho_{SA}\}\right), \quad (10)$$

with jump operators $A_1 = J_S^z \otimes \mathbb{I}_A$, $A_2 = \mathbb{I}_S \otimes J_A^z$, and $A_3 = J_S^z \otimes J_A^z$.

Notice that the correlation matrix Λ has rank one and therefore the GKSL generator can be written with a single effective Lindblad operator $\hat{O}_{\mathrm{eff}} = \sqrt{\lambda}(\mathbb{I}_S \otimes J_A^z + J_S^z \otimes J_A(\alpha))$, where

$$J_A(\alpha) = \mathbb{I}_A + \alpha J_A^z. \quad (11)$$

So, it proves convenient to recast the S-A Lindblad master equation as

$$\partial_t \rho_{SA}(t) = -i[H_0, \rho_{SA}] + \mathbb{M}[\hat{O}_{\mathrm{eff}}]\rho_{SA}. \quad (12)$$

This form allows us to recover Eq. (2) after setting $\alpha = 0$ and tracing over the ancilla Hilbert space.

Now, assuming that the auxiliary system A is prepared in an eigenstate $|\ell\rangle_A$ of J_A^z with $J_A^z|\ell\rangle_A = \ell|\ell\rangle_A$, and considering as initial states $\rho_{SA} = \rho_S \otimes P_A^\ell$, where $P_A^\ell = |\ell\rangle\langle\ell|_A$, then the action of the effective Lindblad operator reduces to

$$\hat{O}_{\mathrm{eff}}(\rho_S \otimes P_A^\ell) = \sqrt{\lambda}[\ell\mathbb{I}_S + (1+\alpha\ell)J_S^z]\rho_S \otimes P_A^\ell, \quad (13)$$

so that the dissipator factorizes as

$$\mathbb{M}[\hat{O}_{\mathrm{eff}}](\rho_S \otimes P_A^\ell) = \lambda\mathbb{M}[\hat{L}_S]\rho_S \otimes P_A^\ell, \quad (14)$$

where $\hat{L}_S \coloneqq [\ell\mathbb{I}_S + (1+\alpha\ell)J_S^z]$. One can choose $\alpha = -1/\ell$, so that $J_A(\alpha)P_A^\ell = P_A^\ell J_A(\alpha) = 0$ and $\hat{L}_S \propto \mathbb{I}_S$. This condition leads to

$$-i[H_0, \rho_{SA}] = -i[H_S - J_S^z, \rho_S] \otimes P_A^\ell, \quad (15)$$

$$\mathbb{M}[\hat{O}_{\mathrm{eff}}]\rho_{SA} = 0. \quad (16)$$

As a result, the superoperator on the right-hand side of (12) acts simply as

$$\mathbb{L}(\rho_S \otimes P_A^\ell) = -i[H_0 - J_S^z, \rho_S] \otimes P_A^\ell. \qquad (17)$$

Consequently, at this optimal coupling, the dynamics $\gamma_t = e^{t\mathbb{L}}$ generated by the exponential action of $\mathbb{L}$ is unitary and leaves the ancilla unaffected,

$$\gamma_t(\rho_S \otimes P_A^\ell) = e^{-itH_S^{eff}} \rho_S e^{itH_S^{eff}} \otimes P_A^\ell, \qquad (18)$$

where $H_S^{\rm eff} = H_S - J_S^z$ is the effective Hamiltonian with a Lamb-shift correction. Notice that, under the above conditions, the Hamiltonian dynamics in (18) holds true *for all* system initial states ρ_S.

Interactions between the main and auxiliary system, such as the ones discussed here, could be implemented in platforms, including main and spectator qubits, such as, for instance, in [11], as well as atomic tweezers for atoms inside a resonator [44], for Rydberg arrays [45–49], including hybrid lattice-tweezer platforms [50]. In the tweezer setups, the correlated noise could be for instance a state-dependent fluctuation of the tweezer trap potential that the atoms experience. Crucial for the viability is the capability of preparing the auxiliary system in a Fock state and to control the system-ancilla interaction. While there is an optimal working point, in what follows we show that the protocol's efficiency is robust against fluctuations of the optimal parameters. For this purpose, we discuss the protocol efficiency for the stabilization of NOON states against dephasing noise. We consider two double-well traps for bosons with long-range interactions [51–54]: one double well representing the system S and the other the ancilla A, where in the double well corresponding to the system S, a NOON state has been prepared, $|\Psi_S^\pm\rangle = |N,0\rangle_S \pm |0,N\rangle_S/\sqrt{2}$, where $|N,0\rangle_S$ is the state of N bosons in the left well and 0 in the right, while $|0,N\rangle$ is the state with the swapped occupations. It satisfies $J_S^z|\Psi_S^\pm\rangle = -(N/2)|\Psi_S^\mp\rangle$ and $(J_S^z)^2|\Psi_S^\pm\rangle = (N^2/4)|\Psi_S^\pm\rangle$. In the presence of dephasing noise, the stochastic Hamiltonian of system and ancilla can be reduced to the form $H'_{SA} = H_{SA} + \frac{1}{2}\alpha w_{SA}(t) K_{SA}$ with α the coupling strength between the two double wells and $K_{SA} = (J_S^z)^2 \otimes \mathbb{I}_A + \mathbb{I}_S \otimes (J_A^z)^2$; see Sec. III A of SM [32]. For $\alpha = 0$, the stochastic Hamiltonian H'_{SA} predicts that quantum coherence is lost over a timescale $1/\lambda$. If instead the ancilla is prepared in the eigenstate $|\ell\rangle_A$ and $\alpha_c = -1/\ell$, the effects of noise cancel exactly for $\alpha = -1/\ell$, as visible in Fig. 2.

