Radiative decay of the $\Xi_c(2923)$ in a hadronic molecule picture

Feng Yang
School of Physical Science and Technology, Southwest Jiaotong University, Chengdu 610031, China

Hong Qiang Zhu
College of Physics and Electronic Engineering, Chongqing Normal University, Chongqing 401331, China

Yin Huang
Asia Pacific Center for Theoretical Physics, Pohang University of Science and Technology, Pohang 37673, Gyeongsangbuk-do, South Korea and School of Physical Science and Technology, Southwest Jiaotong University, Chengdu 610031, China

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In the present work, we study the radiative decay of newly observed $\Xi_c(2923)^0$ based on the successful explanation that $\Xi_c(2923)^0$ is an $S$-wave $D\Lambda - D\Sigma$ molecular state in our previous study [8]. The radiative decay width of $D\Lambda - D\Sigma$ molecular state into $\Xi_c^0\gamma$ final state through hadronic loops are evaluated using effective Lagrangians. We find that decay widths $\Xi_c(2923)^0 \rightarrow \Xi_c^0\gamma$ and $\Xi_c(2923)^0 \rightarrow \Xi_c^0\gamma$ is evaluated to be approximately 1.23-11.66 KeV and 0.30-3.71 KeV, respectively. These are different from the results [3, 14] that obtained by assuming $\Xi_c(2923)^0$ may be conventional charmed baryon. If measurements are in future experimental, these differences will be very useful to help us to test various interpretations of $\Xi_c(2923)^0$.

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I. INTRODUCTION

At present, there are thirty baryons with only one charm quark listed in the review of PDG [1]. Understanding their internal structure is one of the most meaningful topics in particle and nuclear physics. Among them, three newly observed neutral resonance $\Xi_c^{*0}$ named $\Xi_c(2923)^0$, $\Xi_c(2939)^0$, and $\Xi_c(2965)^0$ in the $K^-\Lambda^+_c$ mass spectra by the LHCb Collaboration [2] aroused widespread discussion. Their masses and widths are

\begin{align*}
\Xi_c(2923)^0 : & \quad M = 2923.04 \pm 0.25 \pm 0.20 \pm 0.14 \text{ MeV} \\
& \quad \Gamma = 7.1 \pm 0.8 \pm 1.8 \text{ MeV}, \\
\Xi_c(2938)^0 : & \quad M = 2938.55 \pm 0.21 \pm 0.17 \pm 0.14 \text{ MeV} \\
& \quad \Gamma = 10.2 \pm 0.8 \pm 1.1 \text{ MeV}, \\
\Xi_c(2964)^0 : & \quad M = 2964.88 \pm 0.26 \pm 0.14 \pm 0.14 \text{ MeV} \\
& \quad \Gamma = 14.1 \pm 0.9 \pm 1.3 \text{ MeV},
\end{align*}

respectively.

Due to the uncertainty of spin-parity, many disputes about their internal structure have been proposed. In the chiral quark model, based on the analysis of the two-body Okubo-Zweig-Iizuka (OZI) allowed strong decays, the states $\Xi_c(2923)^0$, $\Xi_c(2939)^0$, and $\Xi_c(2965)^0$ can be considered as $1P$ $\Xi_c$ state with spin-parity $J^P = 3/2^-$ or $J^P = 5/2^-$ [3]. In this paper the authors suggested to search for them in the $\Xi_c(2923)^0/\Xi_c(2939)^0 \rightarrow \Xi_c\pi$ reaction and the $\Xi_c(2965)^0 \rightarrow \Lambda_c K/\Xi_c\pi$ reaction that can well test the $dsc$ nature of the $\Xi_c$. By employing the $^3P_0$ approach in Ref. [4], the two-body strong decays of the $\Xi_c(2923)^0$, $\Xi_c(2939)^0$, and $\Xi_c(2965)^0$ is calculated. The results indicate that the $\Xi_c(2923)^0$ and $\Xi_c(2939)^0$ can also be $1P$ $\Xi_c$ state, while the $\Xi_c(2965)^0$ is suggested to be $2S$ $\Xi_c$ state. The QCD sum rule suggests that the states $\Xi_c(2923)^0$, $\Xi_c(2939)^0$, and $\Xi_c(2965)^0$ are most likely to be considered as the $P$-wave $\Xi_c$ baryons with the spin-parity $J^P = 1/2^-$ or $J^P = 3/2^-$. These are different from the results by Akaev et al. [6] that the $\Xi_c(2923)^0$ and $\Xi_c(2939)^0$ can be considered as $1P$ excitations of the spin-1/2 flavor-sextet and spin-3/2 baryons, respectively, while the $\Xi_c(2965)^0$ may be the exciting $2S$ state of either spin-1/2 flavor-sextet or antitriplet baryon. In addition, the resonance states of the five-quark configuration are possible candidates of these new states with negative parity [7]. However, a completely different conclusion was drawn from Ref. [8] that the $\Xi_c(2923)^0$ can be understood as a $S$ - wave $D\Lambda - D\Sigma$ bound state. Indeed, a $D\Lambda$ or $D\Sigma$ bound state with a mass about 2930 MeV that can be associated to the $\Xi_c(2923)^0$ is supported by Refs. [9–11]. The lattice QCD calculation was also performed and tried to determine their quantum numbers [12].

