Controlling Markovianity with Chiral Giant Atoms

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Giant artificial atoms are promising and flexible building blocks for the implementation of analog quantum simulators. They are realized via a multilocal pattern of couplings of two-level systems to a waveguide, or to a two-dimensional photonic bath. A hallmark of giant-atom physics is their *non-Markovian* character in the form of self-coherent feedback, leading, e.g., to nonexponential atomic decay. The timescale of their non-Markovianity is essentially given by the time delay proportional to the distance between the various coupling points. In parallel, with the state-of-the-art experimental setups, it is possible to engineer complex phases in the atom-light couplings. Such phases simulate an artificial magnetic field, yielding a *chiral* behavior of the atom-light system. Here, we report a surprising connection between these two seemingly unrelated features of giant atoms, showing that the chirality of a giant atom controls its Markovianity. In particular, by adjusting the couplings' phases, a giant atom can, counterintuitively, enter an exact Markovian regime, irrespectively of any inherent time delay. We illustrate this mechanism as an interference process and via a collision model picture. Our findings significantly advance the understanding of giant atom physics, and open new avenues for the control of quantum nanophotonic networks.

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The growing demand to process quantum information for computational purposes underscores the increasing importance of developing scalable quantum networks. These networks consist of spatially distributed nodes interconnected by communication lines. Consequently, investigating the realm where memory effects and quantum feedback are not negligible becomes increasingly crucial in addressing the challenge of quantum computation [1–7]. A notable instance includes multilocal or *giant* atoms [8–12], which are two-level emitters coupled to an environment (such as a field flowing through a waveguide) at multiple spatially separated points. As the light travels among these distinct coupling points, it accumulates a phase (optical length) φ_{WG} proportional to the distance between them. When the coupling points are spaced at distances comparable to the wavelength of the light they interact with, a direct consequence of the phase accumulation is that self-interference effects, absent with ordinary atoms, arise.

A remarkable feature of giant atoms is their *non-Markovian* character. Indeed, a giant atom can reabsorb its own emitted excitation after a time delay proportional to the distance between the coupling points. This phenomenon has been experimentally demonstrated with superconducting qubits coupled to surface acoustic waves [13], and spurred the interest in giant atoms physics. Interestingly, as we detail below, a giant atom coupled to a waveguide at two different coupling points can be described in terms of a small atom (one coupling point) in a semi-infinite waveguide, Figs. 1(a)–1(b1), a typical setup to observe a non-Markovian behavior of the atomic emission [14–16].

Another striking feature of giant atoms is the possibility of engineering dispersive decoherence-free interactions between them [17]. Remarkably, even if the latter effect is inherently related to the phase differences associated with the displacements of the coupling points, it becomes prominent when the travel time of light between coupling points is small compared to the characteristic timescales of the emitters, i.e., in the *Markovian* limit [17,18].

In parallel to memory effects, another important aspect concerns the potential to adjust the propagation direction of light between the nodes. When scattered radiation displays a preferred direction, the interaction between emitters and light is defined as *chiral* [19]. Both theoretically [2] and very recently experimentally [20] it has been shown that introducing light-matter couplings with an additional complex phase can induce a chiral behavior in the radiation emitted by giant atoms. For a giant atom with two coupling points, cf. Fig. 1(a), the atomic emission is chiral whenever the phase difference between the couplings φ_c does not vanish. In particular, when such a phase matches the optical length φ_{WG}, the emission can become maximally chiral [2].

In this work, we bridge these two seemingly unrelated features of giant atoms, namely, their *chirality* and their *(non-) Markovianity*. We show how to make a giant atom enter the Markovian regime, even for non-negligible time delays, by tuning its chirality. Importantly, our result depends *solely* on the complex nature of the atom-light

