# Strong decay modes $\bar{K}\Xi$ and $\bar{K}\Xi\pi$ of the $\Omega(2012)$ in the $\bar{K}\Xi(1530)$ and $\eta\Omega$ molecular scenario

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We study the $\bar{K}\Xi$ decay mode of the newly observed $\Omega(2012)$ assuming that the $\Omega(2012)$ is a dynamically generated state with spin parity $J^P = 3/2^-$ from the coupled channel $S$-wave interactions of $\bar{K}\Xi(1530)$ and $\eta\Omega$. In addition, we calculate its three-body decay width into $K\pi\Xi$. It is shown that the so-obtained total decay width is in fair agreement with the experimental data. We compare our results with those of other recent studies and highlight the differences among them.

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## I. INTRODUCTION

Very recently, the Belle Collaboration observed an $\Omega$ excited state in the $\Xi^0 K^-$ and $\Xi^- K_S^0$ invariant mass distributions [1]. Its mass and width are determined to be $M = 2012.4 \pm 0.7 \pm 0.6$ MeV and $\Gamma = 6.4^{+2.5}_{-2.0} \pm 1.6$ MeV. The existence of such $\Omega$ excited states with a mass around 2000 MeV has already been predicted by various models, such as quenched quark models [2–5], the Skyrme model [6], and lattice gauge theory [7]. On the other hand, the extended quark models [8–10], where the instanton-induced quark-antiquark pair creation or Nambu-Jona-Lasinio interaction was employed, predicted $\Omega$ states with negative parity but lower masses, the lowest $\Omega$ state lying around 1800 MeV, about 200 MeV lower than those predicted in Refs. [2–7]. One of the reasons is that the $\Omega$ states in Refs. [8–10] have large five-quark components.

In Refs. [11–13], the interactions of the $\bar{K}\Xi(1530)$ and $\eta\Omega$ coupled channels were investigated in the chiral unitary approach. An $\Omega$ excited state with a mass around 2012 MeV and $J^P = 3/2^-$ can be dynamically generated by use of a reasonable subtraction constant.

After the observation of the $\Omega(2012)$, its two-body strong decays were studied within the chiral quark model in Ref. [14], where it was shown that the newly observed $\Omega(2012)$ could be assigned to the $J^P = 3/2^-$ three-quark state. In Ref. [15], the mass and residue of the $\Omega(2012)$ were calculated by employing the QCD sum rule method with the conclusion that the $\Omega(2012)$ could be a $1P$ orbitally excited $\Omega$ state with $J^P = 3/2^-$. The analysis of Ref. [15] was extended in Ref. [16] to study the $\Omega(2012) \to K^-\Xi^0$ decay. In Ref. [17], the authors performed a flavor $SU(3)$ analysis and concluded that the preferred $J^P$ for the $\Omega(2012)$ is also $3/2^-$. In Refs. [18,19], its strong decay modes were studied assuming that the $\Omega(2012)$ is a $\bar{K}\Xi(1530)$ hadronic molecule. We note that the hadronic molecular picture plays an important role in understanding the newly observed but unexpected states [20].

In this work, we take the chiral unitary approach and assume that the $\Omega(2012)$ is a dynamically generated state from the $\bar{K}\Xi(1530)$ and $\eta\Omega$ interactions. The coupling of the $\Omega(2012)$ to $\bar{K}\Xi(1530)$ is then obtained from the residue at the pole position. We then calculate its decay into $K\Xi$ via a triangle diagram. We also calculate the three-body partial decay width of the $\Omega(2012)$ into $K\Xi\pi$. The total decay width compares favorably with the experimental data [1] and agrees with other theoretical approaches.

This work is organized as follows. In Sec. II, we briefly explain the chiral unitary approach and calculate the two- and

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three-body decays of the $\Omega(2012)$. The results and discussion are presented in Sec. III, followed by a short summary in Sec. IV.

## II. FORMALISM

### A. $\Omega(2012)$ as a $\bar{K}\Xi(1530)$ and $\eta\Omega$ molecule

The mass of the $\Omega(2012)$ is slightly below the $\bar{K}\Xi(1530)$ threshold. It is natural to treat it as a $\bar{K}\Xi(1530)$ molecular state, dynamically generated from the interaction of the coupled channels $\bar{K}\Xi(1530)$ and $\eta\Omega$ in the isospin $I=0$ sector. However, the possibility to be an $I=1$ molecule cannot be excluded [17]. Within the chiral unitary approach, the interaction of the coupled channels $\bar{K}\Xi(1530)$ and $\eta\Omega$ in the strangeness $-3$ and isospin 0 was first studied in Ref. [12], where a pole at $(2141-i38)$ MeV was found with a natural subtraction constant $a_\mu=-2$ and a renormalization scale $\mu=700$ MeV. Later, it was explicitly shown in Ref. [13] that the pole position of the $3/2^-$ $\Omega$ state can shift by varying $a_\mu$. If we take $a_\mu=-2.5$ and $\mu=700$ MeV, we obtain a pole at $z_R=(2012.7, i0)$ MeV, which can be associated to the newly observed $\Omega(2012)$ state [1]. In the cutoff regularization scheme, the corresponding momentum cutoff is $\Lambda=726$ MeV, which seems quite natural as well.