The successful explanation of $cud$ state [3–7] or molecular structure [8–11] for newly observed $\Xi_c^0$ states naturally leads us to ask what is their real structure? Moreover, there still exist many incomprehensible conclusions, such as the studies of conventional charm baryon cannot agree on their spin-parity [3–7]. More important is that the two-body allowed strong decay widths from different models are consistent with each other within errors. Hence, based on analysis of the two-body allowed strong decays, newly observed $\Xi_c^0$ states are not only considered as conventional three quark state [3–7], but also as $S$-wave molecular configuration [8]. Precise information on the radiative decay mechanism of these $\Xi_c^*$ states will be helpful to determine whether they are conventional charm baryon or molecular state. This idea is original from coupling of photon to molecule constituents is essentially different from that of the quark models, in which the photon couples directly to the quark system [13].

*Electronic address: 20132013@cqnu.edu.cn
Very recently, radiative decay widths of these three states are studied by assuming them as conventional charm baryons [3, 14]. It is helpful if we could estimate the radiative decay widths to make a comparison by considering newly observed \( \Xi^0 \) states as bound state. Thus, we can judge different explanations for structure of \( \Xi^* \) when there exist experimental signals. However, only \( \Xi_c(2923)^0 \) can be interpreted as a bound state [8–11] and, to date, no study addressed the radiative decays of \( \Xi_{c}(2923)^0 \). In the present study, we continue our investigation of \( \Xi_c(2923)^0 \) properties in terms of its radiative decay in hadronic molecule approach developed in our previous study [8].

This paper is organized as follows. In Sec. II, we will present the theoretical formalism. In Sec. III, the numerical result will be given, followed by discussions and conclusions in last section.

II. THEORETICAL FORMALISM

In our previous study [8], \( \Xi_c(2923)^0 \) was interpreted as an \( S \)-wave \( DA \) – \( DS \) molecular state, in which theoretical total decay width is consistent with experimental data [2]. In this work, the radiative decay \( \Xi_{c}(2923)^0 \to \Xi^0 \gamma \) are studied by assuming \( \Xi_{c}(2923)^0 \) as an \( S \)-wave \( DA \) – \( DS \) molecular state. The Feynman diagrams corresponding to the radiative decay are shown in Fig. 1.

\[
\mathcal{L}_{\Xi_c(2923)^0}(x) = \int d^4y \Phi(y) g_{\Xi_c(2923)^0 \gamma} D(x + \omega \gamma) \times Y(x - \omega \gamma) \Xi_c(2923)^0(x),
\]

where \( \omega_{ij} = m_i / (m_i + m_j) \) with \( m_i \) is the masses of \( D \) meson or \( Y \) baryon. Since \( \Sigma \) is an isovector baryon, \( \gamma \) should be replaced with \( \Sigma \cdot \tau \), where \( \tau \) is the isospin matrix. The correlation function \( \Phi(y^2) \) is introduced to describe the distribution of \( D \) and \( Y \) in the hadronic molecule \( \Xi_{c}(2923)^0 \). The more important role of \( \Phi(y^2) \) is to stop the Feynman diagrams ultraviolet infinite. Generally, \( y \) varies from 0 to \( +\infty \) and the amplitudes for the Feynman diagrams shown in Fig. 1 decrease to zero when \( y \to \infty \). Thus, \( \Phi(y^2) \) is often chosen to be of the following form

\[
\Phi(p^2) \equiv \exp(-p_E^2/\Lambda^2)
\]

where \( p_E \) being the Euclidean Jacobi momentum. \( \Lambda \) being the size parameter which characterizes the distribution of the components inside the molecule and it can only be determined from experimental data. In the following, it will be taken as a parameter and discussed later.