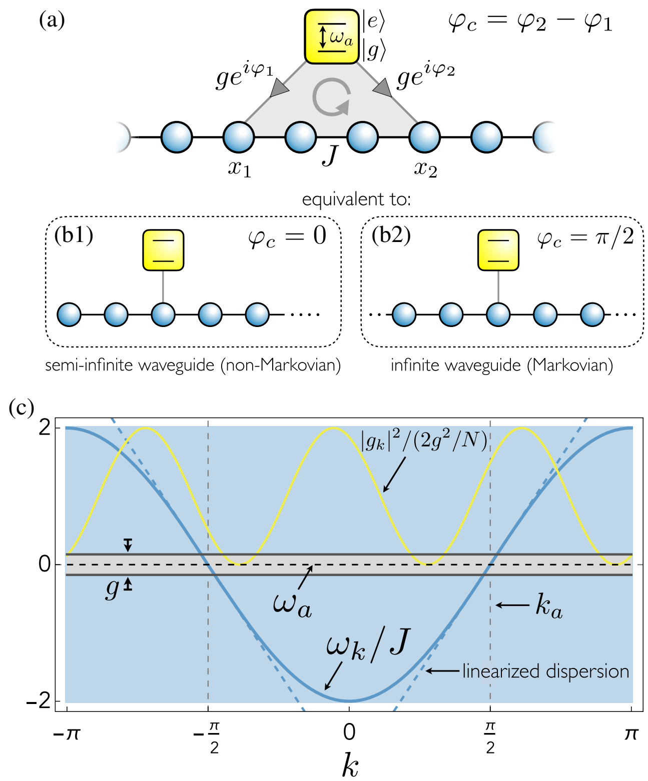

FIG. 1. Setup. (a) Chiral giant atom coupled at the points $x_{1,2}$ to a 1D bidirectional waveguide. The couplings are generally complex with phases $\varphi_{1,2}$, cf. Eq. (2). When $\varphi_c = \pi/2$ [$\varphi_c = 0$], as in (b2) [(b1)], the system is equivalent to a small atom (local coupling) in a [semi-]infinite waveguide and the dynamics is [non-]Markovian. (c) Dispersion law of the waveguide (solid blue), which in weak coupling (gray stripe) can be linearized (dashed blue). The atomic frequency ω_a is resonant with the middle of the band, corresponding to the momentum $k_a = \pi/2$. In yellow we show an instance of a chiral (indeed, $|g_{k_a}| > |g_{-k_a}|$) atom-waveguide coupling with $x_2 - x_1 = 3$ and $\varphi_c = \pi/6$.

couplings. This result definitely shows that the identifying feature of giant atoms is the nonlocality of their couplings, rather than their non-Markovianity.

Setup and Hamiltonian—We consider a single giant atom weakly coupled to a one-dimensional (1D) bidirectional waveguide. The full light-matter Hamiltonian is $\hat{H} = \hat{H}_a + \hat{H}_w + \hat{H}_{\text{int}}$. The free atomic Hamiltonian is $\hat{H}_a = \omega_a \hat{\sigma}^\dagger \hat{\sigma}$, with $\hat{\sigma} = |g\rangle\langle e|$, $|g\rangle$ and $|e\rangle$ being the ground and excited atomic states, respectively. We model the waveguide with a translationally invariant tight-binding array of coupled resonators with Hamiltonian

$$\hat{H}_w = -J \sum_x \hat{a}^\dagger_{x+1} \hat{a}_x + \text{H.c.}, \tag{1}$$

where $\hat{a}_x$ are real space bosonic annihilation operators and $J > 0$. We can Fourier transform $\hat{a}_x = \sum_k e^{-ikx} \hat{a}_k / \sqrt{N}$, where N is the number of resonators, so that Eq. (1) becomes $\hat{H}_w = \sum_k \omega_k \hat{a}^\dagger_k \hat{a}_k$, with $\omega_k = -2J \cos k$ (the lattice constant is set to 1). The atom-waveguide interaction Hamiltonian is

$$\hat{H}_{\text{int}} = g\hat{\sigma}(e^{i\varphi_1} \hat{a}^\dagger_{x_1} + e^{i\varphi_2} \hat{a}^\dagger_{x_2}) + \text{H.c.}, \tag{2}$$

