Exceptional cavity quantum electrodynamics

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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. Based on the recently proposed microcavity with exceptional surface (ES), we report and explain the peculiar quantum dynamics in atom-photon interaction associated with EPs: cavity transparency, decoherence suppression beyond the limitation of Jaynes-Cummings (JC) system, and the population trapping of lossy cavity. An analytical description of the local density of states (LDOS) for ES microcavity is derived from an equivalent cavity quantum electrodynamics (QED) model, which goes beyond the single-excitation approximation and allows exploring the quantum effects of EPs on multiphoton process by parametrizing the extended cascaded quantum master equation. It reveals that a square Lorentzian term in LDOS induced by second-order EPs interferes with the linear Lorentzian profile, giving rise to cavity transparency for atom with special transition frequency in the weak coupling regime. This additional contribution from EPs also breaks the limit on dissipation rate of JC system bounded by bare components, resulting in the decoherence suppression with anomalously small decay rate of the long-time dynamics and the Rabi oscillation. Remarkably, we find that the cavity population can be partially trapped at EPs, achieved by forming a bound dressed state in the limiting case of vanishing atom decay. Our work unveils the exotic phenomena unique to EPs in cavity QED systems, 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 [15–18], and unidirectional invisibility [19, 20]. 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, 21–25], acoustics [26] to macroscopic facilities, such as fiber network [27], electric circuits [28] and heat diffusive system [29]. In recent years, great efforts have been dedicated to accessing the quantum EPs by implementing the non-Hermiticity in quantum systems [11, 30–41], and investigating the quantum states control through EPs [18, 33–35, 40, 42–44]. In this respect, pioneer works have shown the ability of EPs to tune the photon statistics [34, 35], enhance the sensitivity of quantum sensing [45, 46], and steer the evolution of single quantum system [18, 33, 40, 42–44, 47–49].

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 (EM) environments supporting EPs attracts growing attention [40, 42, 44, 47], since the modification of spontaneous emission (SE) features a square Lorentzian profile around EPs, contrast to the nondegenerate resonances in conventional cavity with Lorentzian lineshape, and can lead to a greater enhancement of SE rate [42, 44, 47]. It implies that the formation of EPs in nanophotonic structures significantly alters the local density of states (LDOS), which fully describes the interaction between a QE and arbitrary EM environment [50–53]. Tailoring the LDOS of EM environment is crucial for optimizing the performance of many practical applications, ranging from traditional optoelectronic devices like lasers [54, 55] and solar cells [56], to advanced quantum technologies, such as quantum light sources [57–59] and logical gate [60, 61]. Therefore, a LDOS theory that can tackle the effects of EPs is of both fundamental and applied significance. The formalism of LDOS based on the classical Green function separates the EPs contribution from the usual Lorentzian term [42, 43, 62], while it leaves the origin of EPs unclear, i.e., information of the underlying cavity resonances is ambiguous, the parame-
Bare microcavity Microcavity with exceptional surface

L distance between the original and the mirrored cavities is polarized QE is replaced by a circularly polarized one. The other CCW mode through mirror symmetry, and the linearly with a perfect mirror, where the CW mode is flipped to an-
tion of the equivalent cavity QED model for the system in (b)
the CW mode and the CCW mode. (c) Conceptual illustra-
ate an exceptional surface due to the chiral coupling between
semi-infinite waveguide with a mirror at the one side can cre-
generate CW and CCW modes with equal coupling rate g. (b) A microring cavity coupled to a semi-infinite waveguide with a mirror at the one side can create an exceptional surface due to the chiral coupling between the CW mode and the CCW mode. (c) Conceptual illustration of the equivalent cavity QED model for the system in (b) with a perfect mirror, where the CW mode is flipped to another CCW mode through mirror symmetry, and the linearly polarized QE is replaced by a circularly polarized one. The distance between the original and the mirrored cavities is L.

