Anomalous spontaneous emission dynamics at chiral exceptional points

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An open quantum system operated at the spectral singularities where dimensionality reduces, known as exceptional points (EPs), demonstrates distinguishing behavior from the Hermitian counterpart. Here, we present an analytical description of local density of states (LDOS) for microcavity featuring chiral EPs, and unveil the anomalous spontaneous emission dynamics from a quantum emitter (QE) due to the non-Lorentzian response of EPs. Specifically, we reveal that a square Lorentzian term of LDOS contributed by chiral EPs can destructively interfere with the linear Lorentzian profile, resulting in the null Purcell enhancement to a QE with special transition frequency, which we call EP induced transparency. While for the case of constructive interference, the square Lorentzian term can narrow the linewidth of Rabi splitting even below that of bare components, and thus significantly suppresses the decay of Rabi oscillation. Interestingly, we further find that an open microcavity with chiral EPs supports atom-photon bound states for population trapping and decay suppression in long-time dynamics. As applications, we demonstrate the advantages of microcavity operated at chiral EPs in achieving high-fidelity entanglement generation and high-efficiency single-photon generation. Our work unveils the exotic cavity quantum electrodynamics unique to chiral EPs, which opens the door for controlling light-matter interaction at the quantum level through non-Hermiticity, and holds great potential in building high-performance quantum-optics devices.

I. INTRODUCTION

Exceptional points (EPs), the spectral degeneracies where two or more eigenvalues and the associated eigenstates simultaneously coalesce, are the central concept in non-Hermitian physics [1–5]. A plethora of intriguing effects and exotic phenomena emerge around EPs due to their nontrivial topological properties and the dimensionality reduction, including ultrasensitive sensing [6–12], laser mode selection [13, 14], chiral mode conversion [15–17], and unidirectional invisibility [18, 19]. Harnessing these peculiar features of non-Hermitian degeneracies for building novel devices with unprecedented performance has been experimentally demonstrated in various classical dissipative platforms, ranging from nanophotonics [2, 20–23], acoustics [24], to macroscopic facilities, such as fiber network [25], electric circuits [26] and heat diffusive system [27]. In recent years, great efforts have been dedicated to accessing the quantum EPs by implementing the non-Hermiticity in quantum systems [11, 17, 28–39], and investigating the quantum states control through EPs [17, 31–33, 38, 40–42]. In this respect, pioneer works have shown the ability of EPs to tune the photon statistics [32, 33], enhance the sensitivity of quantum sensing [43, 44], and steer the evolution of single quantum system [17, 31, 38, 40–42, 45–47].

Despite these promising results, quantum effects of EPs beyond wave mechanics is still largely unexplored. Recently, the emission properties of a quantum emitter (QE) in electromagnetic environment supporting EPs attracts growing attention [38, 40, 42, 45], since the modification of spontaneous emission (SE) exhibits a non-Lorentzian feature around EPs, contrast to the Lorentzian response in conventional single-mode cavity, which can lead to a greater enhancement of SE rate [40, 42, 45]. It implies that the formation of EPs in nanophotonic structures significantly alters the local density of states (LDOS), which fully governs the interaction between a QE and arbitrary electromagnetic environment [48–50]. Tailoring the LDOS of electromagnetic environment is crucial for optimizing the performance of many practical applications, ranging from traditional optoelectronic devices like lasers [51, 52] and solar cells [53], to advanced quantum technologies, such as quantum light sources [54–56] and logical gate [57, 58]. Therefore, a LDOS theory that can capture the effects of EPs is of both fundamental and applied significance. The formalism of LDOS based on the classical Green function sep-
Bare microcavity Microcavity with chiral EPs and CCW modes with equal coupling rate quantum emitter (QE) is simultaneously coupled to both CW and counterclockwise (CCW) modes. A linearly polarized croring resonator supports the degenerate clockwise (CW) whispering-gallery-mode (WGM) cavity represented by \( \text{FIG. 1. Schematic diagrams of cavity QED models. (a) A Ljunctions is polarized one. The distance between two waveguide-cavity to another CCW mode through mirror symmetry. Accordingly, the linearly polarized QE is replaced by a circularly polarized one. The distance between two waveguide-cavity junctions is } L. \)

\[ \begin{align*}
\text{Bare microcavity (DP cavity)} & \quad \text{Microcavity with chiral EPs (EP cavity)} \\
\text{Equivalent cavity QED model of EP cavity} & \quad \text{(c)}
\end{align*} \]

arates the EPs contribution from the usual Lorentzian term \( [40, 41, 59]\), while it leaves the origin of EPs unclear, i.e., information of the underlying cavity resonances (frequency and decay rate) is ambiguous, the coupling parameters of coalescent cavity modes that forms EPs are implicit, and thus the EPs condition cannot be given explicitly. As a consequence, there are at least two obvious limitations that hinder the exploration of quantum effects of EPs using macroscopic quantum electrodynamics (QED) based on the classical Green function: the quick and accurate design of nanophotonic structures possessing EPs remains challenging, and the nonlinear regime of cavity QED involving multiphoton interaction is exclusive.