The coupling constant \( g_{\Xi_c(2923)^0 \gamma} \) is determined by the Weinberg compositeness rule [16, 17], which implies that the renormalization constant of wave function \( \Xi_c(2923)^0 \) is zero

\[
Z_{\Xi_c(2923)^0} = \chi_{DE} + \chi_{DA} - \frac{d\Sigma_{c}}{dk_0} |_{k_0=m_{\Xi_c(2923)^0}} = 0,
\]

where \( \chi_{DE} \) is possibility to find \( \Xi_c(2923)^0 \) in the molecule state \( DY \) with normalization \( \chi_{DE} + \chi_{DA} = 1.0 \). \( \Sigma_{\Xi_c(2923)^0} \) is self-energy of \( \Xi_c(2923)^0 \), which can be computed through the Feynman diagrams shown in Fig. 2.

\[
\Sigma_{\Xi_c(2923)^0}(k_0) = \sum_{Y=\Lambda,\Sigma^0,\Sigma^-} C_Y g_{\Xi_c(2923)^0 \gamma} \int \frac{d^4k_1}{(2\pi)^4} \times \frac{1}{p_F^2} \Phi^2 \left[ (k_1 - k_0 \omega) \right]^2 \left( \frac{1}{k_1^2 - m_D^2 (k_1 - k_0)^2 - m_F^2} \right),
\]

where \( k_1 \) and \( m_D \) is the four-momentum of baryon \( Y \) and the mass of meson \( D \), respectively. \( k_2 = m_{\Xi_c(2923)^0} \) is the four-momentum and the mass of \( \Xi_{c}(2923)^0 \), respectively. We set \( m_{\Xi_c(2923)^0} = m_D + m_Y - E_B \) with \( E_B \) being the bind energy of \( \Xi_{c}(2923)^0 \). \( C_Y \) is determined by the isospin symmetry.

\[
C_Y = \left\{ \begin{array}{lc} 1 & Y = \Lambda \\ \frac{1}{\sqrt{2} \bar{\beta}} & Y = \Sigma^0 \\ -\frac{1}{\sqrt{2} \bar{\beta}} & Y = \Sigma^- \end{array} \right.
\]

To compute radiative decay widths shown in Fig. 1, we need the effective Lagrangian densities relevant to vertices involving a photon field are naturally required [18, 19].

\[
\mathcal{L}_{D^0 \to D^+ \gamma} = \frac{e}{4} g_{D^0 \to D^+ \gamma} \epsilon^{\mu \nu \alpha \beta} F_{\mu \nu} D^{\alpha \beta}_0 D^+ + H.c.,
\]

\[
\mathcal{L}_{D^+ \to D^0 \gamma} = \frac{e}{4} g_{D^+ \to D^0 \gamma} \epsilon^{\mu \nu \alpha \beta} F_{\mu \nu} D^0 D^{\alpha \beta}_0 + H.c.,
\]

\[
\mathcal{L}_{\gamma \Lambda} = i e \bar{\gamma}_5 \gamma^\mu \gamma^\rho \gamma^\sigma \gamma^\lambda \bar{\Lambda} F_{\rho \sigma} \bar{D}_\mu D_\lambda + H.c.,
\]
The Lagrangian is expressed as \[ \mathcal{L}_{\gamma \Sigma \Lambda} = \frac{\epsilon_{\gamma \Sigma \Lambda}}{2m_{\Sigma}} \sigma_{\mu \nu} \partial^{\mu} A^{\nu} \Lambda + H.c. \] (10)
\[ \mathcal{L}_{\gamma \Sigma \Sigma} = -\bar{\Sigma}[\gamma_{\mu} A_{\mu} - \frac{\epsilon_{\gamma \Sigma \Sigma}}{2m_{\Sigma}} \sigma_{\mu \nu} \partial^{\mu} A^{\nu}] \Sigma, \] (11)
where the strength tensor is defined as \( F_{\mu \nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \), \( D^\alpha_\beta = \partial_\alpha D^\beta_\rho - \partial_\beta D^\alpha_\rho \), and \( \sigma_{\mu \nu} = i/2(\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu) \). The \( \alpha = e^2/4\pi = 1/137 \) is the electromagnetic fine structure constant. The coupling constants \( g_{D^\alpha D^\beta \gamma} = -0.5 \text{GeV}^{-1} \cdot g_{D^\alpha D^\beta \gamma} = 2\text{GeV}^{-1} \) \[18, 19\]. The anomalous and transition magnetic moments of the baryons are provided by the PDG \[1\] and are shown in Table I.