Fig. 1. Schematic diagrams of cavity QED model. (a) A whispering-gallery-mode (WGM) cavity represented by a microring resonator supports the degenerate clockwise (CW) and counterclockwise (CCW) modes, and is operated at a diabolic point (DP). A linearly polarized quantum emitter (QE) is simultaneously coupled to both CW and CCW modes with equal coupling rate g. (b) A microring cavity coupled to a semi-infinite waveguide with a mirror at the one side can create an exceptional surface due to the chiral coupling between the CW mode and the CCW mode. (c) Conceptual illustration of the equivalent cavity QED model for the system in (b) with a perfect mirror, where the CW mode is flipped to another CCW mode through mirror symmetry, and the linearly polarized QE is replaced by a circularly polarized one. The distance between the original and the mirrored cavities is L.

II. MODEL AND THEORY

Different strategies have been proposed to construct an ES in whispering-gallery-mode (WGM) cavity, by introducing the additional scattering structures [8, 65] or coupling to a waveguide terminated by a mirror at the one side [11, 44]. The cavity QED model for the former can be developed by treating the scattering structures as Rayleigh-type nanoparticles [65, 66], while the latter should be considered as a cascaded system [67, 68] instead in the framework of waveguide QED [69–71] due to the existence of mirror. The formation of ES, a collection of EPs in a high-dimensional parameter space rather a single EP, results from the unidirectional coupling between the clockwise (CW) and the counterclockwise (CCW) modes assisted by the auxiliary reservoir modes. Therefore, the underlying physical mechanism is similar for two implementations. 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 type of ES microcavity.

Fig. 1(a) depicts a basic cavity QED system consisting of a WGM microcavity and a quantum emitter (QE). The QE is linearly polarized, and coupled to a pair of degenerate CW and CCW modes with equal coupling rate g. Since there is no interaction between two cavity modes, the system is operated at a diabolic point (DP) and described by the Jaynes-Cummings (JC) model. 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 photonics structure exhibits a surface of second-order EPs in the parameter space, and offers great tunability to SE rate of a QE. In order to develop an intuitive cavity QED model for this novel ES microcavity, we transform the system to an equivalent one through mirror symmetry, as shown in Fig. 1(c). The original CW 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 cavities in the equivalent model constitute a cascaded system [67, 68, 72], the quantum dynamics is described by the extended cascaded QME \( \hat{h} = 1 \)

\[
\dot{\rho} = -i \{ H_0 + H_1, \rho \} + \gamma L [ \sigma_+ \rho + \kappa \mathcal{L} [ c_L ] \rho + \kappa \mathcal{L} [ c_R ] \rho \}
+ \kappa |\rangle \langle \sigma_+ | \left( e^{i \varphi} [ c_L \rho e^{i \varphi} c_R ] + e^{-i \varphi} [ c_R \rho e^{-i \varphi} c_L ] \right)
\]

(1)

where \( \mathcal{L}[\rho] = \mathcal{O} \rho \mathcal{O}^\dagger - \{ \mathcal{O} \rho, \mathcal{O} \} / 2 \) is the Liouvillian superoperator. The second line of Eq. (1) describes the unidirectional coupling from CCW mode to CW mode, where \( |\rangle \langle \sigma_+ | \) is the reflection of mirror, and \( \varphi = 2\beta L \) is the phase factor, with \( \beta \) and \( L \) being the propagation constant of waveguide and the distance between two cavities, respectively. The free Hamiltonian \( H_0 \) and the interaction Hamiltonian \( H_I \) are given by

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

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

(2)

where \( \sigma_- \) is the lowering operator of QW with transition frequency \( \omega_c \) and SE rate \( \gamma \) in homogeneous medium. \( c_L \) and \( c_R \) are the bosonic annihilation operators for CCW modes of the original and the mirrored microcavities with resonance frequency \( \omega_r \), respectively. \( \kappa \) 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. The phase factor \( \varphi \) originates from the traveling wave nature of WGM modes and depends on the azimuthal orientation of QE location [73, 74]. From Eq. (1), we take the operator expectation values to obtain the equations of motion

\[
\frac{d}{dt} \vec{p} = -i \hat{M} \vec{p}
\]

(3)

where \( \vec{p} = (\langle c_L \rangle, \langle c_R \rangle, \langle \sigma_- \rangle)^T \) and the \( 3 \times 3 \) matrix takes the form

\[
\hat{M} = \begin{pmatrix}
\omega_c - i \frac{\gamma}{2} & 0 & ge^{-i \varphi} \\
-ik |\rangle \langle \sigma_- | & \omega_c - i \frac{\gamma}{2} & ge^{i \varphi} \\
ge^{i \varphi} & ge^{-i \varphi} & \omega_0 - i \frac{\gamma}{2}
\end{pmatrix}
\]

