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Article

Keywords:

Posted Date: March 11th, 2022

DOI: https://doi.org/10.21203/rs.3.rs-1407573/v1

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Version of Record: A version of this preprint was published at Communications Physics on August 22nd, 2022. See the published version at https://doi.org/10.1038/s42005-022-00991-3.
Nonreciprocal and chiral single-photon scattering for giant atoms

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(Dated: March 1, 2022)

In this work, we investigate the nontrivial single-photon scattering properties of giant atoms coupled to waveguides that can be an effective platform for realising nonreciprocal and chiral quantum optics. For the two-level giant-atom setup, we identify the condition for nonreciprocal transmission: the external atomic dissipation is further required other than the breaking of time-reversal symmetry by local coupling phases. Especially, in the non-Markovian regime, unconventional revival peaks periodically appear in the reflection spectrum of such a two-level giant-atom system. To explore more interesting scattering behaviours, we further extend the two-level giant-atom system to Δ-type and V-type three-level giant atoms coupled to double waveguides without external atomic dissipation. We analyse the different physical mechanisms for the nonreciprocal and chiral scattering properties of the Δ-type and V-type giant atoms. Our proposed giant-atom structures have potential applications of high-efficient single-photon targeted router and circulator for quantum information precessing.

INTRODUCTION

Waveguide quantum electrodynamics (QED) studies the interactions between atoms and one-dimensional waveguide modes, providing an excellent platform for constructing long-range interactions and engineering large-scale quantum networks [1–5]. In experiments, typical candidates of implementing waveguide QED systems include quantum dots coupled to photonic crystal waveguides [6, 7], superconducting qubits coupled to transmission lines [8, 9], ultracold atoms coupled to optical fibers [10, 11], etc. To date, waveguide QED has inspired a number of exotic phenomena, such as atom-like mirrors [12, 13], dynamic Casimir effects [14], single-photon routing [15–17], bound states in the continuum [18].

In general, the atom can be viewed as a point when coupled with the waveguide due to its negligible size compared to the wavelength of waveguide modes. Nevertheless, in a recent experiment, a superconducting transmon qubit was designed to interact with surface acoustic waves (SAWs) via multiple coupling points whose separation distances can be much larger than the wavelength of SAWs [19]. Instead, a generalized theory called giant atom has been developed to describe such situations [20]. Since the first theoretical study in 2014 [21], the giant-atom scheme has been broadly investigated with superconducting qubits [22–26], coupled waveguide arrays [27], and cold atoms [28]. With such nonlocal coupling schemes, a series of tempting quantum phenomena have been demonstrated, including frequency-dependent relaxation rate and Lamb shift [21, 25, 29], non-exponential atomic decay [22, 23], decoherence-free interatomic interaction [25, 30, 31], exotic bound states [24, 32], and modified topological effects [33]. Giant atoms have emerged as a new paradigm in quantum optics and require more comprehensive understanding in physics.

On the other hand, controlling the flow of photons, especially realising asymmetric photonic propagations in waveguide QED systems, is crucial for constructing non-reciprocal optical element devices [34–39]. To this end, one could break the time-reversal symmetry of the system such that the interactions between the atoms and the waveguide modes are direction-dependent [16, 40–44]. Such a paradigm, also known as chiral quantum optics [40], can be achieved via several methods, such as the spin-momentum locking effect [45–47], inserting circulators in superconducting circuits [48–50], applying topological waveguides [51, 52], and synthesizing artificial gauge fields [53]. Based on the chiral interaction, targeted photonic routers [17], single-photon circulators [54, 55], cascaded quantum networks [56–58], and enhanced entanglement [59, 60] have been realised. Recently, the concept of giant atom has been introduced to chiral quantum optics, making some advanced functionalities possible, such as chiral bound states [32], dark states without coherent drives [31], and non-Markovicity induced nonreciprocity [61]. These seminal works inspire us to explore more intriguing effects in chiral giant-atom setups, especially with multi-level structure [61–63].

In this paper, we investigate how external atomic dissipations outside the waveguide and local coupling phases affect the single-photon scattering properties of a two-level giant atom with two atom-waveguide coupling points. By taking account of the phase difference between two coupling points, we find that the giant atom behaves
like a chiral small atom in the Markovian regime but exhibits peculiar giant-atom effects in the non-Markovian regime. We physically demonstrate that the breaking of time-reversal symmetry by local coupling phases is not sufficient for realising nonreciprocal photon scatterings. In fact, in the absence of the external atomic dissipation, the scatterings are always reciprocal even if the atomic spontaneous emission becomes chiral \([63, 64]\). In order to realise asymmetric scattering for a giant atom without external dissipation, we propose a \(\nabla\)-type giant atom coupled to two waveguides. In such way, we realise the nonreciprocal and chiral scatterings with single \(\nabla\)-type atom. Targeted routing and circulation schemes can also be realised via such scatterings with proper phases. Finally, we consider a \(\Delta\)-type giant atom and compare its properties with that of \(\nabla\)-type one. We reveal that, the nonreciprocal scatterings stem from the quantum interference effect in the closed-loop atom-level structure for the \(\Delta\)-type giant atom, but from the nontrivial coupling phase difference for the \(\nabla\)-type giant atom.