**TABLE I: Anomalous and transition magnetic moments.**

| \( \kappa_{\Sigma} \) | \( \kappa_{\Sigma^0} \) | \( \kappa_{\Sigma^+} \) |
|-----------------|-----------------|-----------------|
| -0.16           | 0.65            | 1.46            |
| \( \kappa_{\Lambda} \) | \( \mu_{\Sigma} \) | \( \mu_{\Sigma^0} \) |
| -0.61           | 1.61            |

Besides the Lagrangian above, we also require the effective Lagrangian describing coupling of vector meson to charmed baryon. Considering the \( SU(4) \) symmetry and the hidden gauge formalism, the Baryon-Baryon-Vector vertices Lagrangian is expressed as \[8, 20, 21\]
\[ \mathcal{L}_{VBV} = \frac{g_V}{4} \sum_{i,j,k=1}^{4} \bar{B}_{ijk} \gamma^\mu \left( V_{ij}^{\mu} B^{ji} + 2 V_{ij}^{\nu} B^{i\nu} \right), \] (12)
where coupling constant \( g_V = 6.6 \). \( V_\mu \) represents 16-plet vector fields of \( \rho \), and can be expressed as
\[ V_\mu = \begin{pmatrix}
\frac{1}{\sqrt{2}}(\rho^0 + \omega) & \rho^+ & K^+ & D^0 \\
-\frac{1}{\sqrt{2}}(\rho^0 - \omega) & K^0 & -D^- & D^+ \\
K^- & \bar{D}^- & \phi & D^0 \\
D^0 & -D^+ & D^- & \bar{D}^+ 
\end{pmatrix}_{\mu}. \] (13)

\( B \) is the tensor of baryons of the 20-plet of \( p \)
\[ B^{121} = p, \quad B^{122} = n, \quad B^{132} = \frac{1}{\sqrt{2}} \Sigma^0 - \frac{1}{\sqrt{6}} \Lambda, \]
\[ B^{213} = \frac{1}{\sqrt{2}} \Lambda, \quad B^{231} = \frac{1}{\sqrt{2}} \Sigma^0 + \frac{1}{\sqrt{6}} \Lambda, \quad B^{232} = \Sigma^-, \]
\[ B^{233} = \Xi^-, \quad B^{311} = \Sigma^+, \quad B^{313} = \Xi^0, \quad B^{141} = \Sigma^{++}, \]
\[ B^{142} = \frac{1}{\sqrt{2}} \Sigma^+ + \frac{1}{\sqrt{6}} \Lambda, \quad B^{143} = \frac{1}{\sqrt{2}} \Xi^- - \frac{1}{\sqrt{6}} \Xi^-, \]
\[ B^{241} = \frac{1}{\sqrt{2}} \Sigma^+ - \frac{1}{\sqrt{6}} \Lambda, \quad B^{242} = \Sigma^0. \]

**FIG. 2: Self energy of \( \Xi_c \) \((2923)\)**

\[ B^{243} = \frac{1}{\sqrt{2}} \Xi_c^0 + \frac{1}{\sqrt{6}} \Xi_c^-, \quad B^{341} = \frac{1}{\sqrt{2}} \Xi_c^+ + \frac{1}{\sqrt{6}} \Xi_c^-, \]
\[ B^{124} = \sqrt{\frac{2}{3}} \Lambda, \quad B^{234} = \sqrt{\frac{2}{3}} \Xi_c^+, \quad B^{314} = \sqrt{\frac{2}{3}} \Xi_c^-, \]
\[ B^{144} = \Xi_c^{++}, \quad B^{244} = -\Xi_c^+, \quad B^{444} = \Xi_c, \] (14)
where indices \( i, j, k \) of \( B^{ijk} \) denote quark content of a baryon with assignments \( 1 \leftrightarrow u, 2 \leftrightarrow d, 3 \leftrightarrow s, 4 \leftrightarrow c \), and first two of them are antisymmetric.

Moreover, \( \Xi_c^0 \Lambda D^0 \) and \( \Xi_c^0 \Sigma D^+ \) vertices are also required and can be obtained from Ref. \[22\]
\[ \mathcal{L}_{\Xi_c^0 \Lambda D} = \frac{i g_{\Xi_c^0 \Lambda D}}{m_{\Xi_c^0} + m_D} \tilde{\Sigma}^{[i} \gamma^\mu \gamma^5 \Lambda^{j]} \tilde{D} + H.c. \] (15)
\[ \mathcal{L}_{\Xi_c^0 \Sigma D} = \frac{i g_{\Xi_c^0 \Sigma D}}{m_{\Xi_c^0} + m_D} \tilde{\Sigma}^{[i} \gamma^\mu \gamma^5 \tilde{\Sigma} \partial_{[i} \tilde{D} + H.c. \] (16)

where \( \tilde{\Sigma} \) represents Pauli matrix, \( \tilde{\Sigma} \) is \( \Sigma \) triplets, and \( \tilde{D} \) is doubledts of charmed mesons. The relevant coupling constants are listed in Table II.