(4)

The normalized LDOS, i.e., Purcell factor, is expressed as

\[
P(\omega) = 2\pi J(\omega) / \gamma + 1,
\]

where

\[
J(\omega) = \int_{-\infty}^{\infty} dv \omega e^{i \omega v} \left( \langle c_L^\dagger (0) + e^{-i 2\varphi} c_R^\dagger (0) \rangle \langle c_L (\tau) + e^{i 2\varphi} c_R (\tau) \rangle \right)
\]

is the spectral density of EM microcavity. The correlation functions can be calculated by quantum regression theorem (OR) \( \langle O_M O_N (\tau) \rangle = \sum_k M_{jk} O_M (\tau) \langle O_N O_j \rangle \) [75] with the following QME

\[
\dot{\rho} = -i \{ H_0, \rho \} + \kappa \mathcal{L} [ c_L ] + \kappa \mathcal{L} [ c_R ]
+ \kappa |\rangle \langle \sigma_+ | \left( e^{i \varphi} [ c_L \rho e^{i \varphi} c_R ] + e^{-i \varphi} [ c_R \rho e^{-i \varphi} c_L ] \right)
\]

(5)

since the initial state of EM environment is the vacuum state. We arrive at

\[
J(\omega) = -g^2 \text{Im} [2 \chi_{DP}(\omega) + \chi_{EP}(\omega)]
\]

(6)

with the usual linear Lorentzian response, i.e., DP term

\[
\chi_{DP}(\omega) = \frac{1}{(\omega - \omega_0) + i \kappa / 2}
\]

(7)

and the additional EP contribution

\[
\chi_{EP}(\omega) = \frac{-i \kappa |\rangle \langle \sigma_+ | e^{i \Delta \phi}}{[\omega - \omega_e + i \kappa / 2]^2}
\]

(8)

where \( \Delta = \varphi - 2 \phi \) is the phase difference. The factor 2 of \( \chi_{DP}(\omega) \) represents that the QE is coupled to two cavity modes. Eq. (8) clearly indicates the squared Lorentzian profile of \( \chi_{EP} (\omega) \), a signature of second-order EPs. However, it also indicates that any loss of reflection amplitude, i.e., an imperfect mirror with \( |\rangle \langle | < 1 \), will degrade the quantum effects of EPs. In the following study, we take \( |\rangle \langle | = 1 \) 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 scatter in close proximity of microring [8, 20, 66]. As a result, the system is drawn out of ES, 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 EM design. This is important for structure optimization, especially for these consisting of absorbing and dispersing medium like the plasmonic-photoncavity [76–79], where the coupling rate between cavity modes is hard to evaluate, contract to the ES microcavity studied in this work. On the other hand, the quantum-optics properties of a given nanophotonic structure are predictable, by parametrizing the extended cascaded QME (Eq. (1)) using the simple curve fitting of LDOS obtained from EM simulations. Therefore, Eqs. (7) and (8) build the bridge between the EM design of nanophotonic structures and the quantum state control by EPs.

The exotic quantum dynamics at EPs can be understood from the spectral properties of QE. The emission spectrum of QE is defined as

\[
S(\omega, \omega_0) = (2\pi)^{-1} \int_0^\infty dt \int_0^\infty dt_1 e^{-i (\omega t - \omega_0 t_1)} \langle E^-(r, t_1) | \langle E^+(r, t_1) | \langle E^+(r, t_1) \rangle \langle E^- (r, t_1) \rangle \rangle
\]

[75], with \( E^- (r, t_1) = [E^+(r, t_1)]^\dagger \) and \( E^+(r, t) = e^{-i \omega_0 t} \int_0^t dt' G (t, t') \langle \sigma_- (t') \rangle e^{i \omega_0 t'} \), where \( G (t, t') \) is a kind of propagator. 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 dipole spectrum [80, 81], which reflects the local dynamics of a QE. The SE spectrum can be analytically expressed as

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

(9)

with the photonic Lamb shift

\[
\Delta (\omega) = g^2 \text{Re} [2 \chi_{DP}(\omega) + \chi_{EP}(\omega)]
\]

(10)

and the local coupling strength

\[
\Gamma(\omega) = -2g^2 \text{Im} [2 \chi_{DP}(\omega) + \chi_{EP}(\omega)]
\]

(11)
Note that the SE dynamics of QE can be retrieved from \( \mathcal{F}[S(\omega)] \), the Fourier transform of SE spectrum. Eqs. (9)-(11) 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. (10) and (11), which is beneficial to unravel how the emergence of EPs alters the quantum dynamics.