RESULTS AND DISCUSSION

A. Two-level giant atom coupled to a single waveguide

As schematically shown in Fig. 1(a), we consider a two-level giant atom coupled to a waveguide at two separated points \(x = 0\) and \(x = d\). The atom-waveguide coupling coefficients are \(g e^{i\theta_1}\) and \(g e^{i\theta_2}\), respectively, with local coupling phases \(\theta_1\) and \(\theta_2\) for inducing some intriguing interference effects to the scattering properties as will be discussed below. With superconducting quantum devices, the local coupling phases can be introduced with Josephson loops threaded by external fluxes \([64]\).

Under the rotating wave approximation (RWA), the real-space Hamiltonian of the model can be written as
\((\hbar = 1\) hereafter\)

\[
H = H_w + H_a + H_I,
\]

\[H_w = \int_{-\infty}^{+\infty} dx \left[ a_R^\dagger(x) \left( \omega_0 + i v_g \frac{\partial}{\partial x} \right) a_R(x) + a_L^\dagger(x) \left( \omega_0 - i v_g \frac{\partial}{\partial x} \right) a_L(x) \right],\]

\[H_a = (\omega_e - i \frac{\gamma_e}{2}) |e\rangle \langle e|,\]

\[H_I = \int_{-\infty}^{+\infty} dx \left\{ \delta(x) g e^{i\theta_1} [a_R^\dagger(x) + a_L^\dagger(x)] |g\rangle \langle e| + \delta(x - d) g e^{i\theta_2} [a_R^\dagger(x) + a_L^\dagger(x)] |g\rangle \langle e| + H.c. \right\}.\]

(1)

Here \(H_w\) represents the free Hamiltonian of the waveguide modes with \(v_g\) being the group velocity of photons in the waveguide. \(a_{R,L} (a_{R,L}^\dagger)\) are the bosonic annihilation (creation) operators of the right-going and left-going photons in the waveguide, respectively; \(\omega_0\) is the frequency around which the dispersion relation of the waveguide mode is linearised \([1, 65]\). \(H_a\) is for the atom, where \(\omega_e\) describes the transition frequency between the ground state \(|g\rangle\) and the excited state \(|e\rangle\); \(\gamma_e\) is the external atomic dissipation rate due to the non-waveguide modes in the environment. \(H_I\) describes the interactions between the atom and the waveguide, where the Dirac delta functions \(\delta(x)\) and \(\delta(x - d)\) indicate that the atom-waveguide couplings occur at \(x = 0\) and \(x = d\), respectively. Besides, there is an accumulated phase \(\phi = kd\) of photons between two coupling points, where \(k\) is the renormalized wave vector that satisfies the linearised dispersion relation \(E = \omega_e + k v_g\) (\(E\) is the eigenenergy as determined later in Eq. (3)). The relevant physics discussed in this work is valid for \(E \sim \omega_e\), that is, \(k v_g \sim \omega_e - \omega_0\).

Considering that the total excitation number is conserved in RWA, the eigenstate of the system can be expressed in the single-excitation subspace as

\[
|\Psi\rangle = \int_{-\infty}^{+\infty} dx \left[ \Phi_R(x) a_R^\dagger(x) + \Phi_L(x) a_L^\dagger(x) \right] |0, g\rangle + u_e |0, e\rangle,
\]

(2)

where \(\Phi_{R,L}(x)\) are the density of probability amplitudes of creating the right-going and left-going photons at position \(x\), respectively; \(u_e\) is the excitation amplitude of the atom; \(|0, g\rangle\) denotes the vacuum state of the system. The probability amplitudes can be determined by solving

FIG. 1. (a) Schematic configuration of a two-level giant atom coupled to a waveguide at \(x = 0\) and \(x = d\), respectively, with individual local coupling phases \(\theta_1, 2\). (b) Two paths of a single photon propagating from port 1 to port 2 (left) or from port 2 to port 1 (right).
the eigenequation $H|\Psi\rangle = E|\Psi\rangle$, which leads to

$$E \Phi_R(x) = \left(\omega_0 - ivg \frac{\partial}{\partial x}\right) \Phi_R(x) + g\left[e^{i\theta_1} \delta(x) + e^{i\theta_2} \delta(x - d)\right] u_e,$$

$$E \Phi_L(x) = \left(\omega_0 + ivg \frac{\partial}{\partial x}\right) \Phi_L(x) + g\left[e^{i\theta_1} \delta(x) + e^{i\theta_2} \delta(x - d)\right] u_e,$$

$$E u_e = \left(\omega_e - i\frac{\gamma_e}{2}\right) u_e + ge^{-i\theta_1} \Phi_R(0) + \Phi_L(0) + ge^{-i\theta_2} \left[\Phi_R(d) + \Phi_L(d)\right].$$

Assuming that a single photon with wave vector $k$ ($k > 0$) is incident from port 1 of the waveguide. Then the wave functions $\Phi_R(x)$ and $\Phi_L(x)$ can be written in the forms

$$\Phi_R(x) = e^{i k x} \left\{\Theta(-x) + A[\Theta(x) - \Theta(x - d)] + i\theta(x - d)\right\},$$

$$\Phi_L(x) = e^{-i k x} \left\{i \Theta(-x) + B[\Theta(x) - \Theta(x - d)]\right\},$$

where $\Theta(x)$ is the Heaviside step function. Here, $t$ and $r$ denote the single-photon transmission and reflection amplitudes in the regions of $x > d$ and $x < 0$, respectively. We define $A$ and $B$ as the probability amplitudes for the right-going and left-going photons between the two coupling points ($0 < x < d$), respectively.