To overcome the aforementioned limitations of current LDOS theory for EPs and explore the quantum state control through EPs, in this work we propose an equivalent cavity QED model that can provide an intuitive and flexible quantum master equation approach for a class of novel microcavity with chiral EPs \([8, 11, 60–63]\), also known as exceptional surface, which is a collection of EPs in high-dimensional parameter space rather than a single EP. An analytical description of LDOS for such EP cavity is derived in Sec. II, and parameterized by the resonance frequencies, decay rates and dissipative coupling rate of discrete cavity modes, as well as the coupling rates between the cavity and the QE. Therefore, it can be used to determine the parameters in quantum master equation from LDOS of a given photonic structure obtained by electromagnetic simulations, and establishes an explicit connection with cavity QED at EPs. Furthermore, the formulation of SE spectrum is also obtained to understand the peculiar quantum dynamics at EPs. Our LDOS theory not only explains the previously known properties of SE dynamics related to EPs only with the special parameters choice \([41, 42]\), but also allows the entire formalization and exploring more extraordinary effects of EPs in open quantum systems.

Based on the proposed LDOS theory, in Sec. III we perform comprehensive study to show how the chiral EPs tailors the LDOS and affects the SE dynamics of a QE and compare to the conventional Lorentzian cavity. With the equivalent cavity QED model, we further find out the formation of a special kind of quantum state without decay at chiral EPs, and demonstrate its applications in quantum state control. Finally, we conclude in Sec. IV.

II. MODEL AND THEORY

In this section, we build the quantum model to describe the interaction between a QE and the microcavity with chiral EP (we call EP cavity hereafter). In traveling wave optical resonators, such as the whispering-gallery-mode (WGM) microcavity, the chiral EPs emerges via two unidirectionally coupled cavity modes. Different strategies have been proposed to implement the unidirectional coupling between the clockwise (CW) and the counterclockwise (CCW) modes of WGM microcavity, by introducing two or more Rayleigh-type nano-scatters \([8, 62]\) or coupling to a waveguide terminated by a mirror at the one side \([11, 42, 63]\). Despite different realizations, the underlying physical mechanism is essentially the same, where the unidirectional coupling is induced by auxiliary reservoir modes offered by either nano-scatters or waveguide. Here we employ the latter scheme for the sake of theoretical simplicity, but the results and conclusions presented in this work are suitable for other types of EP cavity.

Fig. 1(a) depicts a basic cavity QED system consisting of a QE and a WGM microcavity with infinite waveguide. The WGM microcavity has a decay rate \( \kappa \) due to the coupling to waveguide. Since there is no coupling between the CW and CCW modes, the cavity has two degenerate modes, i.e., is operated at diabolic points (DPs) \([60, 64]\), and thus called DP cavity hereafter. The QE is linearly polarized, and coupled to a pair of degenerate CW and CCW modes with equal coupling rate \( g \). The
system is described by the two-mode Jaynes-Cummings (JC) model. While by coupling the WGM microcavity to a semi-infinite waveguide with a mirror at the right end, a unidirectional coupling from CCW mode to CW mode is created, as Fig. 1(b) illustrates. This photonic structure supports chiral EPs, and offers great tunability to SE rate of a QE.

In order to develop an intuitive cavity QED model for this novel EP cavity, we transform the system to an equivalent one through mirror symmetry, as shown in Fig. 1(c). The original CCW mode is flipped to a mirrored CCW mode, and the corresponding left-handed emission of QE is converted to the right-handed emission into the mirrored CCW mode as well. As a result, the linearly polarized QE becomes circularly polarized in the equivalent cavity QED model. Two CCW modes in the equivalent model constitute a cascaded system [65–67], the quantum dynamics of a QE is described by the extended cascaded quantum master equation (ℏ = 1)

\[ \dot{\rho} = -i \{ H_0 + H_I, \rho \} + \gamma \mathcal{L} [\sigma_-] \rho + \kappa \mathcal{L} [c_L] \rho + \kappa \mathcal{L} [c_R] \rho + \kappa |r| \left( e^{i\phi} \left[ c_L \rho, c_R^\dagger \right] + e^{-i\phi} \left[ c_R, c_L^\dagger \rho \right] \right) \]

with the free Hamiltonian

\[ H_0 = \omega_0 \sigma_+ \sigma_- + \omega_c c_L^\dagger c_L + \omega_c c_R^\dagger c_R \]

and the interaction Hamiltonian

\[ H_I = g \left( e^{-i\phi} c_L^\dagger \sigma_- + h.c. \right) + g \left( e^{i\phi} c_R^\dagger \sigma_- + h.c. \right) \]

where \( \sigma_- \) is the lowering operator of QE with transition frequency \( \omega_0 \) and SE rate \( \gamma \) in homogeneous medium. \( \gamma = \mu^2 \omega_0^3 n_b / (3\pi\varepsilon_0 c^3) \), where \( \varepsilon_0 \) is the permittivity of vacuum, \( \mu \) is the dipole moment of QE, and \( n_b \) is the refractive index of background medium. \( c_L^\dagger c_R \) is the bosonic annihilation operator for CCW mode in the left/right cavity with resonance frequency \( \omega_c \). \( \kappa \) is the dissipative coupling rate between the cavity modes and the waveguide. The intrinsic decay of two CCW modes is omitted, considering that the evanescent coupling to waveguide dominates the cavity dissipation. \( \mathcal{L}[\rho] = \rho \mathcal{O} \mathcal{O}^\dagger - \{ \mathcal{O}^\dagger \mathcal{O}, \rho \} / 2 \) is the Liouvillian superoperator for operator \( \mathcal{O} \). The second line of Eq. (1) describes the unidirectional coupling from the left CCW mode to the right CCW mode, where \( |r| \) is the reflectivity of mirror, and \( \varphi = \beta L \) is the phase factor, with \( \beta \) and \( L \) being the propagation constant of waveguide and the distance between two cavities, respectively. In the interaction Hamiltonian, the phase factor \( \phi \) originates from the traveling wave nature of WGM modes and depends on the azimuthal orientation of QE location [64, 68].