III. RESULTS AND DISCUSSION

A. Cavity 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 ES microcavity is typically \( 10^4 \) due to the dissipative coupling to external structures, therefore, we will focus on the case for \( \kappa \gg \gamma \) in the following discussion. Fig. 2(a) shows the SE dynamics of QEs with different \( \gamma \), while the coupling rate \( g \) and the cooperativity \( C = 8g^2/\kappa\gamma = 0.2 \) are set to be fixed. With resonant QE-cavity coupling, the Purcell factor is equal to \( (C + 1) \), and thus QEs in DP cavity experience the Purcell enhancement and the corresponding SE dynamics manifests a faster decay, as shown by the dashed lines in Fig. 2(a). On the contrary, the SE dynamics of QEs in ES microcavity 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 cavity transparency. The same effect has been reported in Ref. [44] and explained as a consequence of decoupling between the QE and the cavity modes due to the special electric field pattern forming at EPs. Here, we unravel from the perspective of LDOS that the cavity transparency in ES microcavity results from the precise cancellation of EP and DP contributions of \( J(\omega) \) at special frequency point, giving arise to null Purcell enhancement for a QE with the same transition frequency.

The frequency of null Purcell enhancement can be found from Eqs. (6)-(8) by letting \( J(\omega) = 0 \), the solution takes a simple form

\[
\Delta \omega = \Delta \omega_m \equiv -\frac{\kappa}{2} \tan \left( \frac{\Delta \phi}{2} \right)
\]

where \( \Delta \omega = \omega - \omega_c \) is the frequency detuning. Fig. 2(b) displays the normalized \( J(\omega) \) 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 \( \Delta \phi \in [\pi, 2\pi] \), while the frequency of zero point 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. [44]. Fig. 2(c) plots the decomposition of \( J(\omega) \) for \( \Delta \phi = 0 \) and \( \pi/2 \), where we reverse the EP term by taking a negative sign to display the cancellation of EP and DP contributions. It shows that the EP term can be negative and weaken the Purcell effect. For \( \Delta \phi = 0 \), the EP response features an even symmetry and a narrow profile. 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 cavity resonance, and results in the vanishing Purcell effect for a QE resonantly coupled to the ES microcavity. The lower panel of Fig. 2(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 disappear Purcell effect occurs at the left side of cavity resonance \( (\Delta \omega_m < 0) \), while the enhanced Purcell effect is observed at the whole frequency range of \( \omega > \omega_c \), leading to a strongly asymmetrical lineshape of \( J(\omega) \).

Fig. 2(d) plots the Purcell effect enhancement of EPs as the function of \( \Delta \phi \) for various \( \Delta \omega \) and \( |r| \), which is defined as \( \eta = (P(\omega) - 1)/C \). It shows that EPs attain...
greatest tunability to Purcell effect at resonant QE-cavity coupling, where a double increase of Purcell enhancement compared to a DP cavity can be realized with a perfect mirror and $\Delta \phi = \pi$, and thus permit the stronger light-matter interaction. Detuned from the cavity resonance, 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. 2(b). Fig. 2(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. It is worth noting that $|r| \sim 0.97$ is achievable [44, 58], and thus a practical mirror will not significantly weaken the ability of EPs to tune the Purcell effect in experiment.