Substituting Eq. (4) into Eq. (3), one obtains

$$0 = -ivg(A - 1) + ge^{i\theta_1} u_e,$$

$$0 = -ivg(t - A)e^{i\phi} + ge^{i\theta_2} u_e,$$

$$0 = -ivg(r - B) + ge^{i\theta_1} u_e,$$

$$0 = ivg Be^{-i\phi} + ge^{i\theta_2} u_e,$$

$$0 = \frac{g}{2} e^{-i\phi} (A + B + r + 1) + \frac{g}{2} e^{-i\phi} \times (Ae^{i\phi} + Be^{-i\phi} + te^{i\phi}) - (\Delta + i\frac{\gamma_e}{2}) u_e$$

with $\Delta = E - \omega_e$ being the detuning of frequency between the incident photon and the atomic transition frequency. From the condition $E \sim \omega_e$, we work in the dispersive regime where the detuning is much smaller than the atomic transition frequency (deducting the offset $\omega_0$ in the linear dispersion) as $|\Delta/(\omega_e - \omega_0)| \ll 1$. Then the transmission and reflection amplitudes can be obtained from solving Eq. (5) as

$$t = \frac{\Delta + i\frac{\gamma_e}{2} - 2\Gamma e^{i\phi} \sin \phi}{\Delta + i\frac{\gamma_e}{2} + 2\Gamma (1 + e^{i\phi} \cos \theta)},$$

$$r = \left[\frac{2\Gamma (1 + e^{i\phi} \cos \theta) + 2e^{i\phi} \sin \phi}{\Delta + i\frac{\gamma_e}{2} + 2\Gamma (1 + e^{i\phi} \cos \theta)}\right] \left[1 + e^{-i(\theta + \phi)}\right].$$

where $\theta = \theta_2 - \theta_1$ is the phase difference between the two atom-waveguide coupling channels and $\Gamma = g^2/v_g$ is the rate of the atomic emission into the waveguide. Comparing Eqs. (6a) and (7a). On one hand, the trans-
mission amplitudes $t$ and $t'$ share the same denominator that is an even function of $\theta$. On the other hand, the numerators of $t$ and $t'$ in Eqs. (6a) and (7a) can be rewritten as

$$
\Delta - 2\Gamma \sin \phi \cos \theta + i(\gamma_e/2 - 2\Gamma \sin \phi \sin \theta),
$$

$$
\Delta - 2\Gamma \sin \phi \cos \theta + i(\gamma_e/2 + 2\Gamma \sin \phi \sin \theta).
$$

Equation (8) clearly shows that nonreciprocal single-photon transmissions ($|t|^2 \neq |t'|^2$) can be achieved only if a finite external atomic dissipation rate is taken into account ($\gamma_e > 0$). This can be observed by the transmission spectra shown in Figs. (c) and (d).

When $\gamma_e = 0$, Fig. 3(a) depicts the transmission rates $T_{1\rightarrow2}$ and $T_{2\rightarrow1}$ versus the detuning $\Delta$ with various $\theta$.

For $\theta = \pi/2$, we find $T_{1\rightarrow2} = T_{2\rightarrow1} = 1$ over the whole range of the detuning, implying that reflections are prevented for both directions. For $\theta = \pi$, however, the transmission spectrum exhibits the Lorentzian line shape with phase-dependent Lamb shift and linewidth (decay rate) [21]. In both cases ($\theta = \pi/2, \pi$), the transmission rates are reciprocal, yet the atomic excitation probabilities are different as will be discussed below. When $\gamma_e \neq 0$, as shown in Figs. (b) and (c), the scattering becomes nonreciprocal if $\theta = \pi/2$; however, with $\theta = \pi$, the scatterings are still reciprocal even in the presence of the external dissipation. The yellow dot-dashed, red dotted, and blue dashed lines in Fig. 3(d) depict the contrast ratio

$$
I = \frac{T_{2\rightarrow1} - T_{1\rightarrow2}}{T_{2\rightarrow1} + T_{1\rightarrow2}}.
$$

versus the coupling phase difference $\theta$ with different atomic dissipation rates. It can be seen that this nonreciprocal transmission behaviour is phase-dependent. Indeed, the above behaviours are easy to be analysed with Eqs. (6a) and (7a).

Furthermore, the underlying physics of the reciprocal and nonreciprocal scatterings can be understood via examining the atomic excitation by the single photon. To this end, we define the contrast ratio $D$ of the atomic excitation probabilities for two opposite propagating directions as

$$
D = \frac{|u_{e_{1\rightarrow2}}|^2 - |u_{e_{1\rightarrow1}}|^2}{|u_{e_{2\rightarrow1}}|^2 + |u_{e_{1\rightarrow2}}|^2},
$$

with

$$
u_{e_{1\rightarrow2}} = \frac{t - 1}{-i \frac{\gamma_e}{\gamma} e^{i\theta_1} + e^{i(\theta_2 + \phi)}},$$

$$
u_{e_{2\rightarrow1}} = \frac{t' - 1}{-i \frac{\gamma_e}{\gamma} e^{i\theta_2} + e^{i(\theta_1 + \phi)}.}
$$

According to Eqs. (6a) and (7a), parameters $t - 1$ and $t' - 1$ have the same denominator containing $\gamma_e$ but different numerators without $\gamma_e$. Furthermore, because the denominator that contains $\gamma_e$ is eliminated when calculating Eq. (10), the contrast ratio $D$ is independent of dissipation rate $\gamma_e$. Note that the contrast ratio $D$ can be used to capture the difference of the atomic excitation probabilities for opposite directions even if the eigenstate Eq. (2) is unnormalized.
the reflection rates

\( \theta = \pi/2 \) and \( \phi_0 = \pi/2 \).