where we assume g to be real and $x_2 = x_1 + d$. We focus here on a two-legged giant atom, though our result generalizes to the case of multiple coupling points, as we detail in [21]. In Fourier space the interaction Hamiltonian (2) reads $\hat{H}_{\text{int}} = \sum_k \hat{h}_{\text{int}}(k)$, where $\hat{h}_{\text{int}}(k) = g_k \hat{\sigma} \hat{a}^\dagger_k + \text{H.c.}$ and $g_k = g[e^{i(\varphi_1 + kx_1)} + e^{i(\varphi_2 + kx_2)}]/\sqrt{N}$. When all these phases are zero, we refer to the giant atom as *nonchiral*. By contrast, we will call the giant atom *chiral* whenever we take into account nonzero phases. This is because, in the latter case, time-reversal symmetry is broken [$g_k \neq g_{-k}$, or $T\hat{h}_{\text{int}}(k)T^{-1} \neq \hat{h}_{\text{int}}(-k)$, T being the time-reversal symmetry operator]. The assumption of weak coupling makes our system equivalent to a giant atom coupled to a continuous waveguide with linear dispersion [22], see Fig. 1(c). Therefore, from an experimental point of view, our waveguide Hamiltonian can be implemented with a continuous transmission line [20], as well as with an array of coupled superconducting circuits [23–25].

Result—Our result can be condensed in the following sentence: the chirality of a giant atom controls its Markovianity. Remarkably, Markovianity can be achieved irrespectively of any time delay. This implies that such a chiral giant atom undergoes spontaneous emission even when the coupling points are significantly far apart, when reabsorption would occur in the nonchiral case. Despite the atom being *giant*, in the sense that a non-Markovian behavior is expected, it behaves as if it were *small* (i.e., single-legged). Thus, we argue in favor of the *nonlocality of the couplings* as a defining feature of giant atoms.

We derive this result through the analytic calculation of the atomic dynamics and further check it through the Lindblad master equation. We then provide two mechanisms for this phenomenon, based on an interference argument and on a collision model picture [10,26,27].

Assume the initial state is $|\Psi(0)\rangle = |e\rangle|0\rangle$, $|0\rangle$ being the vacuum state of the field. Then at time t the full atom-waveguide state is $|\Psi(t)\rangle = \varepsilon(t)|e\rangle|0\rangle + \sum_k c_k(t) \hat{a}^\dagger_k |g\rangle|0\rangle$. Imposing the Schrödinger equation, the dynamics of an initially excited chiral giant atom follows the delay differential equation [21]

$$\dot{\varepsilon}(t) = -\frac{\Gamma}{2}\varepsilon(t) - \frac{\Gamma}{2}\cos(\varphi_c) e^{i\varphi_{\text{WG}}} \Theta(t - t_d) \varepsilon(t - t_d). \tag{3}$$

Here, $\Theta(t)$ is the Heaviside step function, φ_{WG} is the optical length between the coupling points, t_d is the corresponding time delay, and Γ is the decay rate. More specifically, the optical length is given by $\varphi_{\text{WG}} = k_a d$, where d is the distance between the coupling points, k_a is the momentum corresponding to the atomic transition frequency ω_a ($k_a = \pi/2$ in our case), $v = 2J\sin(k_a)$ is the speed of light in the waveguide, and $\Gamma = 4g^2/v$.

For $\varphi_c = 0$, Eq. (3) is well known [14], and indeed shows the analogy between a nonchiral giant atom and a small atom in front of a mirror, cf. Figs. 1(a)–1(b1). The only difference with Ref. [14] is the minus sign in front of the second term in the right-hand side of Eq. (3). Such π phase difference is due to the fact that for an atom in front of a mirror the optical length is proportional to twice the distance between the atom and the mirror.

Interestingly, we observe that the atomic decay is exactly exponential at $\varphi_c = \pi/2$, regardless of the time delay t_d, matching the behavior of a small atom coupled to a waveguide. By contrast, the non-Markovianity of a chiral giant atom is prominent when $\varphi_{\mathrm{WG}} = m\pi$ and $\varphi_c = (m+1)\pi$ with integer $m \geq 0$. At these values, part of the emitted light is trapped between the coupling points forming a bound state in the continuum (BIC) [28,29]. In Fig. 2, we show the dynamics of an initially excited chiral giant atom, which indeed decays exponentially, regardless of the distance between the coupling points, for $\varphi_c = \pi/2$.