The protocol is robust against small variations of α from the optimal value α_c. Indeed, given the fidelity $\mathcal{F}(\alpha,N)(t) = \mathrm{Tr}[\rho_{SA}(t)\rho_{\rm NOON} \otimes P_A^\ell]$, its time average $\overline{\mathcal{F}(\alpha,N)}$ signals how frequently, for a given α, the time-evolving NOON state differs from itself. Figure 3 displays the behavior of $\overline{\mathcal{F}(\alpha,N)} = \lim_{T\to\infty}(1/T)\int_0^T dt \mathcal{F}(\alpha,N)(t)$ as a function of α and for different values of N. For any finite N, the residual discrepancy $1 - \overline{\mathcal{F}(\alpha_c,N)}$ is due to the non-vanishing time-averaged contribution of the unitary oscillations generated by the effective Hamiltonian $H_S^{\rm eff}$ in Eq. (18). These residual oscillations scale as $1/N$ and thus vanish in the $N \to \infty$ limit. The central width of the curves decreases with increasing strength λ of the dephasing, as shown in Fig. 3. Yet, in the weak-coupling limit we require $\lambda \ll 1$; the same figure shows that for $0.001 \le \lambda \le 0.01$ the narrowing of the resonance is not dramatic.

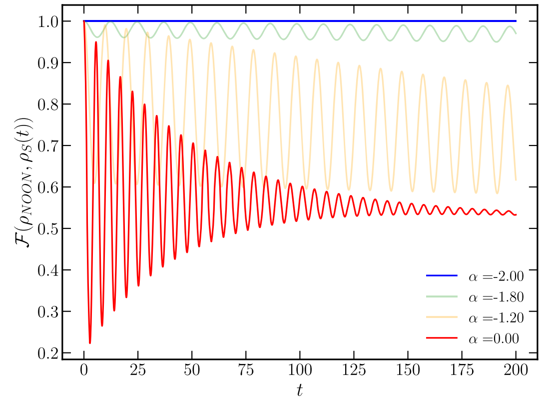

FIG. 2. Fidelity of staying NOON at time t for a system S with one qubit, $N_S = N_A = 1$, with fixed $\ell = 1/2$, for different values of α. Optimal cancellation of decoherence is achieved for $\alpha = -1/\ell = -2$. Here, we used $\gamma_S = 0.5$, $\gamma_A = 0$, and $\Delta_S = \Delta_A = -1$, which also guarantees cancellation of the residual unitary oscillations.

We emphasize that perfect cancellation requires (i) spatially correlated noise between system and ancilla, (ii) control over the state of the ancilla, and (iii) the possibility of tuning the coupling constant α to an optimal value. This value depends on the initial state of the ancilla, which gives some room for optimization as one could tune jointly both α and ℓ so that the parameter $\alpha\ell = -1/c_{SA}$. For the specific ancilla operators of the example in Eq. (5), one could enhance the number of available dark channels by independently increasing the number of bosons N_A in the ancillary sector. Specifically, in the Jordan-Schwinger representation of the $\mathfrak{su}(2)$ algebra with N_A bosons, the operator J_A^z admits $N_A + 1$ distinct eigenstates. Each of these can serve as a candidate ancillary dark state, thereby providing a larger set of interference channels that can be exploited to suppress noise more effectively.