From Eq. (1), we take the operator expectation values to obtain the equations of motion

\[ \frac{d}{dt} \langle \tilde{p} \rangle = -i \mathcal{M} \langle \tilde{p} \rangle \]

where \( \mathcal{M} = \begin{pmatrix} \omega_c - i\frac{\kappa}{2} & 0 & ge^{-i\phi} \\ -ik|r|e^{i\phi} & \omega_c - i\frac{\kappa}{2} & ge^{i\phi} \\ ge^{i\phi} & ge^{-i\phi} & \omega_0 - i\frac{\kappa}{2} \end{pmatrix} \)

The above equations do not include the time delay of itinerant photon. This is justified if the propagation time \( L/v_\gamma \), where \( v_\gamma \) is the group velocity of waveguide, is much smaller than the timescale of system evolution set by \( \min \{ g^{-1}, \kappa^{-1}, \gamma^{-1} \} \). In the following study, we consider the optical frequency and focus on the strong coupling regime, and thus the timescale of system evolution is determined by \( g^{-1} \). For WGM cavity, \( g \) is typically smaller than 1meV, and we evaluate that \( g^{-1} \gg L/v_\gamma \) for \( L < 20 \mu m \); in this case, we can assume that the photon propagates without time delay.

To study the quantum dynamics of a QE located in a realistic photonic structure, especially for the interactions involving multiphoton process, we need to determine the system parameters in quantum master equation from LDOS obtained by electromagnetic simulation. Therefore, an analytical LDOS theory is crucial to extract the system parameters. The normalized LDOS of electromagnetic environment, i.e., Purcell factor, is expressed as \( F_P(\omega) = |J(\omega) + J_0(\omega)| / J_0(\omega) \), where \( J(\omega) \) and \( J_0(\omega) = \gamma/2\pi \) are the spectral density of photonic structures and background medium (free space), respectively. For EP cavity shown in Fig. 1(b), the spectral density of cavity is given by \( J(\omega) = \int_0^\infty d\omega e^{i\omega \tau} g^2 \left[ \left( c_L^\dagger (0) + e^{-i2\phi} c_R^\dagger (0) \right) \left[ c_L(\tau) + e^{i2\phi} c_R(\tau) \right] \right] \)

[69–71], where the environmental correlation functions \( \langle c_L^\dagger (0)c_L(\tau) \rangle \), \( \langle c_L^\dagger (0)c_R(\tau) \rangle \) and \( \langle c_R^\dagger (0)c_R(\tau) \rangle \) can be calculated by applying the quantum regression theorem \( \langle O_i(t)O_j(t+\tau) \rangle = \text{Tr} \left[ O_i(0)O_j(t+\tau) \right] \) [72], with \( L' \rho' = \rho', \) and \( \bar{\rho}' \) is given by

\[ \bar{\rho}' = -i\omega_c \left[ c_L^\dagger c_L + c_R^\dagger c_R, \rho \right] + \kappa \mathcal{L} [c_L] \rho + \kappa \mathcal{L} [c_R] \rho + \kappa |r| \left( e^{i\phi} \left[ c_L \rho, c_R^\dagger \right] + e^{-i\phi} \left[ c_R, c_L^\dagger \rho \right] \right) \]

From the above quantum master equation, we have

\[ \frac{d}{dt} \langle \tilde{p} \rangle = \begin{pmatrix} \omega_c - i\frac{\kappa}{2} & 0 & ge^{-i\phi} \\ -ik|r|e^{i\phi} & \omega_c - i\frac{\kappa}{2} & ge^{i\phi} \\ ge^{i\phi} & ge^{-i\phi} & \omega_0 - i\frac{\kappa}{2} \end{pmatrix} \langle \tilde{p} \rangle \]

where \( \langle \tilde{p} \rangle = \begin{pmatrix} \langle c_L^\dagger (0)c_L(\tau) \rangle, \langle c_L^\dagger (0)c_R(\tau) \rangle \end{pmatrix}^T \). Since the initial state of electromagnetic environment is the vacuum state, we can obtain the analytical expression of \( \langle c_R^\dagger (0)c_L(\tau) \rangle \) with \( \langle c_L^\dagger (0)c_L(0) \rangle = 1 \); other correlation functions can be solved in similar fashion. We finally arrive at