B. Anomalous decay and population trapping

The above analysis reveals that EPs have the ability to significantly modify the SE process, from completely suppressed to enhanced Purcell effect, and thus the ES microcavity 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 study the long-time dynamics of QE decaying in ES microcavity, where the eigenenergy with minimum decay ($\omega_m$) dominates the system evolution. In a DP cavity, the decoherence always takes place with a rate greater than $\gamma$ for $\kappa \gg \gamma$. For the ES microcavity, the minimum decoherence rate can be obtained by finding an optimal QE-cavity detuning ($\Delta \omega_{\text{opt}}$) for achieving a smallest decay of $\omega_m$, denoted as $\Gamma = -\text{Im} [\omega_m]$. From the denominator of SE spectrum (Eq. (9)), we obtain

$$\Delta \omega_{\text{opt}} = \frac{2g^2}{\kappa} \sin(\Delta \phi) + \Delta \omega_m$$  (13)

where the limitation $\gamma \rightarrow 0$ has been taken. We can see that Eq. (13) encompasses the case of cavity transparency, with $g \ll \kappa$ in the weak coupling regime. As $g$ increases, the contribution of the first term in Eq. (13) becomes significant, the resultant $\Delta \omega_{\text{opt}}$ is no longer overlapped with the location of null Purcell enhancement for $\Delta \phi \neq 0$. As shown by Fig. 3(a), the greatest deviation occurs around $\Delta \phi = \pi/2$. Fig. 3(b) plots the minimum decoherence rate $\Gamma$ corresponding to $\Delta \omega_{\text{opt}}$ for various $g$, where we can see that the decoherence rate reaches the minimum at $\Delta \phi = 0$, and increases as $\Delta \phi$ varies from 0 to $\pi$. Remarkably, the decoherence rate $2\Gamma$ in ES microcavity can be smaller than $\gamma$, which is below the limit of JC system for $\kappa \gg \gamma$, and thus is not possible for a DP cavity. We can also see that enhancing the QE-cavity interaction is beneficial to reduce the decoherence rate, and a tenfold reduction ($\Gamma \rightarrow \gamma/20$) can be achieved with a moderate QE-cavity coupling $g = 20\gamma$ at $\Delta \phi = 0$. While for any value of $g$, $\Gamma \rightarrow \gamma/2$ as $\Delta \phi$ approaches to $\pi$, it implies that the minimum decoherence rate at $\Delta \phi = \pi$ is achieved by decoupling the QE from the cavity.

Fig. 3(c) compares the long-time dynamics of initial excited QEs in the free space and in ES microcavity with optimal detuning $\Delta \omega_{\text{opt}}$ from the cavity. We can see from the slope of decay curves that the EPs can offer different degrees of suppression on the decoherence process, according to the value of $\Delta \phi$. As a result, the population of QEs in ES microcavity can be larger than the bare QEs for $t > 50\gamma^{-1}$, as well as the QEs in ES microcavity but with detuning $\Delta \omega_m$ for $\Delta \phi \neq 0$. The corresponding $\Gamma$ of QEs with detuning $\Delta \omega_m$ is plotted as the dashed black line in Fig. 3(b). It shows that though smaller than $\gamma/2$, $\Gamma$ of QEs with detuning $\Delta \omega_m$ evidently increases at both $\Delta \phi = \pi/4$ and $\pi/2$ compared to that of QEs with optimal detuning $\Delta \omega_{\text{opt}}$ (curve of $g = 5\gamma$). Therefore, the resultant population is even lower than a bare QE for $t > 100\gamma^{-1}$, due to the fast depletion by cavity at the initial stage of decay.

The significantly reduced decay of eigenenergy makes
FIG. 4. Population trapping of ES microcavity with vanishing atom decay. (a) Time dynamics of cavity population as the function of QE-cavity detuning $\Delta_{0c}$, with parameters $\Delta \phi = \pi/2$ and $g = 10\gamma$. The inset shows the time dynamics at $\Delta_{0c} = 0$. The right panel plots $\Gamma$ as the function of $\Delta_{0c}$.

(b) Steady-state cavity population versus $g$ and $\Delta_{0c}$. (c) Steady-state populations of QE (solid lines), cavity (thin solid lines with shading), and the system (QE + cavity, dashed lines) as the function of $g$ for various $\Delta \phi$. In all figures, $\kappa = 20\gamma$.