Because of these features, the protocol is in principle scalable. We provide an illustration of its efficiency for the case of the S NOON state. Figure 3 shows that increasing the dimension of S and A improves the protocol efficiency against noise. We note, moreover, that the condition on state preparation of the ancilla can be relaxed. In fact, in Sec. V of SM [32] it is shown that, if the ancilla is prepared in a mixed state with mean excitation $\overline{\ell}$ and width $\Delta\ell$, white-noise decoherence is canceled when $\alpha = -1/\overline{\ell}$ and $\lambda(\Delta\ell)^2 \ll 1$.

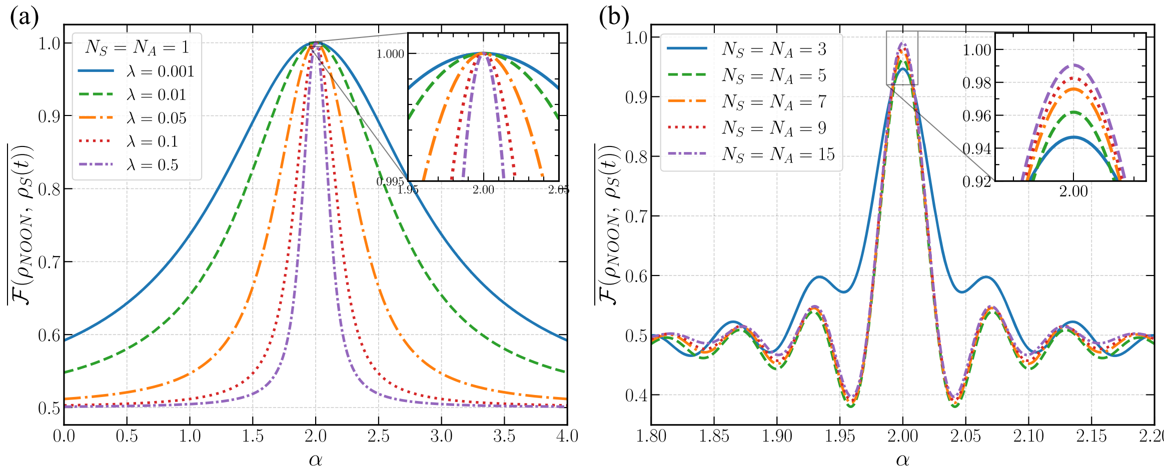

FIG. 3. Time-averaged fidelity $\bar{\mathcal{F}}$ for preserving the NOON state as a function of the interaction parameter α. Each curve corresponds to a distinct configuration with $N_S = N_A$, where the ancilla is prepared in the projector P_A^ℓ onto the eigenstate of J_A^z with $\ell = -\frac{1}{2}$. (a) For $N_S = N_A = 1$, the system reduces to a single qubit S, exhibiting a resonance peak centered at $\alpha_c = 2$. The resonance width decreases as the dephasing strength λ increases. (b) For odd $N_S = N_A > 1$ and $\lambda = 0.1$, $\bar{\mathcal{F}}$ displays a pronounced peak near $\alpha_c = 2$. The inset provides a magnified view, showing that the peak averaged fidelity approaches unity as the number of bosons increases. Data for even $N_S = N_A$ follow the same trend and are omitted for clarity.

In conclusion, our proposal exploits dynamically generated interferences to protect any state of a dephasing system S from noisy effects. It differs from existing techniques such as those based on the identification of decoherence-free subspaces or on feedback; rather, it resembles the mechanisms behind classical stochastic resonance by virtue of a resonant response that leads to noise cancellation [55]. Because of its structure, which is independent from the noise due to classical fluctuations, the method appears extendable to true decoherence as generated by quantum environments. Furthermore, it is expected to work for any spectral property of the dephasing noise, as long as it is spatially correlated. For its properties it can provide a key strategy for high-precision metrology or sensing, and can increase the quantum volume of a circuit when the noise is non-Markovian and typical quantum error-correction codes become inefficient.

Acknowledgments—F. B. acknowledges financial support from PNRR MUR Project No. PE0000023-NQSTI. G. M. acknowledges funding from the German Ministry of Education and Research (BMBF, Project “NiQ: Noise in Quantum Algorithms“), by the Deutsche Forschungsgemeinschaft (DFG, German Research Foundation)—Project No. 429529648—TRR 306 QuCoLiMa (“Quantum Cooperativity of Light and Matter”), and by the QuantERA II Programme (project “QNet: Quantum transport, metastability, and neuromorphic applications in Quantum Networks”), which has received funding from the EU’s Horizon 2020 research and innovation programme under Grant Agreement No. 101017733, as well as from the Deutsche Forschungsgemeinschaft DFG (Project No. 532771420). The authors thank Andrea Trombettoni and Johannes Zeiher for insightful discussions regarding the critical aspects and possible experimental implementations.

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.

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