The couplings of the bound state to the coupled channels $\bar{K}\Xi(1530)$ (channel 1) and $\eta\Omega$ (channel 2) can be obtained from the residue of the scattering amplitude at the pole position $z_R$, which reads

$$T_{ij}=\frac{g_{ii}g_{jj}}{\sqrt{s}-z_R}, \tag{1}$$

where $g_{ii}$ is the coupling of the state to the $i$th channel. One finds with $a=-2.5$ and $\mu=700$ MeV,

$$g_{11}=1.65, \qquad g_{22}=2.80. \tag{2}$$

By taking these coupling constants $g_{11}^2$ and $g_{22}^2$ that we obtained above and the loop functions $G_{\bar{K}\Xi^*}$ and $G_{\eta\Omega}$, we find that $X_{11}=-g_{11}^2\frac{\partial G_{\bar{K}\Xi^*}(\sqrt{s})}{\partial\sqrt{s}}|_{\sqrt{s}=z_R}=0.48$ and $X_{22}=-g_{22}^2\frac{\partial G_{\eta\Omega}(\sqrt{s})}{\partial\sqrt{s}}|_{\sqrt{s}=z_R}=0.29$. Thus, about 50% of the sum rule comes from the $\bar{K}\Xi^*$ channel, while the $\eta\Omega$ channel is also important. These results are different from those values obtained in Ref. [12] where the cutoff method for the loop functions is taken. However, if we used the cutoff method, we get $X_{11}=0.59$ and $X_{22}=0.15$, which is similar with those results obtained in Ref. [12]. In addition, the product $G_{11}G_{22}$ can be evaluated from the pole position [12]. If we take $\mu=700$ MeV, and $\alpha_\mu=-2.17$ and $\alpha_\mu=-2.65$ for the $\bar{K}\Xi^*$ and $\eta\Omega$ channels, respectively, we get a pole at $z_R=(2012.3, i0)$ MeV and $X_{11}=0.61$ and $X_{22}=0.19$. We can see that the $\bar{K}\Xi^*$ channel is always dominant, but the $\eta\Omega$ channel also gives a non-negligible contribution. However, one should note that without the $\eta\Omega$ channel, there will be no dynamically generated state because the interaction in the $K\Xi(1530)$ and $\eta\Omega$ coupled channels is off diagonal in the chiral unitary approach.

Then, one can write down the effective interaction of $\Omega(2012)\bar{K}\Xi(1530)$ ($\equiv\Omega^*\bar{K}\Xi^*$) and $\Omega(2012)\eta\Omega$ ($\equiv\Omega^*\eta\Omega$),

$$v_{\Omega^*\bar{K}\Xi^*}=\frac{g_{11}}{\sqrt{2}}\bar{\Xi}^{*\mu}\Omega^*_\mu\phi_{\bar{K}}, \tag{3}$$

$$v_{\Omega^*\eta\Omega}=g_{22}\bar{\Omega}^\mu\Omega^*_\mu\phi_\eta. \tag{4}$$

It is worth to mention that the $g_{\Omega^*(2012)\bar{K}\Xi(1530)}$ obtained in Ref. [19] with the assumption that the $\Omega^*(2012)$ is a pure $S$-wave $\bar{K}\Xi(1530)$ hadronic molecule with spin parity $3/2^-$ is about 2.24, which is different from ours, since we have also taken into account the $\eta\Omega$ channel. For the $\Omega^*(2012)\to\bar{K}\pi\Xi$ three-body decay, since only the $\bar{K}\Xi(1530)$ component contributes at tree level and the partial decay width is proportional to $g^2_{\Omega^*(2012)\bar{K}\Xi(1530)}$, our three-body decay width is almost the same as that of Ref. [19].

We note that based on the Weinberg-Salam compositeness condition [21–24], a fully consistent quantum-field approach has been developed by the Tübingen-Beijing group for the study of the exotic meson [25–39] and baryon states [40–43]. In our present work based on the unitary chiral theory, the $\eta\Omega$ channel is important and cannot be ignored. It will be interesting to see a future study of the $\Omega(2012)$ in the hadronic molecular approach of Refs. [25–44].