This behavior can be captured as well through the atomic master equation [30] for small distances between the coupling points. In this case the Markov approximation, which makes the evolution of the density matrix time local, is still valid. Notwithstanding, the interaction Hamiltonian still keeps track of the spatial nonlocality of the atom-light interaction. The atomic master equation reads $\dot{\rho} = -i[\hat{H}_a, \rho] + \gamma(\hat{\sigma}\rho\hat{\sigma}^\dagger - \{\hat{\sigma}^\dagger\hat{\sigma}, \rho\}/2)$, where [21]

$$\gamma = \Gamma[1 + \cos(\varphi_c)\cos(k_a d)]. \tag{4}$$

This matches the analytical result, predicting an exponential decay with rate Γ for $\varphi_c = \pi/2$. The same rate is obtained for $k_a d = \pi/2$, which, for our discretized waveguide, corresponds to odd distances d. Indeed, Fig. 2 shows that for odd d's the atomic deexcitation slightly deviates from the exponential decay one would get from the Lindblad master equation for any complex phase φ_c. We notice though that this discrepancy increases with d as the Markov approximation breaks down.

Mechanism—First, the exponential atomic decay at $\varphi_c = \pi/2$ can be explained as an interference effect, Fig. 3(a). The key observation is that the phases $\varphi_{1,2}$ and φ_{WG} have to be considered with and without their signs, respectively. Indeed, φ_{WG} is always positive regardless of the interference path, while φ_c is positive (negative) when going from the atom (field) to the field (atom).

At the left and right coupling points, (x_1, x_2), we can write the emitted field amplitude at time t as $c_{1,2}(t)$ [21]. We can further divide these terms into backward (to the left) and forward (right) emitted field amplitude

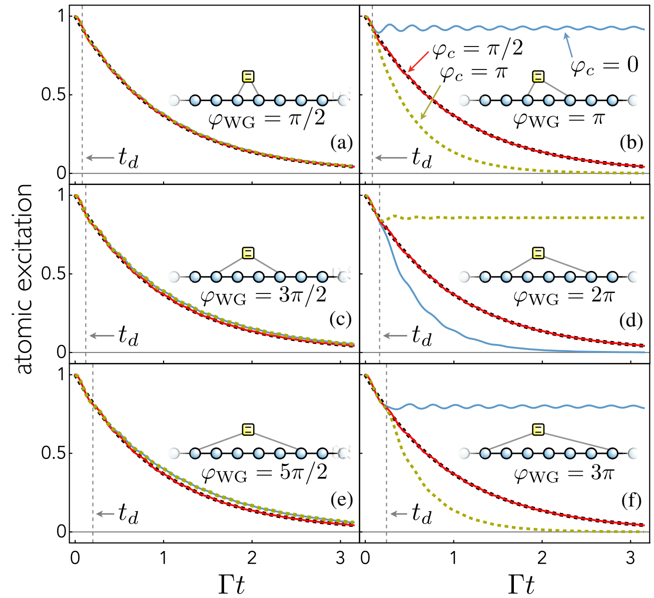

FIG. 2. Markovianity for any time delay. Dynamics of the atomic excitation, $|\varepsilon(t)|^2$, of an initially excited chiral giant atom for various distances d between the coupling points, corresponding to various optical lengths $\varphi_{\mathrm{WG}} = k_a d$ and phase differences φ_c between the couplings to the waveguide. Specifically, $d = 1, 2, 3, 4, 5, 6$ in (a)–(f), respectively. In all panels (a)–(f) $\varphi_c = 0, \pi/2, \pi$ correspond to blue, red, and dashed yellow, respectively. The black dotted lines represent the exponential decay $e^{-\Gamma t}$, while the other curves are obtained numerically. Regardless of the time delay $t_d = d/v$, vertical dashed gray line, increasing from (a) to (f), at $\varphi_c = \pi/2$ the decay is exactly exponential. When the optical length is an integer multiple of π, the atom never fully decays and correspondingly a BIC occur. On top of the BICs occurring at odd multiples of π for a nonchiral giant atom, (b) and (f), as in Ref. [14], complex couplings allow the appearance of BICs at even multiples of π as well [21]. Other parameter values: $k_a = \pi/2$, $g = 0.2J$, $v = 2J$, $N = 90$ (number of resonators), the first (second) coupling point is at $N/2$ $(N/2 + d)$.