\[ J(\omega) = J_{DP}(\omega) + J_{EP}(\omega) \]
with the DP and EP contributions
\[ J_{\text{DP}}(\omega) = -g^2 \text{Im} [\chi_{\text{DP}}(\omega)] \]
\[ J_{\text{EP}}(\omega) = -g^2 \text{Im} [\chi_{\text{EP}}(\omega)] \]
where \( \chi_{\text{DP}}(\omega) \) is the usual linear Lorentzian response given by \( \langle c_\uparrow(0)c_L(\tau) \rangle \) and \( \langle c_\uparrow(0)c_R(\tau) \rangle \)
\[ \chi_{\text{DP}}(\omega) = \frac{1}{\pi} \frac{2}{(\omega - \omega_c) + i\kappa/2} \]
and \( \chi_{\text{EP}}(\omega) \) is the EP term contributed by \( \langle c_\uparrow(0)c_R(\tau) \rangle \)
\[ \chi_{\text{EP}}(\omega) = \frac{1}{\pi} \frac{-i\kappa |r| e^{i\Delta \phi}}{[(\omega - \omega_c) + i\kappa/2]^2} \]
where \( \Delta \phi = \varphi - 2\phi \) is the phase difference. The factor 2 in the numerator of \( \chi_{\text{DP}}(\omega) \) represents that the QE is coupled to two cavity modes. Eq. (11) clearly indicates the square Lorentzian profile of \( \chi_{\text{EP}}(\omega) \), a signature of second-order EPs. However, it also indicates that any loss of reflection amplitude, i.e., an imperfect mirror with \( |r| < 1 \), will degrade the quantum effects of EPs. In the following study, we take \( |r| = 1 \) unless specially noted. We highlight that our approach can also be applied to a more general case where the coupling between two cavities are bidirectional but asymmetrical, for example, introducing a Rayleigh-type nano-scatter in close proximity of microring [8, 19, 73]. As a result, the system is drawn out of chiral EPs, and such configuration can be utilized to study the emergence of quantum effects of EPs. Furthermore, one can obtain the LDOS of target nanophotonic structure that owns desirable quantum effects of EPs for electromagnetic design. This is important for structure optimization, especially for these consisting of absorbing and dispersing medium like the plasmonic-nanophotonic cavity [74–77], where the coupling rate between cavity modes is hard to evaluate due to their distinct features, contract to the EP cavity studied in this work. On the other hand, the quantum-optics properties of a given nanophotonic structure are predictable, by parametrizing the extended cascaded quantum master equation (Eq. (1)) using the simple curve fitting of LDOS obtained from electromagnetic simulations. Therefore, Eqs. (8)-(11) build the bridge between the electromagnetic design of nanophotonic structures and the quantum state control by EPs.

To validate our theory and as an example, in Fig. 2 we compare the spectral density given by Eqs. (8)-(11) with the electromagnetic simulations using a realistic EP cavity. The QE is Z-oriented (perpendicular to the cavity plane), and located at one-quarter of the perimeter from the waveguide-cavity junction \( X = 0 \) in the clockwise direction. For QE with a dipole moment of \( \mu = 60 \) Debye, we can evaluate \( (\omega_c, \kappa, g) = (0.78122 \text{eV}, 152.8 \mu\text{eV}, 24.9 \mu\text{eV}) \) by curve fitting using Eq. (10) from the corresponding DP cavity, i.e., EP cavity without the mirror in waveguide; subsequently, the spectral density of EP cavity can be analytically obtained from Eqs. (8)-(11). Fig. 2 shows the spectral density of EP cavity, where we can see the good accordance between the analytical expressions and the numerical simulation for \( \Delta \phi = 0 \) and \( \pi \). The results confirm the validity of our LDOS theory for EP cavity. With the obtained parameters, the quantum dynamics of system can be investigated by the extended cascaded quantum master equation.

The exotic quantum dynamics at EPs can be well understood from the spectral properties of QE. The emission spectrum of QE is defined as \( S(\omega) = (2\pi)^{-1} \int_0^\infty dt_1 \int_0^\infty dt_2 e^{i\omega(t_2-t_1)} \langle E^- (r, t_1) \cdot E^+ (r, t_2) \rangle \) [72], with \( E^- (r, t_1) = [E^+ (r, t_1)]^\dagger \) and \( E^+ (r, t) = \)
$e^{-i\omega t} \int_{t_1}^{t_2} dt' G(t, t') \sigma_+(t') e^{i\omega t'}$, where $G(t, t')$ is a kind of propagator dependent on the detection method. For assuming an initially excited QE and $r = r_0$, where $r_0$ is the QE position, $S(\omega)$ becomes the SE spectrum, also called the polarization spectrum [78, 79], which reflects the local dynamics of a QE, and is given by $S(\omega) = \int_{-\infty}^{\infty} dt_2 \int_{-\infty}^{\infty} dt_1 e^{i\omega(t_2-t_1)} \langle \sigma_+ t_1 \sigma_- t_2 \rangle$. Denoting $\tau = t_2 - t_1$ and taking the limit of $t_1 = t \to \infty$, the SE spectrum can be expressed as $S(\omega) = \lim_{t \to \infty} \text{Re} \left[ \int_{-\infty}^{\infty} d\tau \left( \sigma_+(t+\tau) \sigma_-(t) \right) e^{i\omega \tau} \right]$, where $\langle \sigma_+(t+\tau) \sigma_-(t) \rangle$ can be calculated using the quantum regression theorem. Then the SE spectrum can be analytically obtained

$$S(\omega) = \frac{1}{\pi} \frac{\gamma + \Gamma(\omega)}{[\omega - \omega_0 - \Delta(\omega)]^2 + \left(\frac{\gamma + \Gamma(\omega)}{2}\right)^2}$$  \hspace{1cm} (12)

with the photonic Lamb shift

$$\Delta(\omega) = g^2 \text{Re} \left[ \chi_{\text{DP}}(\omega) + \chi_{\text{EP}}(\omega) \right]$$  \hspace{1cm} (13)

and the local coupling strength

$$\Gamma(\omega) = -2g^2 \text{Im} \left[ \chi_{\text{DP}}(\omega) + \chi_{\text{EP}}(\omega) \right]$$  \hspace{1cm} (14)

Note that the SE dynamics of QE can be retrieved from $F[\mathcal{S}(\omega)]$, the Fourier transform of SE spectrum. Eqs. (12)-(14) indicate that LDOS is crucial for understanding the exotic behavior of quantum dynamics at EPs. The contributions of DP and EP are separated in Eqs. (13) and (14), which is beneficial to unravel how the emergence of EPs alters the quantum dynamics.