**TABLE II: Values of relevant meson-baryon coupling constants.**

| \( g_{\Xi_c^0 \Lambda D} \) | \( g_{\Xi_c^0 \Sigma D} \) | \( g_{\Xi_c^0 \Sigma D} \) | \( g_{\Xi_c^0 \Sigma D} \) |
|-----------------|-----------------|-----------------|-----------------|
| -5.38           | 6.43            | 9.31            | 3.71            |

Putting all the pieces together, we obtain amplitudes for Feynman diagrams shown in Fig. 2.

\[ M_d = \tilde{\mu}(p_1) \left( i \frac{\sqrt{3} e}{16} C_{\Sigma} g_{\Xi_c} g_{\Sigma} g_{D^0} g_{D^+} \right) \int \frac{d^4 k_1}{(2\pi)^4} \]
\[ \times \Phi((p \omega_D - q \omega_\Sigma)^2) \epsilon_{\mu \nu \rho \sigma} g_{\mu} g_{\nu} \frac{1}{p^2 - m_{\Sigma}^2} \frac{1}{q^2 - m_{D^+}^2} \]
\[ \times ((p_{2\mu} - q_{2\nu}) g_{\rho} g_{\sigma} - k_{1\mu} g_{\rho} g_{\sigma} - k_{1\rho} g_{\nu} g_{\sigma}) \frac{\epsilon_{\mu \nu \rho \sigma}}{k_1^2 - m_{D^0}^2} \]
\[ \times e_{\rho}(p_2) \mu(k_0), \] (17)
\[ M_b = \tilde{\mu}(p_1) \left( -i \frac{e}{(m_{\Xi_c^0} + m_D)} C_{\Xi_c} g_{\Xi_c} g_{\Sigma} g_{D^0} g_{D^+} \right) \int \frac{d^4 k_1}{(2\pi)^4} \]
\[ \times \Phi((p \omega_D - q \omega_\Sigma)^2) \epsilon_{\mu \nu \rho \sigma} g_{\mu} g_{\nu} \frac{1}{p^2 - m_{\Sigma}^2} \frac{1}{q^2 - m_{D^+}^2} \]
\[ \times \Phi((p \omega_D - q \omega_\Sigma)^2) \epsilon_{\mu \nu \rho \sigma} g_{\mu} g_{\nu} \frac{1}{p^2 - m_{\Xi_c}^2} \frac{1}{q^2 - m_{D^0}^2} (q^\rho + k_1^\rho). \]
\begin{align}
&\times \frac{1}{k_1^4 - m_{D^*}^2 - i\sqrt{6}e}{C_{\Sigma D^*} \mathcal{G}(\mathcal{G}_8 \not\partial \not\partial D^* \not\partial) \int \frac{d^4 k_1}{(2\pi)^4}} \\
&\times \Phi((p_2 \not\partial \not\partial - q_2 \not\partial \not\partial) (k_1 \not\partial \not\partial - k_2 \not\partial \not\partial) - g_{\not\partial \not\partial} + k_1 k_2 \not\partial \not\partial/m_{D^*}^2)} \\
&\times e_\gamma^\prime(p_2) \mu(k_0),
\end{align}

where \( \mathcal{M} \) and \( \mathcal{M}' \) is the amplitude for \( \Xi_c(2923)^0 \to \Xi_c^0 \not\partial \not\partial \) and \( \Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial \), respectively.

Summing up all the individual amplitudes, we obtain the total amplitude of \( \Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial \).

\begin{align}
\mathcal{M}(\Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial) &= \mathcal{M}_b(\Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial) + M_c(\Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial) + M_f(\Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial)
\end{align}

However, the \( \mathcal{M}' \) obtained currently is not satisfying the gauge invariance of the photon filed. Therefore, the contact diagram in Fig. 3 must be included to ensure the gauge invariance of total amplitudes. We employ the following form to satisfy \( p_2^2 \mathcal{M}_{\text{Born}}(\Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial) = 0 \).

\begin{align}
\Xi_c^{(0)}(k_0) &\rightarrow D(q) \\
\Lambda, \Sigma \not\partial \not\partial &\rightarrow \gamma(p_2)
\end{align}