$c_1(t) = c_{1,\mathrm{b}}(t) + c_{1,\mathrm{f}}(t)$, analogously for $c_2(t)$. The backward emitted field at coupling point x_1 has (i) a contribution coming directly from the atomic exponential decay, and (ii) a contribution coming from the backward emitted field at coupling point x_2. Thus, we can further divide $c_{1,\mathrm{b}}(t)$ into its *exponential* (exp) and *delay* (del) contributions as $c_{1,\mathrm{b}}(t) = c^{\mathrm{del}}_{1,\mathrm{b}}(t) + c^{\mathrm{exp}}_{1,\mathrm{b}}(t)$ and $c_{2,\mathrm{f}}(t) = c^{\mathrm{exp}}_{2,\mathrm{f}}(t) + c^{\mathrm{del}}_{2,\mathrm{f}}(t)$ (the same goes for the forward emitted field at coupling point x_2). Note that $c^{\mathrm{del}}_{1,\mathrm{b}}(t) = c^{\mathrm{del}}_{2,\mathrm{f}}(t) = 0$ for $t < t_d$.

The only way to increase the atomic amplitude ε is via the field contribution coming from the delay, which is (we drop the time dependence to lighten notation)

$$c^{\mathrm{del}}_{1,\mathrm{b}} e^{-i\varphi_1} + c^{\mathrm{del}}_{2,\mathrm{f}} e^{-i\varphi_2}. \tag{5}$$