### III. RESULTS AND DISCUSSION

#### A. Exceptional point induced transparency

We first consider the weak coupling regime, where the Purcell effect is expected to modify the SE rate of a QE. Note that the quality factor of modes in EP cavity is typically $10^4$ due to the dissipative coupling to waveguide, therefore, we will focus on the case of $\kappa \gg \gamma$ in the following discussion. Fig. 3(a) shows the SE dynamics of QEs with different $\gamma$, while the coupling rate $g$ and the cooperativity $C = \gamma g^2/\kappa \gamma = 0.2$ are set to be fixed. For resonant QE-cavity coupling, the maximum Purcell factor of DP cavity is equal to $(C + 1)$, obtained by substituting Eq. (8) into $F_P(\omega_c) = J(\omega_c)/J_0(\omega_c) + 1$. Therefore, QEs experience the Purcell enhancement and the corresponding SE dynamics manifests a faster decay, as shown by the dashed lines in Fig. 3(a). On the contrary, the SE dynamics of QEs in EP cavity is counterintuitive, which shows good accordance with QEs in the free space as if the cavity is absent. Therefore, we call this intriguing phenomenon as **EP induced transparency**. The same effect has been reported in Ref. [42] and explained as a consequence of decoupling between the QE and the cavity modes due to the formation of stand wave pattern of electric field at EPs. Different from the previous work, here we unravel from the perspective of LDOS that the EP induced transparency results from the precise cancelation of EP and DP contributions of $J(\omega)$ at special frequency point, giving rise to null Purcell enhancement for a QE with the same transition frequency.

The frequency corresponding to null Purcell enhancement can be found from Eqs. (8)-(11) by letting $J(\omega) = 0$. The solution takes a simple form

$$\Delta \omega = \Delta \omega_m \equiv -\frac{\kappa}{2} \tan \left( \frac{\Delta \phi}{2} \right)$$  \hspace{1cm} (15)

where $\Delta \omega = \omega - \omega_c$ is the frequency detuning. Fig. 3(b) displays the normalized LDOS of cavity $J(\omega)/J_0(\omega)$ versus $\Delta \omega$ and $\Delta \phi$, which can be easily tuned by varying the mirror location, i.e., the waveguide length $L$. We can see that the zero point of $J(\omega)$ goes away from cavity resonance as $\Delta \phi$ increase from 0 to $\pi$, and the opposite tendency is observed for $\Delta \phi \in [\pi, 2\pi]$, where the zero point of frequency is larger than cavity resonance in this case. Especially, the null Purcell enhancement is achieved at the cavity resonance with $\Delta \phi = 0$, which is just the
circumstance studied in Ref. [42]. Fig. 3(c) plots the decomposition of $J(\omega)$ for $\Delta \phi = 0$ and $\pi/2$, where we plot $-J_{EP}(\omega)$, i.e., we reverse the curve of EP term by taking a negative sign, to clearly display the cancellation of EP and DP contributions, see the shaded circles for indicating $\Delta \omega_{\min}$. It shows that the EP term can be negative and weakens the Purcell effect. For $\Delta \phi = 0$, the EP response features an even symmetry and a narrow linewidth, compared to a DP (Lorentzian) cavity. The upper panel of Fig. 3(c) shows $J(\omega) > 0$ for $|\omega - \omega_c| > \kappa/2$, while $J(\omega) < 0$ for $|\omega - \omega_c| < \kappa/2$. The resultant $J(\omega)$ is slightly enhanced at a wide frequency range far detuned from the cavity, but strongly suppressed around the cavity resonance. The EP response cancels the DP contribution at $\Delta \omega$, resulting in the vanishing Purcell effect for a QE resonantly coupled to EP cavity. The lower panel of Fig. 3(c) shows that the EP response exhibits a totally different profile for $\Delta \phi = \pi/2$, which becomes odd symmetry and changes its sign at cavity resonance. As a result, the disappeared Purcell effect occurs at the left side of cavity resonance ($\Delta \omega = -\kappa/2 < 0$), while the enhanced Purcell effect is observed at the whole frequency range of $\omega > \omega_c$, leading to a strongly asymmetrical line-shape of $J(\omega)$.

Fig. 3(d) plots the Purcell effect enhancement of EPs as the function of $\Delta \phi$ for various $\Delta \omega$ and $|r|$, which is defined as $\eta \equiv J(\omega_c)/J_{DP}(\omega_c) = J_{EP}(\omega_c)/J_{DP}(\omega_c) + 1$. It shows that EP cavity attains greatest tunability to Purcell effect from chiral EPs. Particularly, with resonant QE-cavity coupling, a double increase of Purcell enhancement compared to a DP cavity can be realized with a prefect mirror and $\Delta \phi = \pi$, permitting the stronger light-matter interaction. For a QE detuned from cavity, the null Purcell enhancement still exists with $|r| = 1$, but the maximum $\eta$ decreases. $\eta > 1$ can be achieved inside the parameter region indicated by the gray dashed line in Fig. 3(b). Fig. 3(d) also shows that the maximum $\eta$ drops as the loss of reflection amplitude increases, but a mirror with $|r| > 0.8$ still holds great tunability of Purcell effect, where, for example, the maximum $\eta$ is greater than 1.75 at $\Delta \omega = 0$. It is worth noting that $|r| \sim 0.98$ is achievable [42, 55], and thus a practical mirror will not significantly weaken the ability of chiral EPs to tune the Purcell effect in experiment.