the ES microcavity-QE system advantageous for single-photon generation exploiting the photon blockade [83, 84], where the photon antibunching takes place at the energy of eigenenergy levels in the single-excitation subspace with a coherent input. Therefore, single-photon blockade can attain efficiency improvement by utilizing the ES microcavity. In the CW mode (cR) driven case, Fig. 3(d) compares the single-photon purity $g^{(2)}(0) = \langle c_X^\dagger c_X^\dagger c_X c_X \rangle / n_X$ and population $n_X = \langle c_X^\dagger c_X \rangle$ of DP and ES microcavities, with $X = R, L$, and parameters $\Delta \phi = 0$, $g = 5\gamma$ and $\kappa = 20\gamma$. We evaluate $g \equiv g_c = \sqrt{(\kappa^2 + \gamma^2)/4}$, where $g_c$ is the critical coupling rate of strong coupling for a DP cavity [85]. It implies that the strong antibunching is absent in DP cavity, since the system just reaches the strong coupling regime. As shown in Fig. 3(d), $g^{(2)}(0)$ 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. (13), the CCW mode of ES microcavity (cL) 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 in ES microcavity. Therefore, ES microcavity shows great potential in building high-efficiency single-photon source.

Inspired by the results of Fig. 3(c) that the long-time decay rate of a QE can be smaller than its SE rate in the free space, we then consider a limiting case of $\gamma \rightarrow 0$. The eigenenergy corresponding to $\Delta \omega_{\text{opt}}$ is found as $\omega_m = \Delta \omega_{\text{opt}}$, which is purely real in this case, signifying the formation of an atom-photon bound state at EPs, except for $\Delta \phi = \pi$, where the QE is decoupled from the cavity and remains in the excited state. The physical mechanism of atom-photon bound state at EPs is similar to the accidental bound states in the continuum (BIC) [86–88], also known as Friedrich-Wintgen BIC, which results from the destructive interference of two coupling pathways between cavities, the direct coupling and the indirect coupling mediated by dissipative channel. While in ES microcavity, the difference is that the direct coupling is replaced by the QE-mediated coupling, and the waveguide-mediated dissipative coupling is unidirectional.

Fig. 4(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. (13), as well as $\Gamma$, indicates that the bound state achieved at $\Delta_{0c} = 0$, see the blue line in the right panel of Fig. 4(a). As a result, the decay of cavity population is obviously slower around $\Delta_{0c} = 0$, and the populations of QE and two CCW modes are partially trapped after a few cycles of Rabi oscillation at $\Delta_{0c} = 0$, as shown by the inset of Fig. 4(a). It shows that though the steady-state populations of two cavities are the same, cR cavity exhibits stronger Rabi oscillation due to the unidirectional energy transfer from cL cavity. The maximum population of cR cavity reaches 0.4, while that of cL, cavity is about 0.2. By contrast, the population trapping cannot be realized in a DP cavity without the bound state, as shown by the green line with shading.

Fig. 4(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 $P^\text{opt}(\Delta \phi)$, which is maximal at $\Delta \phi = 0$, i.e., $P^\text{opt}(0)$ is the upper bound of steady-state cavity population. $P^\text{opt}(\Delta \phi)$ 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. (3), which are $P^\text{es} = \kappa^4 / (8g^2 + \kappa^2)^2$ for QE and $P^\text{es} = 8g^2\kappa^2 / (8g^2 + \kappa^2)^2$ for cavity. We can see that $P^\text{es}$ reduces as $g$ increase, exhibiting distinguishing feature from $P^\text{opt}$. Furthermore, we can obtain $P^\text{opt}(0) = 0.25$, and the expression of leaky
energy $P_{\text{ss}} = 8g^2/(8g^2 + \kappa^2)$, which is monotonically increasing with respect to $g$. It implies that the optimal $g$ for maximum cavity population is the result of balance between the population transfer from QE and the system dissipation. Fig. 4(c) plots the steady-state populations versus $g$ for various $\Delta \phi$. It shows that for $\Delta \phi > 0.9\pi$, trapping with high population can be achieved for cavity at a wide range of $g$.

C. Decoherence suppression in transient dynamics

The decoherence suppression and population trapping at EPs require optimizing the decay of one of eigenenergies in the system, and thus neglect the effects of EPs on the coherent QE-cavity interaction in the transient dynamics. In the following, we investigate the impact of EPs on the coherent energy exchange between the QE and the cavity, i.e., the Rabi oscillation.