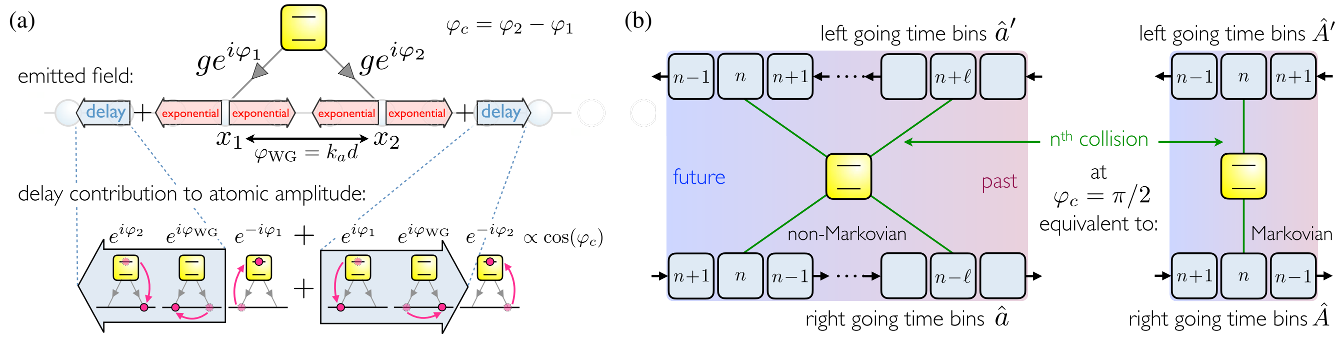

FIG. 3. Mechanism of Markovianity for any time delay. (a) Interference picture. At each coupling point (x_1, x_2) the emitted field can be divided into forward and backward components (right- and left-pointing blue and red arrows, respectively) $c_{1,\mathrm{f}}(t)$ and $c_{1,\mathrm{b}}(t)$ (same for x_2). For $t \geq t_d = d/v$, $d = x_2 - x_1$, the forward [backward] component at the coupling point x_2 [x_1] acquires a delay contribution $c^{\mathrm{del}}_{2,\mathrm{f}}(t)$ [$c^{\mathrm{del}}_{1,\mathrm{b}}(t)$] (light blue arrows), as the atomic excitation (small fuchsia circle) is transferred by hopping to the x_1 [x_2] coupling point and traveling along the waveguide for a distance d. A non-Markovian behavior takes place only if the atom gets reexcited. The amplitude for this process to occur is the sum of the two possible ways for the reexcitation to happen, as illustrated at the bottom part. Considering the phases acquired along these two paths, the final amplitude is proportional to $\cos(\varphi_c)$. (b) Collision model picture. Left- and right-going field modes (top and bottom, respectively) are mapped into trains of time-bin ancillae moving in opposite directions. At each time step the system interacts with two separated ancillae from each bath and the chains are shifted by one position. After the first collision, each ancilla will interact again with the system in ℓ steps. When the complex phase $\varphi_c = \pi/2$ (and odd multiples) this picture is unitarily equivalent to a Markovian collision model where the system interacts locally with the two left- and right-going time bins of a bidirectional waveguide.

On the other hand, the atomic population contribution to these field components is

$$c^{\mathrm{del}}_{1,\mathrm{b}} = e^{i\varphi_2} e^{i\varphi_{\mathrm{WG}}} \varepsilon, \qquad c^{\mathrm{del}}_{2,\mathrm{f}} = e^{i\varphi_1} e^{i\varphi_{\mathrm{WG}}} \varepsilon. \tag{6}$$

By plugging Eqs. (6) into Eq. (5), the delay contribution to the atomic amplitude turns out to be proportional to $\cos(\varphi_c)$. Therefore, for $\varphi_c = \pi/2$ there is no delay contribution, and the atomic excitation decays exactly exponentially irrespectively of any distance between coupling points.

Second, in the framework of collision models [26], the time evolution of the atom and the waveguide is described as a sequence of discrete interactions (collisions) involving the system (the atom) and discretized field modes (the ancillae or time bin modes). In the case of a giant atom, the interaction involves two separated ancillae, see Fig. 3(b). A non-Markovian behavior typically arises due to the double interaction between the system and the same ancilla after a finite time. For a two-legged giant atom, the coupling Hamiltonian (2) in interaction picture with respect to the waveguide and the atom reads

$$\hat{H}_{\mathrm{int}}(t) = g\hat{\sigma}^\dagger \big[e^{-ik_a x_1} \hat{a}_{t-\tau_1} + e^{ik_a x_1} \hat{a}'_{t+\tau_1} + e^{i\varphi_c}\big(e^{-ik_a x_2} \hat{a}_{t-\tau_2} + e^{ik_a x_2} \hat{a}'_{t+\tau_2}\big)\big] + \mathrm{H.c.} \tag{7}$$

Without loss of generality we have set $\varphi_1 = 0$ and thus $\varphi_2 = \varphi_c$. In Eq. (7), $\hat{a}_{t-\tau_{1,2}}$ $(\hat{a}'_{t+\tau_{1,2}})$ are the right-(left-) going time bin operators corresponding to coupling points $x_{1,2}$ [18], and $\tau_{1,2} = x_{1,2}/v$ are the time-domain coordinates corresponding to the coupling points' positions. The left-going operators have a prime to stress the distinction with the right-going ones. For an infinitesimal evolution time Δt, the related propagator is approximated as $\hat{\mathcal{U}}_n \simeq \mathbb{1} - i(\hat{\mathcal{H}}^{(0)}_n + \hat{\mathcal{H}}^{(1)}_n)\Delta t - (\hat{\mathcal{H}}^{(0)}_n)^2 \Delta t^2/2$, where $\hat{\mathcal{H}}^{(0)}_n = (1/\Delta t)\int_{t_{n-1}}^{t_n} \mathrm{d}s \hat{H}_{\mathrm{int}}(s)$ and $\hat{\mathcal{H}}^{(1)}_n = (i/2\Delta t)\int_{t_{n-1}}^{t_n} \mathrm{d}s \int_{t_{n-1}}^{s} \mathrm{d}s' [\hat{H}_{\mathrm{int}}(s'), \hat{H}_{\mathrm{int}}(s)]$ are the 0th and the 1st order terms of the Magnus expansion of the generator [31]. The operators of both the emitter and the waveguide modes are present only in the 0th-order term, i.e., only $\hat{\mathcal{H}}^{(0)}_n$ describes the interaction while $\hat{\mathcal{H}}^{(1)}_n$ is a Lamb-shift term. We thus write the former as