B. Decay suppression of Rabi oscillation at chiral EPs

The above analysis reveals that the chiral EPs have the ability to significantly modify the SE process of a QE, from completely suppressed to enhanced Purcell effect, and thus a EP cavity can provide greater degrees of freedom to control the light-matter interaction than a DP cavity. We now go beyond the weak coupling regime and investigate the effects of chiral EPs on the coherent energy exchange between the QE and the cavity, i.e., the Rabi oscillation.

$$J(\omega) \approx \frac{\kappa}{\omega - \omega_c} - \frac{\kappa}{\omega - \omega_{c2}}$$

where the Rabi oscillation is evident. The parameters are $g = 10\gamma$, $\kappa = 20\gamma$, and $\Delta \phi = \pi$. It also shows that the maximum population of $c_L$ cavity is slightly lower than the DP cavity, due to the lack of direct energy blackflow from $c_R$ cavity. On the other hand, the Rabi oscillation of $c_L$ cavity sustains for a longer time and four cycles

![FIG. 4. Decay suppression of Rabi oscillation in EP cavity.](image)

(a) Population dynamics for QE (blue line) and cavity modes of EP (yellow and pink lines) and DP (thin green line with shading) cavities. SE dynamics of a bare QE (in free space) with the same decay rate is also shown for comparison (dashed red line). The parameters are $g = 10\gamma$, $\kappa = 20\gamma$, $\omega_0 = \omega_c$, and $\Delta \phi = \pi/2$. The inset plots the corresponding normalized LDOS of cavity. (b) is the same as (a) but for $g = 100\gamma$. The upper panels of (c) and (d) plot the normalized SE spectra corresponding to (a) and (b) for EP (blue lines) and DP (red lines) cavities, respectively. While the lower panels show the SE spectrum versus $\Delta \phi$. The inset in the upper panel shows the peak of Rabi doublet with lower energy, where the SE spectrum of a bare QE is also presented for comparison (black dashed line). The white dashed lines in the lower panel track the eigenenergy levels given by Eqs. (16) and (17).
can be observed, indicating less energy dissipation in EP cavity. While the Rabi oscillation of DP cavity manifests a faster decay, and the population of second-cycle oscillation is much smaller than that of EP cavity and hard to recognize. Meanwhile, the maximum population of $c_R$ cavity is more than threefold compared to the DP cavity, and reaches $\sim 0.66$. However, though the higher population, the period of Rabi oscillation is not obviously changed. The SE spectra shown in the upper panel of Fig. 4(c) reveals that, it is because the splitting width changed. The SE spectra shown in the upper panel of Fig. 4(d), where the SE spectrum of a bare QE is shown for comparison. Eq. (16) indicates that the linewidth narrowing in this case is anomalous, since the minimum decay achieved at $\Delta \phi = \pi$ is $-2\text{Im} \left[ \omega_{1,2} \right] \approx \gamma/2$, a half of a bare QE, see the inset in the upper panel of Fig. 4(d), where the SE spectrum of a bare QE is shown for comparison. Eq. (16) also indicates that the reduction of decay rate originates from the unidirectional coupling ($\kappa \cos(\Delta \phi)$) term between two cavity modes for chiral EPs. The condition $\Delta \phi = \pi$ implies that the decay suppression of Rabi oscillation can be understood as a result of the backflow of partial leaky energy through the reflection of mirror, and thus unique to EP cavity. The lower panels of Fig. 4(c) and (d) display the SE spectra as the function of $\Delta \phi$ for $g = 10\gamma$ and $100\gamma$. It is interesting to see that though both are under the strong coupling regime with the equal splitting seen in SE spectrum at $\Delta \phi = \pi$, the underlying evolution of Rabi splitting is distinguishing. The Rabi splitting in the case of $g = 100\gamma$ is nearly unchanged as $\Delta \phi$ varies, while for $g = 10\gamma$, the peak of Rabi splitting with lower energy evolves from cavity resonance, and the corresponding linewidth is much narrower than the peak with higher energy until $\Delta \phi$ approaches to $\pi$.

The above results indicate that EP cavity together with sufficiently strong QE-cavity interaction will have significant advantages in quantum-optics applications, such as the spontaneous entanglement generation (SEG) between qubits, benefiting from the decoherence suppression by chiral EPs. Here we demonstrate the EP enhanced two-qubit SEG in EP cavity, see Fig. 5(a) for schematic. Extending our model into the multi-QE case is straightforward

$$
\dot{\rho}_M = -i \left[ H^M, \rho \right] + \gamma \sum_i \mathcal{L} \left[ \sigma^-_i \right] \rho + \kappa \mathcal{L} \left[ c_L \right] \rho + \kappa \mathcal{L} \left[ c_R \right] \rho + \kappa \rho \mathcal{L} \left[ c_R \right] \rho + \kappa \rho \mathcal{L} \left[ c_L \right] \rho
$$

(18)
with full Hamiltonian $H^M = H_0^M + H_I^M$, where the free Hamiltonian $H_0^M$ and the interaction Hamiltonian $H_I^M$ are given by

$$H_0^M = \omega_0 \sum_i \sigma_i^+ \sigma_i^- + \omega_c c_L^\dagger c_L + \omega_c c_R^\dagger c_R$$

where we assume the identical qubits. With an initially excited qubit, the generated entanglement is quantified by the concurrence $C(t)$ between qubits, which in our system is given by the simple expression $C(t) = 2 |C_{qe}(t)C_{qe}(t)|$ [71, 81], with $C_{qe}(t)$ and $C_{qe}(t)$ being the probability amplitudes of two single-excitation states that one qubit in the excited state while another in the ground state.