We plot in Fig. 3(d) the contrast ratio \( D \) (black solid line) as a function of the phase difference \( \theta \) with \( \phi_0 = \pi/2 \). For \( \theta = \pi/2 \), \( D = -1 \) means that the atom can only be excited by the left-incident photon, and thus the atom-waveguide interaction becomes ideally chiral [64, 66]. In this case, the right-incident photon is guided transparently because it does not interact with the atom. While in the Markovian regime, the reflections are lacking for both directions under the ideal chiral coupling [16, 43], this is in fact not true in the non-Markovian regime as will be discussed in Sec. IIIB. For \( \theta = 3\pi/2 \), \( D = 1 \) corresponds to the ideal chiral case [32]. When the atom can only be excited by the right-incident photon, for other cases of \( D = 0 \) and \( 0 < |D| < 1 \), \( \phi_0 = \pi/2 \) in the non-Markovian regime is able to simulate a chiral atom-waveguide system. However, in the non-Markovian regime, due to the non-Markovian features in the transmission and reflection spectra [23, 61]. Here we just consider our system in the non-Markovian regime and demonstrate the reflection with \( \phi_0 = \pi/2 \) and \( \theta = \pi/2 \). Note that the reflection is totally prevented in the small-atom case with an ideal chiral coupling, which has been demonstrated in Ref. [16].

We plot in Fig. 4 the reflection rates \( R = |r|^2 \) for the left-incident photon in the Markovian and non-Markovian regimes. The yellow solid curve shows that the reflection in the Markovian regime disappears completely. Such a reflectionless behavior occurs in the case where the atom-waveguide couplings are nonchiral and nondiagonal chirality [64, 66]. In addition, we demonstrate in Fig. 3(e) that the propagation phase \( \Delta \) is also sensitive to the propagating phase. This provides an alternative way to tune the chiral/nondiagonal scattering on demand.

The results above can also be interpreted from the aspect of Hermitian and non-Hermitian scattering centers [67–69]. In our system with \( \gamma_e = 0 \) (\( \gamma_e \neq 0 \)), the giant atom can be regarded as a Hermitian (non-Hermitian) scattering center of the Aharonov-Bohm structure supporting two spatial interference paths. For the Hermitian case, the scattering remains reciprocal; however, when introducing an imaginary potential, e.g., the external atomic dissipation, the combination of the non-Hermiticity and the broken time-reversal symmetry gives rise to nonreciprocal scatterings [67, 68]. It is noted that, as discussed in the case in Figs. 3(b) and 3(c) \( \theta = \pi \), not all non-Hermitian scattering centers can demonstrate nonreciprocal transmissions. The exceptions include, e.g., \( \mathcal{P}, \mathcal{T}, \) or \( \mathcal{PT} \)-symmetric scattering centers [67, 69]. In our model, although the giant atom can exhibit chiral spontaneous emission corresponding to the time-reversal symmetry breaking if \( \theta \neq n\pi \) (\( n \) is an arbitrary integer) [64], the scatterings are still reciprocal unless the additional non-Hermiticity (such as external dissipations) are introduced.

2. Non-Markovian regime

With nontrivial local coupling phases, as demonstrated above, the current giant-atom model (in the Markovian regime) is able to simulate a chiral atom-waveguide system. However, one important characteristic of the giant atom is the peculiar scattering behaviours arising in the non-Markovian regime, where the propagating phase accumulation \( \phi = \phi_0 + \tau \Delta \) is sensitive to the detuning \( \Delta \) due to the large enough \( \tau \) that is comparable to or larger than the lifetime of the atom [27]. Such a detuning-dependent phase will undoubtedly result in the non-Markovian features in the transmission and reflection spectra [23, 61]. Here we just consider our system in the non-Markovian regime and demonstrate the reflection with \( \phi_0 = \pi/2 \) and \( \theta = \pi/2 \). Note that the reflection is totally prevented in the small-atom case with an ideal chiral coupling, which has been demonstrated in Ref. [16].

In this section, we extend the giant-atom model to a multi-level version and demonstrate the possibility of realising nonreciprocal scatterings without the additional non-Hermiticity (i.e., external atomic dissipation). Specifically, we introduce here an additional atomic transition coupled to other waveguide modes. As shown in Fig. 5(a), we propose a \( \nabla \)-type giant atom coupled to two waveguides via two different atomic transitions, respectively. Each transition is coupled to the correspond-
ing waveguide at two points, and an external microwave field is applied to drive the magnetic dipole transition between two excited states [70]. Such system allows for, without the help of external dissipation, high-efficiency single-photon routing and circulating. Furthermore, at the end of this section, we will also consider a Δ-type scheme and compare the differences between these two three-level structures.