FIG. 3: Contact diagram for \( \Xi_c(2923)^0 \to \Xi_c^{0} \not\partial \not\partial \). We also indicate definitions of the kinematics \((p_1, p_2, k_1, k_2, \not\partial \not\partial)\) used in the calculation

\begin{align}
\mathcal{M}_\text{com}(\Xi_c^{0} \to \Xi_c^{(0)} = \int \frac{e}{16 m_{\Xi^{0}} + m_{D^*}} \int C_{\Sigma D^*} \int \frac{d^4 k_1}{(2\pi)^4}} \\
\times \Phi((p_2 \not\partial \not\partial - q_2 \not\partial \not\partial) (k_1 \not\partial \not\partial - k_2 \not\partial \not\partial) - g_{\not\partial \not\partial} + k_1 k_2 \not\partial \not\partial/m_{D^*}^2)} \\
\times e_\gamma^\prime(p_2) \mu(k_0),
\end{align}

Where

\begin{align}
T_1 &= \exp[-\frac{1}{\beta^2}[\eta(\not\partial \not\partial^2 + m_{D^*}^2) + a m_{D^*}^2 + \not\partial \not\partial^2]], \\
T_2 &= \exp[-\frac{1}{\beta^2}[\eta(\not\partial \not\partial^2 + m_{D^*}^2) + a m_{D^*}^2 + \not\partial \not\partial^2]],
\end{align}

\begin{align}
C^{\prime} &= \int \frac{m^2 H^3}{4y^3} p_1^0 \frac{H^2 H^1}{4y^3} (m m^2 - m^2 \not\partial \not\partial^2) + \not\partial \not\partial^2] \\
&\times (m^2 - m m^2) p_1^0 + \frac{H^2 H^1}{2y^2} (m m^2 - m m^2) p_1^0 \\
&\times \frac{H^2 H^1}{2y^2} (m m^2 - m m^2) p_1^0
\end{align}
+ \frac{1}{y} (2p_1^\rho - m_1 y^\rho), \quad (28)
\frac{\mathcal{C}_2}{\Lambda} = \frac{\mathcal{H}_1 \mathcal{H}_2 m_\Sigma}{\Lambda_1^2} (m^2 - m m_1) y^\rho - \frac{\mathcal{H}_1 \mathcal{H}_2}{\Lambda_1^2} \left( m^2 - m m_1 \right) p_1^\rho
+ \frac{\mathcal{H}_1 \mathcal{H}_2}{\Lambda_1^2} p_1 \cdot p_2 (2p_1^\rho - m_1 y^\rho) - \frac{\mathcal{H}_1^2 \mathcal{H}_1}{\Lambda_1^2} p_1 \cdot p_2 (2p_1^\rho - m_1 y^\rho)
- m_1 y^\rho + \frac{m y}{\Lambda} y^\rho - \frac{3 \mathcal{H}_1}{\Lambda_1^2} (2p_1^\rho - m_1 y^\rho) - m_1^2 (2p_1^\rho - m_1 y^\rho)
+ \frac{\mathcal{H}_1^2 m_\Sigma^2 m_1}{\Lambda_1^2} y^\rho - \frac{\mathcal{H}_1 m_\Sigma m_1}{\Lambda_1^2} p_1^\rho - \frac{\mathcal{H}_1^2 m_\Sigma^2 m_1}{\Lambda_1^2} p_1^\rho
- \frac{\mathcal{H}_1^2 m_1^2 (2p_1^\rho - m_1 y^\rho)}{\Lambda_1^2} + \frac{\mathcal{H}_1^2 m_1^2 (2p_2^\rho - m_1 y^\rho)}{\Lambda_1^2}, \quad (29)

where \( y = 1 + \sigma + \eta + \zeta, \mathcal{H}_1 = 2(\eta + \omega_D), \mathcal{H}_1 = 2(\zeta + \omega_D) \), \( m \) and \( m_1 \) are the masses of \( \Xi_1^0 \) and \( \Xi_2^1 \), respectively.

Once the amplitudes are determined, the corresponding partial width can be obtained immediately with the following formula

\[ d\Gamma(\Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma) = \frac{1}{2J + 1} \frac{1}{32\pi^2} \frac{|\bar{p}_1|}{M_0^2} d\Omega. \quad (30) \]

Where \( J \) is the total angular momentum of \( \Xi_1(2923)^0 \), \( |\bar{p}_1| \) is the three-momentum of the decay products in the center of mass frame, and the overline indicates the sum over the polarization vectors of the final hadrons.