In the case of resonant qubit-cavity coupling, the upper bound of maximum concurrence is $0.5$. As shown in Fig. 5(b), the concurrence exponentially decays with a rate of $\gamma$ for a DP cavity; while in the EP cavity, the evolution of $C(t)$ manifests two distinguishing stages with different decay rates: the decay of $C(t)$ is identical as DP cavity for $t < 1.6 \gamma^{-1}$; after that, it is found to decay with a rate $\sim \kappa^3/32g^2$, obtained from the eigenenergies of Eq. (18) in the single-excitation subspace [71]. With a large coupling rate $g = 100\gamma$, $\kappa^3/32g^2 = 1/40\gamma \ll \gamma$, and thus the decay of $C(t)$ for $t > 1.6 \gamma^{-1}$ is much slower than a DP cavity. It also indicates that for a larger $g$, the decay suppression of quantum entanglement will be more evident. Fig. 5(e) displays the $C(t)$ corresponding to the maximum concurrence $C_{\text{max}} = \max [C(t)] = 0.9866$, achieved with the optimal qubit-cavity detuning $\Delta_{\text{opt}} = \omega_0 - \omega_c = 2.32g$. In this case, the concurrence decays much faster in a DP cavity than the EP cavity, exhibiting the great potential of EP cavity to preserve the generated entanglement.

C. Population trapping at chiral EPs with atom-photon bound states

Inspired by the above results that the spectral linewidth can be smaller than a bare QE, in the following we investigate the SE dynamics of an ideal QE at chiral EPs, i.e., $\gamma = 0$. Fig. 6(a)-(c) plot the time evolution of QE and two cavity modes versus $\Delta \phi = 0$ for resonant QE-cavity coupling, where we can observe the population trapping at $\Delta \phi = 0$ and $0.77\pi$, while it cannot be achieved in a DP cavity, see the green shaded lines. We
can also see that for $\Delta \phi = 0$, the steady-state population of QE is almost zero and smaller than that of two cavity modes; however, the situation is reversed for $\Delta \phi = 0.77\pi$, and the trapped population is also higher than the case of $\Delta \phi = 0$. These phenomena indicate the formation of atom-photon bound states in the system, but the underlying physical mechanisms are different. For the former case ($\Delta \phi = 0$), the destructive interference of two cavity modes through the reflection of mirror gives rise to standing wave inside the cavity for population trapping. While the mechanism of bound state for $\Delta \phi \neq 0$ is similar to the accidental bound states in the continuum (BIC) [82–84], also known as Friedrich-Wintgen BIC, which results from the destructive interference of two coupling pathways between QE and $c_R$ cavity, the direct coupling and the indirect coupling mediated by $c_L$ cavity. The explicit expression of non-zero $\Delta \phi$ for BIC can be obtained by finding purely real eigenvalues of characteristic matrix Eq. (5), which is given by

$$\Delta \phi_{BIC} = 2 \arccos \left( \frac{\kappa}{2\sqrt{2g}} \right)$$

(21)

In Fig. 6(d), we plot the imaginary part (decay) of eigenenergies in the single-excitation subspace as the function of $\Delta \phi$, where it shows that bound states for $\Delta \phi = 0$ and $\Delta \phi_{BIC}$ are achieved at different eigenenergy levels, $\omega_3$ and $\omega_2$, respectively. Fig. 6(e) shows the Hopfield coefficients (eigenvector components) of $\omega_2$ versus $\Delta \phi$, where we can see that the coefficients for cavity modes and QE are both equal to 0.5 at $\Delta \phi_{BIC}$; by contrast, the Hopfield coefficient of QE for $\omega_2$ is less than 0.1 at $\Delta \phi = 0$ and much smaller than cavity modes. It explains the different distributions of steady-state population for two bound states shown in Fig. 6(a)-(c).

Since BIC exhibits better energy preservation at chiral EPs, we then consider a general case of arbitrary $\Delta \phi$, and find the QE-cavity detuning corresponding to BIC. From the denominator of SE spectrum (Eq. (12)), we can obtain

$$\Delta \omega_{BIC} = \frac{2g^2}{\kappa} \sin(\Delta \phi) + \Delta \omega_m$$

(22)

We can see that Eq. (22) encompasses the case of EP induced transparency, with $g \ll \kappa$ in the weak coupling regime. As $g$ increases, the contribution of the first term in Eq. (22) becomes significant, the resultant $\Delta \omega_{BIC}$ is no longer overlapped with the location of null Purcell enhancement for $\Delta \phi \neq 0$.

Fig. 7(a) shows the time evolution of cavity population versus QE-cavity detuning $\Delta_{0c} = \omega_0 - \omega_c$, with parameters $\Delta \phi = \pi/2$, $g = 10\gamma$, and $\kappa = 20\gamma$. Eq. (22), as well as the decay of eigenenergies, indicates that BIC is achieved at $\Delta_{0c} = 0$, see the blue line in the right panel of Fig. 7(a). As a result, the decay of cavity population is obviously slower around $\Delta_{0c} = 0$, and the populations of QE and two cavities are partially trapped after a few cycles of Rabi oscillation at $\Delta_{0c} = 0$, as shown by the inset of Fig. 7(a). It shows that though the steady-state populations of two cavities are the same, $c_R$ cavity exhibits stronger Rabi oscillation due to the unidirectional energy transfer from $c_L$ cavity. The maximum population of $c_R$ cavity reaches 0.4, while that of $c_L$ cavity is about 0.2.