As shown in Fig. 5(a), the atomic transition \( |e_1\rangle \leftrightarrow |g\rangle \) of frequency \( \omega_{e_1} \) is coupled to waveguide \( W_a \) with complex coupling coefficient \( g_1 \text{e}^{i\theta_1} \) at two separated points \( x = 0 \) and \( x = d_a \), respectively; the transition \( |e_2\rangle \leftrightarrow |g\rangle \) of \( \omega_{e_2} \) is coupled to \( W_b \) with \( g_2 \text{e}^{i\theta_2} \) at \( x = 0 \) and \( x = d_b \), respectively. The excited states \( |e_1,2\rangle \) are coupled to an external coherent field of Rabi frequency \( \Omega \) and initial phase \( \alpha \). The atom is initialized on the ground state \( |g\rangle \).

The Hamiltonian of the \( \nabla \)-type giant atom coupled to two waveguides can be written as

\[
H' = H'_w + H'_a + H'_1,
\]

\[
H'_w = \int_{-\infty}^{\infty} \! \! dx \left[ a_{L}^\dagger(x) \left( \omega_0 + iv_g \frac{\partial}{\partial x} \right) a_L(x) + a_{R}^\dagger(x) \left( \omega_0 - iv_g \frac{\partial}{\partial x} \right) a_R(x) \right] + \int_{-\infty}^{\infty} \! \! dx \left[ b_{L}^\dagger(x) \left( \omega_0 + iv_g \frac{\partial}{\partial x} \right) b_L(x) + b_{R}^\dagger(x) \left( \omega_0 - iv_g \frac{\partial}{\partial x} \right) b_R(x) \right],
\]

\[
H'_a = \left( \omega_{e_1} - i \frac{v_{e_1}}{2} \right) |e_1\rangle \langle e_1| + \left( \omega_{e_2} - i \frac{v_{e_2}}{2} \right) |e_2\rangle \langle e_2| + \left( \Omega \text{e}^{i\theta_1} \right) |e_1\rangle \langle e_2| + \text{H.c.},
\]

\[
H'_1 = \int_{-\infty}^{\infty} \! \! dx \left\{ \delta(x - d) g_1 \text{e}^{i\theta_1} \left[ a_{R}^\dagger(x) + a_{L}^\dagger(x) \right] |g\rangle \langle e_1| + \delta(x + d) g_2 \text{e}^{i\theta_2} \left[ b_{R}^\dagger(x) + b_{L}^\dagger(x) \right] |g\rangle \langle e_2| + \delta(x - d) g_2 \text{e}^{i\theta_2} \left[ b_{R}^\dagger(x) + b_{L}^\dagger(x) \right] |g\rangle \langle e_2| + \text{H.c.} \right\}.
\]

where \( a_{R,L}/b_{R,L}^\dagger \) (\( a_{R,L}/b_{R,L}^\dagger \)) annihilates (creates) right-going and left-going photons in the waveguide \( W_a/W_b \), respectively. In the single-excitation subspace, the eigenstate of the system can be expressed as

\[
|\Psi\rangle = \int_{-\infty}^{\infty} \! \! dx \left[ \Phi_{aR}(x) a_{R}^\dagger(x) + \Phi_{aL}(x) a_{L}^\dagger(x) \right]|0, g\rangle + u_{e_1}|0, e_1\rangle + u_{e_2}|0, e_2\rangle,
\]

where \( \Phi_{aR,L} \) (\( \Phi_{bR,L} \)) are the probability amplitudes of creating the right-going and left-going photons in \( W_a \) (\( W_b \)), respectively.

Assuming that a photon with wave vector \( k_a \) is emitted from port 1 of \( W_a \), the probability amplitudes for this case can be written as

\[
\Phi_{aR}(x) = e^{ik_a x} \left\{ \Theta(-x) + M \left[ \Theta(x) - \Theta(x - d_a) \right] \right\},
\]

\[
\Phi_{aL}(x) = e^{-ik_a x} \left\{ s_{1\rightarrow2} \Theta(x - d_a) \right\},
\]

\[
\Phi_{bR}(x) = e^{ik_a x} \left\{ s_{1\rightarrow4} \Theta(x - d_b) \right\},
\]

\[
\Phi_{bL}(x) = e^{-ik_a x} \left\{ s_{1\rightarrow3} \Theta(x) + W \left[ \Theta(x) - \Theta(x - d_b) \right] \right\},
\]

where the wave vectors \( k_a = (E' - \omega_0)/v_g \) with the eigenenergy \( E' \) in \( W_a \) and \( k_b = k_a + (\omega_{e_2} - \omega_{e_1})/v_g \) in \( W_b \). When excited to state \( |e_1\rangle \) by the incident photon from port 1, the atom can either re-emit a photon with the same frequency to \( W_a \) via decaying back to state \( |g\rangle \) directly, or radiate a photon with frequency \( \omega_{e_2} \) to \( W_b \) via first transferring from state \( |e_1\rangle \) to state \( |e_2\rangle \) due to the external driving and then decaying to state \( |g\rangle \) [62, 71].