### III. RESULTS

To make a reliable prediction for the radiative decay widths of \( \Xi_1(2923)^0 \), two issues we need to clarify are, respectively, the relation of the parameter \( \Lambda \) to the correlation function \( \Phi(y) \) and the coupling of \( \Xi_1(2923)^0 \) with its molecular compositions. Unfortunately, the parameter \( \Lambda \) could not be determined by first principles, it is usually set to be about 1.0 GeV to reproduce the experimentally observed decay width in the literature [8, 15, 22–25] (and their references). In this work, we adopt the value \( \Lambda = 1.0 \) GeV because it is determined from the experimental data of Refs. [8, 15, 22–25] (and their references) within the same correlation function \( \Phi(y) \) adopted in current work.

The coupling constants corresponding to the effective Lagrangians listed in Eq. (1) have been determined. Although detailed results can find in Ref. [8], we will review them here again. With a value of \( \chi_{DS} = 0.0-1.0 \) and the compositeness condition that we introduced in Eq. (3), the \( \chi_{DS} \) dependence of the coupling constants \( g_{\Xi_1^0\Sigma} \) and \( g_{\Xi_1^0\Xi} \) are computed and are shown in Fig. 4. We find that the coupling constant \( g_{\Xi_1^0\Sigma} \) mono-nously decreases with increasing \( \chi_{DS} \), while the coupling constant \( g_{\Xi_1^0\Xi} \) increases with increasing \( \chi_{DS} \). The opposite trend can be easily understood, as the coupling constants \( g_{\Xi_1^0\Sigma} \) and \( g_{\Xi_1^0\Xi} \) are directly proportional to the corresponding molecular compositions [25, 26]. And a simple relation between the coupling constants \( g_{\Xi_1^0\Sigma} \) and \( g_{\Xi_1^0\Xi} \) can be deduced as

\[ (\mathcal{A} g_{\Xi_1^0\Sigma}^2)^2 = 1 - (B g_{\Xi_1^0\Sigma})^2 \]

with \( \mathcal{A} = 0.0885 \) and \( B = 0.0891 \).

FIG. 4: (color online) The coupling constant of \( g_{\Xi_1^0\Sigma} \) as a function of the parameter \( \chi_{DS} \).

Once the model parameter \( \Lambda \) and coupling constants \( g_{\Xi_1^0\Sigma} \) and \( g_{\Xi_1^0\Xi} \) are determined, the radiative decay width \( \Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma \) can be calculated straightforwardly. In Fig. 5, the radiative decay width \( \Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma \) of \( \chi_{DS} \) is presented, where we restrict the value of \( \chi_{DS} \) from 0.0 to 1.0. The results show that the value of the radiative decay width \( \Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma \) verses \( \chi_{DS} \) is the smallest for \( \chi_{DS} = 0.0 \) case, in this case \( \Xi_1(2923)^0 \) is a pure \( \Lambda \) molecular state. With the increase of \( \chi_{DS} \), the radiative decay width \( \Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma \) increases. However, the decay width for \( \Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma \) and \( \Xi_1(2923)^0 \rightarrow \Xi_2^1 \gamma \) decreases when \( \chi_{DS} \) varies from 0.92 to 1.0 and 0.94 to 1.0, respectively. Such dependence of the radiative decay width on \( \chi_{DS} \) can be easily understood due to the interference between \( \Lambda \) and \( \Sigma \) channels is sizable (see

FIG. 5: (color online) The partial decay widths from \( \Sigma \) and \( \Lambda \) contribution as a function of the parameter \( \chi_{DS} \).
Fig. 5), leading to a total radiative decay width is the largest. It means that the DA channel strongly couples to the DS channel, and the same conclusion can also find in Refs. [8–11].

The contributions of the DA and DS components for \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) decays are calculated and also presented in Fig. 5. We find that the DA channel contribution decreases with increasing \( \chi_{DE} \), while the contribution from DS channel increases with increasing \( \chi_{DE} \). Moreover, the DA channel plays a predominant role at small \( \chi_{DE} \), while the contribution from the DS channel becomes the most important when \( \chi_{DE} \) is larger than 0.11 for the process \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) and 0.07 for the process \( \Xi_c(2923)^0 \rightarrow \Xi^0_\gamma \).

The individual contributions of the \( D, D^*, \Sigma, \Lambda \) exchanges, and the contact term for \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) decays are calculated and shown in Fig. 6. According to the Lagrangian above, the relative signs of the Feynman diagrams shown in Fig. (1) are well defined. And the total decay widths obtained are the square of their coherent sum. We can find that the \( \Sigma \) exchange and contact term play a dominant role, while the \( D, D^*, \) and \( \Lambda \) exchanges give minor contributions. However, as \( \chi_{DE} \) gets bigger, the contribution from the contact term becomes the most important. The interferences among them are still sizable. And it makes the total radiative decay width at the order of 1.23–11.66 KeV. For \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) transition, the predicted total decay width increases from 0.30 to 3.71 KeV. The \( \Sigma \) and \( D \) exchanges provide the dominant contribution, while the contact term contribution that play a dominant role in the \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) reaction is small. The sizable interference between them still exists. We also find that the \( D^* \) contribution that is considered in the \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) reaction is not included for the process \( \Xi_c(2923)^0 \rightarrow \Xi^0_\gamma \).