Fig. 7(b) displays the steady-state cavity population versus $\Delta \phi$ and $g$. It reveals that for a given $\Delta \phi$, there is an optimal $g$ for maximum population, denoted as $P_c^{opt}(\Delta \phi)$. The maximal $P_c^{opt}(\Delta \phi)$ is found to achieve at $\Delta \phi = 0$, i.e., $P_c^{opt}(0)$ is the upper bound of steady-state cavity population. $P_c^{opt}(\Delta \phi)$ is evaluated to decrease by 0.03 as $\Delta \phi$ varies from 0 to 0.9$\pi$, and thus is robust against the variation of $\Delta \phi$. For $\Delta \phi = 0$, the steady-state populations can be analytically given from Eq. (4), which are $P_c^{ss} = \kappa^2 / (8g^2 + \kappa^2)^2$ for QE and $P_c^{ss} = 8g^2\kappa^2 / (8g^2 + \kappa^2)^2$ for cavity. We can see that $P_c^{ss}$ reduces as $g$ increase, exhibiting distinguishing feature from $P_c^{ss}$. Furthermore, we can obtain the analyt-
FIG. 8. Decay suppression in long-time dynamics with atom-photon bound states at chiral EPs. (a) Logarithmic plot of QE dynamics with detuning $\Delta \omega_{\text{BIC}}$ (solid lines) from cavity for various $\Delta \phi$. The dynamics of bare QEs are shown for comparison (dashed lines). The parameters are $g = 5\gamma_0$ and $g^2/\kappa \gamma = 5/4$, where $\gamma_0$ is the unit rate of QE decay. Note that the SE rate $\gamma$ is taken to be $\gamma_0$, $\gamma_0/2$, and $\gamma_0/4$ for $\Delta \phi = 0$, $\pi/4$, and $\pi/2$, respectively, to distinguish different curves, but the corresponding $\Gamma_m$ is still the curve of $g = 5\gamma$ in (b). (b) $\Gamma_m$ as the function of $\Delta \phi$ for various $g$. (c) $g^{(2)}(0)$ (upper panel) and population (lower panel) of $c_L$ cavity as the function of frequency detuning in the case of $c_R$ cavity drive. The results are obtained by implementing a driving Hamiltonian $H_d = \Omega \left( c_R^\dagger + c_R \right)$ in the extended cascaded quantum master equation (Eq. (1)) and numerically calculating using QuTip [80], with parameters $g = 5\gamma$, $\Delta \phi = 0$, $\Delta \omega_c = \Delta \omega_{\text{BIC}}$ and driving strength $\Omega = 0.2\gamma$. $\kappa = 20\gamma$ for all figures.

As applications, the significantly reduced decay of eigenenergies makes the EP cavity-QE system advantageous for single-photon generation exploiting the photon blockade [85, 86], where the photon antibunching takes place at the frequency of one of eigenenergy levels in the single-excitation subspace with a coherent input. Therefore, single-photon blockade can attain efficiency improvement at chiral EPs. In the case of $c_R$ cavity drive, Fig. 8(c) compares the single-photon purity $g^{(2)}(0) = \langle c_L^\dagger c_L c_L c_L \rangle/n_L^2$ and population $n_L = \langle c_L^\dagger c_L \rangle$ of DP and EP cavities, with parameters $\Delta \phi = 0, g = 5\gamma$ and $\kappa = 20\gamma$. We evaluate $g \approx g_c = \sqrt{(\kappa^2 + 2\gamma^2)/4}$, where $g_c$ is the critical coupling rate of strong coupling for a DP cavity [87]. It implies that the strong antibunching is absent in DP cavity, since the system just reaches the strong coupling regime. As shown in Fig. 8(a), $g^{(2)}(0)$ curve of DP cavity is flat, and the minimum $g^{(2)}(0)$ is $\sim 0.1$. With the resonant QE-cavity coupling, which is also the optimal configuration detuning according to Eq. (22), the $c_R$ cavity demonstrates a great enhancement of single-photon purity by over an order of magnitude at $\Delta \phi = 0$, with the minimal $g^{(2)}(0) \sim 0.005$, accompanied by a hundredfold improvement of population. While for $\Delta \phi = \pi/4$, the single-photon purity and population are both improved by about an order of magnitude at chiral EPs. Therefore, EP cavity shows great potential in building high-efficiency single-photon source.

IV. CONCLUSION

In this work, we study the spontaneous emission of a quantum emitter located in microcavity supporting chiral EPs, and find the intriguing phenomena of EP in-
duced transparency and the anomalously small decay rate of transient and long-time dynamics. An analytical description of LDOS and SE spectrum is presented to unveil that the non-Lorentzian response and atom-photon bound states are responsible for these peculiar quantum dynamics at chiral EPs. The results demonstrate the striking ability of chiral EPs to suppress the decoherence process occurred in both weak- and strong-coupling regimes. Therefore, we envision that EPs can offer the robustness against the dephasing for various quantum-optics applications, including but not limited to high-efficiency single-photon sources [88, 89], nonlinear interaction at the single-photon level [90], and high-fidelity entanglement generation and transport [91, 92]. Besides the nanophotonic structures, the quantum effects of chiral EPs demonstrated in this work can also be implemented in other kind of platforms, such as superconducting [93, 94], cavity optomechanics [95, 96], and open cavity magnonic systems [28, 97]. We believe that our work can provide insights on the effects of EPs in diverse quantum systems and is instructive for harnessing the non-Hermiticity for building novel quantum devices.

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