If a photon with wave vector \( k_b \) is sent from port 4 of \( W_b \), the probability amplitudes can be written as

\[
\Phi_{aR}(x) = e^{ik_a x} \left\{ s_{4\rightarrow2} \Theta(x - d_a) + M' \left[ \Theta(x) - \Theta(x - d_a) \right] \right\},
\]

\[
\Phi_{aL}(x) = e^{-ik_a x} \left\{ N' \left[ \Theta(x) - \Theta(x - d_a) \right] + s_{4\rightarrow1} \Theta(-x) \right\},
\]

\[
\Phi_{bR}(x) = e^{ik_a x} \left\{ Q' \left[ \Theta(x) - \Theta(x - d_b) \right] + s_{4\rightarrow3} \Theta(x - d_b) \right\},
\]

\[
\Phi_{bL}(x) = e^{-ik_a x} \left\{ \Theta(x - d_b) + W' \left[ \Theta(x) - \Theta(x - d_b) \right] + s_{4\rightarrow3} \Theta(-x) \right\}.
\]
1. Nonreciprocal scattering

For simplicity, we start by supposing $\phi_b = \theta_b = \theta_1 = \theta_2 = 0$, i.e., the transition $|g\rangle \leftrightarrow |e_2\rangle$ is coupled to $W_b$ at a single point. Then, the scattering probabilities can be calculated from $S_{1\rightarrow 2} = |s_{1\rightarrow 2}|^2$ and $S_{1\rightarrow 3(4)} = |s_{1\rightarrow 3(4)}|^2$ with the scattering amplitudes given by

$\begin{align*}
S_{1\rightarrow 2} &= \frac{\Delta' + i\frac{\gamma_{e_2}}{2} - \Omega^2 f - 2\Gamma_1 e^{i\theta} \sin \phi_0}{\Delta' + i\frac{\gamma_{e_2}}{2} - \Omega^2 f + 2\Gamma_1 (1 + e^{i\theta} \cos \theta)}, \\
S_{1\rightarrow 3(4)} &= \frac{g_2 \Omega e^{-i\alpha} (s_{1\rightarrow 2} - 1)}{g_1 (\Delta' + \gamma_{e_2} / 2 + i\Gamma_2) \left[ e^{i\theta_1} + e^{i(\theta_2 - \phi_0)} \right]},
\end{align*}$

with the detuning $\Delta' = E' - \omega_\epsilon$ and the atomic emission rates $\Gamma_{1,2} = g_{1,2}^2 / v_g$. As discussed above, we make the substitution $\phi_0 = k_a d_a \simeq \phi_0$ in the Markovian regime. Likewise, one can also obtain $S_{2\rightarrow 1} = |s_{2\rightarrow 1}|^2$ and $S_{2\rightarrow 3(4)} = |s_{2\rightarrow 3(4)}|^2$ with

$\begin{align*}
S_{2\rightarrow 1} &= \frac{\Delta' + i\frac{\gamma_{e_1}}{2} - \Omega^2 f - 2\Gamma_1 e^{-i\theta} \sin \phi_0}{\Delta' + i\frac{\gamma_{e_1}}{2} - \Omega^2 f + 2\Gamma_1 (1 + e^{-i\theta} \cos \theta)}, \\
S_{2\rightarrow 3(4)} &= \frac{g_2 \Omega e^{-i\alpha} (s_{2\rightarrow 3} - 1)}{g_1 (\Delta' + \gamma_{e_1} / 2 + i\Gamma_2) \left[ e^{i\theta_1} + e^{i(-\theta_2 - \phi_0)} \right]},
\end{align*}$

which are achieved via exchanging $\theta_1$ and $\theta_2$ in Eq. (16).

It is found that $s_{4\rightarrow 1} = s_{2\rightarrow 3}$ and thus $S_{4\rightarrow 1} = |s_{4\rightarrow 1}|^2 = S_{2\rightarrow 3}$. Compared with Eq. (6a) and Eq. (7a) of the two-level giant atom, both $s_{1\rightarrow 2}$ and $s_{2\rightarrow 1}$ include an additional coupling term $\Omega^2 f$ with

$$f = \frac{1 - i(\gamma_{e_2} + \Gamma_2)}{\Delta'^2 + (\frac{\gamma_{e_2}}{2} + \Gamma_2)^2}.$$ 

which describes the photon transfer from $W_a$ to $W_b$. It can be seen from Eqs. (16)-(18) that, in contrast to the two-level giant-atom scheme, the transmission between ports 1 and 2 in $W_a$ is nonreciprocal even if the external dissipations are not considered. In fact, for $\gamma_{e_1} = \gamma_{e_2} = 0$, the imaginary part of $f$ describing the decay of $|e_2\rangle \rightarrow |g\rangle$ into $W_b$ plays the role of an external dissipation for the transition $|e_1\rangle \rightarrow |g\rangle$.

It is worth noting that the scattering probabilities of the $\nabla$-type system are independent of the phase $\alpha$ of the external coherent field $\Omega$ in spite of the closed-loop atom-level structure. This is because the $\nabla$-type atom cannot provide the inner two-path quantum interference. For instance, when excited to state $|e_1\rangle$ by an incident photon from port 1, the atom may be pumped to state $|e_2\rangle$ by the external field $\Omega$ and then return to state $|g\rangle$ after emitting a photon into $W_b$, which is the only path for the photon transferring from $W_a$ to $W_b$. This is radically different from the $\Delta$-type structure as will be discussed in Sec. IIC. In fact, the photon cannot be routed from $W_a$ to $W_b$ in the absence of the field $\Omega$, implying that the $\nabla$-type three-level giant atom reduces to a two-level one. This is also consistent with the fact that $S_{1\rightarrow 3(4)} = S_{2\rightarrow 3(4)} = 0$ when $\Omega = 0$.