It is interesting to compare our results with those in Ref. [3, 14]. According to Ref. [14], the \( \Xi_c(2923)^0 \) may be conventional charmed baryons with \( P \)-wave. The electromagnetic decay width \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) is found to be small and of the order of 1.1 KeV, whereas the decay width is about 742.0 KeV and 0.4 KeV for the transition \( \Xi_c(2923)^0 \rightarrow 2\Xi^0_\gamma \) and \( 4\Xi^0_\gamma \), respectively. The conclusion of the conventional charmed baryon for \( \Xi_c(2923)^0 \) is also supported by Ref [3]. They find that the radiative decay width of \( \Xi_c(2923)^0 \) into the ground state \( 2\Xi_c \) is about 0.0 KeV, while the transition \( \Xi_c(2923)^0 \rightarrow 2\Xi^0_\gamma \) play a dominant role. These are quite different from our results that the total radiative decay widths for \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) and \( \Xi_c(2923)^0 \rightarrow \Xi^0_\gamma \) is of the order of 1.23–11.66 KeV and 0.30–3.71 KeV, respectively, by assuming \( \Xi_c(2923)^0 \) as an \( S \)-wave \( DA – DS \) bound state. Reader can find these differentia in Tab. III. If measurements are in future experimental, these differences will be very useful to help us to test various interpretations of \( \Xi_c(2923)^0 \).

![Figure 6](image)

**FIG. 6:** (color online) The partial decay widths from \( D \) (red dash line), \( \Sigma \) (blue dot line), \( D^* \) (purple dash dot dot line), \( \Lambda \) (green dash dot line) exchange contribution and the contact term (black line) for the \( \Xi_c(2923)^0 \rightarrow \gamma \Xi^0 \) as a function of the parameter \( \chi_{DE} \).

| \( \Xi_c(2923)^0 \rightarrow \) | \( 2\Xi_c^0 + \gamma \) | \( 4\Xi_c^0 + \gamma \) | \( 2\Xi^0 + \gamma \) |
|------------------|------------------|------------------|------------------|
| Ref. [3]         | 472.0            | 1.0              | 0.0              |
| Ref. [14]        | 451.0            | 0.4              | 1.1              |
| This work        | 0.30 – 3.71      | 0.0              | 1.23 – 11.66     |

### IV. SUMMARY

Presently, there is not sufficient experimental information to determine the spin-parity of \( \Xi_c(2923)^0 \) state. Its properties, however, such as the spectroscopy and the strong decay widths, can be well explained in the context of the conventional charm baryon [3–7] state or molecular state [8–11]. Precise information on the radiative decay mechanism of \( \Xi_c(2923)^0 \) will be helpful to determine whether it is a conventional charm baryon or molecular state. Fortunately, the radiative decay widths of \( \Xi_c(2923)^0 \) have been studied by assuming \( \Xi_c(2923)^0 \) as conventional charm baryon [3, 14]. It is helpful if we could estimate the radiative decay width to make a comparison by assuming \( \Xi_c(2923)^0 \) as a molecular state. Thus, we can judge the different explanations for the structure of \( \Xi_c(2923)^0 \) if there exist experimental signals.

In the present study, we estimated the partial widths for the radiative decay from the \( \Xi_c(2923)^0 \) to the \( \Xi_c^0 \) and \( \Xi^0_\gamma \) state in a molecular scenario, in which the \( \Xi_c(2923)^0 \) is considered a \( DA – DS \) hadronic molecule. In the relevant parameter region, the partial widths are evaluated as

\[
\Gamma(\Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma) = 1.23 – 11.66 \text{ KeV},
\]
\[
\Gamma(\Xi_c(2923)^0 \rightarrow \Xi^0_\gamma) = 0.30 – 3.71 \text{ KeV}.
\]

Our estimations indicate that the partial widths for the transition \( \Xi_c(2923)^0 \rightarrow \Xi_c^0 \gamma \) are approximately one order of magnitude smaller than those of \( \Xi_c(2923)^0 \rightarrow \Xi^0_\gamma \). Our results are quite different from the results [3, 14] that are obtained by assuming \( \Xi_c(2923)^0 \) may be conventional charmed baryon. And can be used to test the (molecular) nature of the \( \Xi_c(2923)^0 \).
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