$$\hat{\mathcal{H}}^{(0)}_n = g\hat{\sigma}^\dagger \big[\hat{a}_n + e^{i\varphi_c} e^{-i\varphi_{\mathrm{WG}}} \hat{a}_{n-\ell} + \hat{a}'_n + e^{i\varphi_c} e^{i\varphi_{\mathrm{WG}}} \hat{a}'_{n+\ell}\big] + \mathrm{H.c.}, \tag{8}$$

where we introduced the right-going time-bin operators $\hat{a}_n = (1/\sqrt{\Delta t})\int_{t_{n-1}}^{t_n} \mathrm{d}t \hat{a}_t$ (the same holding for the left-going ones $\hat{a}'_n$). Without loss of generality, we set $x_1 = 0$, $x_2 = d$ and $\ell \Delta t = \tau_2$ (which is nothing but the time delay t_d). The relation (8) captures all the relevant physics of the atom-waveguide crosstalk in all regimes and makes clear the (general) non-Markovian nature of the dynamics.

Consider now a separate (Markovian) collision model of an emitter coupled to a bidirectional waveguide, whose nth collision is described by the Hamiltonian

$$\hat{h}^{(0)}_n = \sqrt{2} g\hat{\sigma}^\dagger (\hat{A}_n + \hat{A}'_n) + \mathrm{H.c.}, \tag{9}$$

where $\hat{A}_n$ and $\hat{A}'_n$ are the right- and left-going time-bin operators, respectively. Setting $\hat{A}_n = (\hat{a}_n + e^{i(\varphi_c - \varphi_{\mathrm{WG}})}\hat{a}_{n-\ell})/\sqrt{2}$ and $\hat{A}'_n = (\hat{a}'_n + e^{i(\varphi_c + \varphi_{\mathrm{WG}})}\hat{a}'_{n+\ell})/\sqrt{2}$, these transformations define a *unitary* transformation of the field time bins *if and only if* $\varphi_c = \pi/2$ [21,32].

Therefore, for $\varphi_c = \pi/2$, there exists an exact mapping between Eq. (8) (generally non-Markovian) and Eq. (9) (Markovian). Under this condition, the transformation does not alter the reduced dynamics of the system. Thus, the collision model describing a chiral giant atom with an arbitrary delay line between its legs becomes equivalent to a collision model where the atom interacts with the field at a single point with a rescaled coupling strength. Significantly, this implies that the reduced dynamics of the system is *exactly* Markovian, even in the presence of any time delay.

Conclusion—An artificial atom coupled at multiple points to a waveguide is a paradigmatic setup to observe memory effects due to self-interference. We have found that this paradigm can break down when the atom-light couplings are allowed to be complex. By properly adjusting the coupling phases, the artificial atom has an exact Markovian behavior, regardless of any inherent time delay involved in the dynamics. This unexpected effect enriches the already exotic physics of giant atoms, opening new theoretical and experimental avenues. Despite that most efforts are devoted to the engineering of decoherence-free Hamiltonians [11,33], our results show how unexpected phenomena can occur in the opposite regime, that is far from protecting the atom from decoherence. Also, the effect we find is relevant from the theoretical standpoint on its own. Indeed, many works righteously point out that time delays need to be neglected to derive a master equation for giant atoms in a waveguide [8,10,17,18]. Our result shows that, at least for a single giant atom, there is no need to make such approximation. We note that the effect we described could be tested, in principle, by coupling a transmon qubit either to a microwave photonic waveguide [8,20], or to an array of superconducting LC circuits [24,25]. Finally, the connection we find between chirality and Markovianity could be generalized to more structured photonic environments [16,34,35], which is a promising direction for future work.

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