Figure 6 shows the single-photon scattering spectra as functions of the detuning $\Delta'$ and the phase difference $\theta$. As discussed above, it can be seen from Figs. 6(a) and 6(b) that the nonreciprocal scattering can still be realised in $W_a$ ($S_{1\rightarrow 2} \neq S_{2\rightarrow 1}$) with $W_b$ playing the role of the external thermal reservoir in the two-level scheme as analysed above. According to the conclusion in Sec. II, for $\theta \neq n\pi$ and $\phi_0 = \pi / 2 + 2n\pi$, the excitation probabilities $|\alpha_{e_1}|^2$ for two opposite directions are unequal, i.e., the effective interaction between the atom and $W_a$ is chiral. Then, as shown in Figs. 6(c) and 6(d), the non-reciprocal scattering between ports 1 and 4 can be led to by the chiral coupling, since the scattering probability $S_{1\rightarrow 4}$ ($S_{2\rightarrow 1}$) is related to the coupling between the atomic transition $|e_1\rangle \leftrightarrow |g\rangle$ and the right-going (left-going) mode in $W_a$. When $\theta = \pi / 2 (3\pi / 2)$, $S_{1\rightarrow 4}$ ($S_{2\rightarrow 1}$) approaches 0.5 and $S_{2\rightarrow 1}$ ($S_{1\rightarrow 4}$) falls to 0. This corresponds to the ideal chiral case where the atom is only coupled to the right-going (left-going) modes effectively in $W_a$. When $\theta = \pi$, the scatterings between ports 1 and 4 are reciprocal, similar to the results of the non-chiral case in Sec. II.

2. Chiral scattering

Next, we turn to study another kind of asymmetric scattering phenomenon proposed recently called chiral scattering. Specifically, the transmission from port 1 to port 4 and that from port 2 to port 3 are different. Quantitatively, the chiral scattering can be evaluated by the chirality defined as

$$C = \frac{S_{1\rightarrow 4} - S_{2\rightarrow 3}}{S_{1\rightarrow 4} + S_{2\rightarrow 3}}.$$
Figure 7(a) shows the chirality as a sinusoidal function of the phase difference $\theta$. In view of this, chiral scatterings can be observed as long as $\theta \neq n\pi$, where the chirality $C \neq 0$ means $S_{1\to4} \neq S_{2\to3}$. This can be further verified by the scattering spectra as shown in Figs. 7(b) and 7(c).

Note that $C = 1$ ($C = -1$) corresponding to $\theta = \pi/2$ ( $\theta = 3\pi/2$), implies that only the scattering from port 2 (1) to port 3 (4) is prevented, as shown in Fig. 7(b) [Fig. 7(c)].

The underlying physics of the chiral scattering can also be attributed to the difference between the atomic excitation probabilities for two incident directions as discussed above. The excitation probabilities $|u_{1\to}|^2$ by the photon incident from port 1 and port 2 can be unequal, and thus the atom is pumped from $|e_1\rangle$ to $|e_2\rangle$ with unequal probabilities. This leads to different probabilities of routing photons from $W_a$ to $W_b$. Furthermore, as shown in Fig. 7, the chiral scattering scheme here shows the in-situ tunability that the scattering chirality can be controlled by tuning the phase difference $\theta$.

More interestingly, the $\nabla$-type giant atom is also a promising candidate of realizing a single-photon circulator. When turning on the external field and setting $\theta = \pi/2$ and $\theta' = 3\pi/2$, the two waveguides are coupled to the atom with ideal chiral couplings in opposite manners, respectively. That is to say, the atom is only based on the asymmetric scatterings above. Specifically, one can send a single photon deterministically from port 1 to one of the three ports on demand. Note that $S_{2\to1} = S_{3\to4} = 1$ over the whole frequency range and the router and circulator can run with very high efficiency.

In this subsection, we would like to demonstrate how to realise a single-photon targeted router and circulator based on the asymmetric scatterings above. Specifically, one can send a single photon deterministically from port 1 to one of the three ports on demand. Note that $S_{2\to1} = S_{3\to4} = 1$ over the whole frequency range and the router and circulator can run with very high efficiency.

3. Targeted router and circulator

In this subsection, we would like to demonstrate how to realise a single-photon targeted router and circulator based on the asymmetric scatterings above. Specifically, one can send a single photon deterministically from port 1 to one of the three ports on demand. Note that $S_{2\to1} = S_{3\to4} = 1$ over the whole frequency range and the router and circulator can run with very high efficiency.

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The mechanism of the targeted router can be understood in the quantum networks [54, 55].
4. Comparison with the \( \Delta \)-type scheme

Finally, we consider a \( \Delta \)-type giant-atom scheme where the \( \nabla \)-type atom in Fig. 5(a) is replaced by a \( \Delta \)-type one in Fig. 5(b) and compare the single-photon scatterings of these two schemes. The \( \Delta \)-type structure is constructed with an external coherent field \( \epsilon e^{i\phi} \) which couples the two ground states \(|g_1, 2\rangle |) of a \( \Delta \)-type atom that has been broadly studied to demonstrate quantum interference phenomena, such as coherent population trapping [73] and electromagnetically induced transparency [74].