Fig. 5(a) shows the time evolution of populations with initially excited QE resonantly coupled to ES microcavity, where the Rabi oscillation is evident. The parameters are $g = 10\gamma$, $\kappa = 20\gamma$, and $\Delta \phi = \pi$. It shows that the maximum population of $c_L$ cavity is slightly lower than the DP cavity, but it sustains for a longer time and four cycles of Rabi oscillation can be observed. While the Rabi oscillation of DP cavity manifests a faster decay, and the population of second-cycle oscillation is much smaller than that of ES microcavity. Meanwhile, the maximum population of $c_R$ cavity is more than threefold compared to the DP cavity, and reaches $\sim 0.66$. However, the period of Rabi oscillation is not obviously changed. The SE spectra shown in the upper panel of Fig. 5(c) reveals that, it is because the splitting width is slightly enlarged, but the linewidth of Rabi peaks are greatly reduced. We highlight that the reduction of linewidth and the resultant slow decay of Rabi oscillation are attributed to the modification of LDOS by the EP term, see the inset of Fig. 5(a). In this sense, the decoherence suppression in ES microcavity is natural: EPs create a new kind of cavity with simultaneously narrower linewidth and greater Purcell enhancement than the usual DP cavity with Lorentzian response, and thus the ES microcavity can exhibit the superior property.

Fig. 5(b) shows the time evolution of populations for $g = 100\gamma$, where the decay of Rabi oscillation is slower than not only the DP cavity but also a bare QE. This is counterintuitive, since the coupling to a lossy cavity should increase the dissipation, and thus result in a faster decay. To gain insight into the problem, we obtain the approximate expressions of eigenenergies

$$\omega_{1,2} \approx \pm \sqrt{2}g + \frac{\kappa}{4} \sin(\Delta \phi) - i \left[ \frac{\cos(\Delta \phi) + 1}{4} \kappa + \gamma \right]$$  \(14\)

\[ \omega_3 \approx \frac{\kappa}{2} \sin(\Delta \phi) \left[ \frac{\cos(\Delta \phi)}{C} - 1 \right] - i \frac{\kappa}{2} \left[ \cos(\Delta \phi) - 1 + \frac{\cos(2\Delta \phi)}{C} \right] \]

where we expand the eigenenergies up to second order with respect to $\kappa$ and $\gamma$. The white dashed lines in the lower panel of Fig. 5(d) track the energies of eigenen-
Concurrence
dynamical concurrence between two qubits for anomalous, since the minimum decay achieved at \( \Delta \) is significant narrowing of Rabi peaks. However, Eq. \( \kappa \) case of \( g \) shown in the upper panel of Fig. 5(d). Same as the other kind of platforms, such as superconducting \[95, 96\], demonstrated in this work can also be implemented in other kind of quantum devices, including but not limited to high-efficiency single-photon sources \[90, 91\], nonlinear interaction at the single-photon level \[92\], and high-fidelity entanglement generation and transport \[93, 94\]. Besides the nanophotonic structures, the quantum effects of EPs demonstrated in this work can also be implemented in other kind of platforms, such as superconducting \[95, 96\], cavity optomechanics \[97, 98\], and open cavity magnonic systems \[30, 99\]. We believe that our work can provide insight into the effects of EPs in a wide range of quantum systems and harnessing the non-Hermiticity for building novel quantum devices.

IV. CONCLUSION

In this work, we unravel some intriguing phenomena related to EPs in the cavity QED system based on a kind of WGM microcavity supporting an exceptional surface, and demonstrate the associated quantum-optics applications. An analytical description of LDOS and SE spectrum is presented to study the quantum effects of EPs, and understand the peculiar features of quantum dynamics at EPs. The results demonstrate the striking ability of EPs to suppress the decoherence occurred in the SE dynamics and the coherent light-matter interaction in the strong coupling regime. We thus envision that the EPs can offer the robustness again the dephasing in the SE dynamics and the coherent light-matter interaction in the strong coupling regime. We thus envision that the EPs can offer the robustness again the dephasing occurred in the SE dynamics and the coherent light-matter interaction in the strong coupling regime. We thus envision that the EPs can offer the robustness again the dephasing occurred in the SE dynamics and the coherent light-matter interaction in the strong coupling regime. We thus envision that the EPs can offer the robustness again the dephasing occurred in the SE dynamics and the coherent light-matter interaction in the strong coupling regime.
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