For the \( \Delta \)-type giant-atom system, the Hamiltonians of the atom and the atom-waveguide interaction become

\[
H_a' = \left( \omega_{g_2} - \frac{i\gamma_{g_2}}{2} \right) |g_2\rangle \langle g_2| + \left( \omega_c - \frac{i\gamma_c}{2} \right) |c\rangle \langle c| \\
+ (\epsilon e^{i\phi} |g_1\rangle |g_2\rangle + \text{H.c.},
\]

\[
H'_{L} = \int_{-\infty}^{+\infty} dx \left\{ \delta(x) |g_1\rangle e^{i\theta_1} [a_{L}^\dagger(x) + a_{L}^\dagger(x)] |g_1\rangle \langle e| \\
+ \delta(x - d) |g_2\rangle e^{i\theta_2} [a_{R}^\dagger(x) + a_{L}^\dagger(x)] |g_1\rangle \langle e| \\
+ \delta(x) |g_2\rangle e^{i\theta_3} [b_{R}^\dagger(x) + b_{L}^\dagger(x)] |g_2\rangle \langle e| + \text{H.c.} \right\}.
\]

The single-excitation eigenstate of the system takes the form

\[
|\Psi\rangle = \int_{-\infty}^{+\infty} dx \left\{ [\Phi_{aR}(x)a_{R}^\dagger(x) + \Phi_{aL}(x)a_{L}^\dagger(x)] |0, g_1\rangle \\
+ [\Phi_{bR}(x)b_{R}^\dagger(x) + \Phi_{bL}(x)b_{L}^\dagger(x)] |0, g_2\rangle \} + u_e |0, e\rangle.
\]

With the same procedure above (see Supplementary for more details), one can obtain the scattering probabilities in this case.

Setting the atom on the ground state \(|g_1\rangle \) initially, we plot in Fig. 9 the scattering spectra of \( \tilde{S}_{1\rightarrow 4} \) and \( \tilde{S}_{4\rightarrow 1} \). It is worth noting that, even in the absence of the local couplings, the nonreciprocal scatterings for the two phase-sensitive closed-loop three-level giant-atom structures with tunable phase on each atom-waveguide coupling. We found that the atomic excitation in the two-level giant-atom structure depends on the propagation direction of waveguide modes and can be tuned by the nontrivial coupling phase difference. In such scenarios, our two-level giant atom in the Markovian regime is equivalent to a two-level small atom chirally coupled to the waveguide mode. However, it is worth noting that the nonreciprocal and chiral scatterings without external dissipations is achieved by the local coupling phases and the non-Hermiticity induced by the external atomic dissipation due to the surrounding non-waveguide modes. Moreover, in the non-Markovian regime, the reflection spectra exhibit peculiar features with multiple reflection peaks that are absent in the chiral small-atom case.

For exploring more interesting asymmetric scatterings properties and applications with such giant-atom structures, we have extended the two-level structure to the three-level \( \nabla \)-type and \( \Delta \)-type ones coupled to two waveguides via different atomic transitions. We found that, for the atomic transition coupled to one waveguide, the nonreciprocity of the \( \nabla \)-type case stems from the different atomic structures and the transition to the other waveguide can serve as the external field. However, the nonreciprocity of \( \Delta \)-type structure is constructed with an external coherent field \( \epsilon e^{i\phi} \) which couples the two ground states \(|g_1, 2\rangle |) of a \( \Delta \)-type atom that has been broadly studied to demonstrate quantum interference phenomena, such as coherent population trapping [73] and electromagnetically induced transparency [74].

In summary, we have investigated step-by-step the conditions of single-photon nonreciprocal and chiral scatterings in the two-level and three-level giant-atom structures with tunable phase on each atom-waveguide coupling. We found that the atomic excitation in the two-level giant-atom structure depends on the propagation direction of waveguide modes and can be tuned by the nontrivial coupling phase difference. In such scenarios, our two-level giant atom in the Markovian regime is equivalent to a two-level small atom chirally coupled to the waveguide mode. However, it is worth noting that the nonreciprocal and chiral scatterings without external dissipations is achieved by the local coupling phases and the non-Hermiticity induced by the external atomic dissipation due to the surrounding non-waveguide modes. Moreover, in the non-Markovian regime, the reflection spectra exhibit peculiar features with multiple reflection peaks that are absent in the chiral small-atom case.

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efficient single-photon optical elements for quantum net- work engineering and optical communications.

In this theoretical work, the methods used are solving the stationary Schrödinger equation with the Hamiltonian and single-excitation eigenstate (as described in the main text [Eq. (3)] and Supplementary).

All data are available in the main text or in the sup- plementary materials.

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C.Y.T. conceived the study, performed the calculation, and wrote the first version of the manuscript. D.L., G.L.Z., and Z.Y. checked the calculations, analyzed the data, and redrafted the manuscript. W.Z.H., L.Y., and W.J.H. helped in the interpretation of the results and in the writing of the manuscript. All authors discussed the results and reviewed the manuscript.

This work is supported by the National Natural Science Foundation of China (Grants No. 12074030 and No. U1930402), the Fundamental Research Funds for the Central Universities (Grant No. 2412019FZ045), and the Science Foundation of the Education Department of Jilin Province during the 14th Five-Year Plan Period (Grant No. JJKH20211279KJ).

The authors declare no competing interests.

This work is supported by the National Natural Science Foundation of China (Grants No. 12074030 and No. U1930402), the Fundamental Research Funds for the Central Universities (Grant No. 2412019FZ045), and the Science Foundation of the Education Department of Jilin Province during the 14th Five-Year Plan Period (Grant No. JJKH20211279KJ).
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