### B. The $\Omega(2012)\to\bar{K}\Xi$ and $\bar{K}\Xi\pi$ decays

In the present work, we assume that the $\Omega(2012)$ ($\equiv\Omega^{*-}$) exists and has a mass as that reported by the Belle Collaboration, and we study its strong decays to the two-body final state $\bar{K}\Xi$ assuming that it is a molecular state of $\bar{K}\Xi(1530)$ and $\eta\Omega$, as predicted by the chiral unitary approach [12]. Then, the $\Omega(2012)\to\bar{K}\Xi$ decay can proceed through the triangular diagrams as shown in Figs. 1(a)–1(e), where the $\Sigma$, $\Lambda$, $\rho$, $\omega$, $\phi$, and $K^*$ exchanges are considered. Note that by considering those diagrams of Fig. 1, we can easily show that the partial decay widths of $\Omega^{*-}\to K^-\Xi^0$ and $\Omega^{*-}\to K^0\Xi^-$ are the same, and we will explicitly calculate the decay of $\Omega^{*-}\to K^-\Xi^0$ in the following. Compared with the decays to the two final states, the contribution to the three-body decay $\Omega(2012)\to\bar{K}\Xi\pi$ only comes from the $\bar{K}\Xi(1530)$ component. The decay width of the $\Xi(1530)$ listed in Ref. [45] is around 9 MeV, and the partial decay width $\Xi(1530)\to\Xi\pi$ is the largest and almost saturates its total decay width. Therefore, we only compute the three-body decay through the decay of the $\Xi(1530)$, and the simplest Feynman diagram is shown in Fig. 1(f). Considering the quantum numbers and

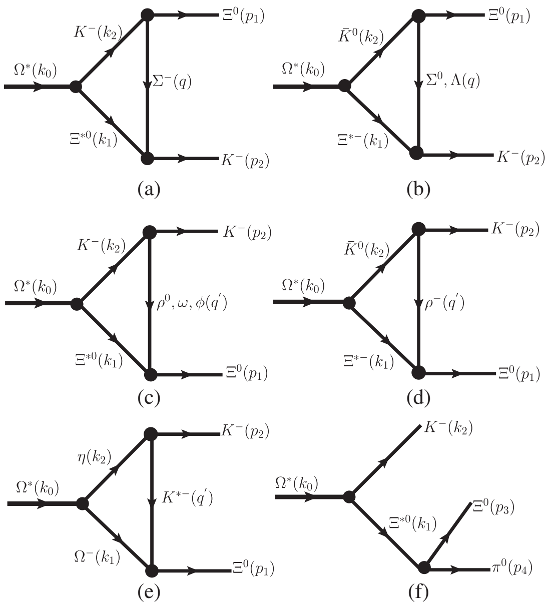


FIG. 1. Feynman diagrams contributing to the decays of the $\Omega(2012)$ to $\bar{K}\Xi$ and $\bar{K}\Xi\pi$.

phase space, in addition to the final state $K^-\Xi^0\pi^0$ shown in Fig. 1(f), the final states $\bar{K}^0\Xi^-\pi^0$, $\bar{K}^0\Xi^0\pi^-$, and $K^-\Xi^-\pi^+$ should also be calculated.

In order to evaluate the decay amplitudes of the diagrams shown in Fig. 1, we need the following effective Lagrangians to calculate the relevant vertices [46–48],

$$\mathcal{L}_{\Xi\Lambda K} = \frac{ig_{\Xi\Lambda K}}{m_\Lambda + m_\Xi}\partial^\mu \bar{K}\,\bar{\Lambda}\gamma_\mu\gamma_5\Xi + \text{H.c.}, \tag{5}$$

$$\mathcal{L}_{\Xi\Sigma K} = \frac{ig_{\Xi\Sigma K}}{m_\Sigma + m_\Xi}\partial^\mu \bar{K}\,\bar{\Sigma}\cdot\tau\gamma_\mu\gamma_5\Xi + \text{H.c.}, \tag{6}$$

$$\mathcal{L}_{\Xi^*\Sigma K} = \frac{f_{\Xi^*\Sigma K}}{m_K}\partial^\nu \bar{K}\bar{\Xi}^*_\nu\gamma_5\tau\cdot\Sigma + \text{H.c.}, \tag{7}$$

$$\mathcal{L}_{\Xi^*\Lambda K} = \frac{f_{\Xi^*\Lambda K}}{m_K}\partial^\nu \bar{K}\bar{\Xi}^*_\nu\gamma_5\Lambda + \text{H.c.}, \tag{8}$$

$$\mathcal{L}_{KK\rho} = ig_{\rho KK}[\bar{K}\partial_\mu K - \partial_\mu\bar{K}K]\vec{\tau}\cdot\vec{\rho}^\mu, \tag{9}$$

$$\mathcal{L}_{KK\omega} = ig_{\omega KK}[\bar{K}\partial_\mu K - \partial_\mu\bar{K}K]\omega^\mu, \tag{10}$$

$$\mathcal{L}_{KK\phi} = ig_{\phi KK}[\bar{K}\partial_\mu K - \partial_\mu\bar{K}K]\phi^\mu, \tag{11}$$

$$\mathcal{L}_{\eta KK^*} = ig_{\eta KK^*}[\bar{K}\partial_\mu\eta - \partial_\mu\bar{K}\eta]K^{*\mu}, \tag{12}$$

$$\mathcal{L}_{\rho\Xi\Xi^*} = i\frac{g_{\rho\Xi\Xi^*}}{m_\rho}\bar{\Xi}^{*\mu}\gamma^\nu\gamma_5[\partial_\mu\vec{\tau}\cdot\vec{\rho}_\nu - \partial_\nu\vec{\tau}\cdot\vec{\rho}_\mu]\Xi + \text{H.c.}, \tag{13}$$

$$\mathcal{L}_{\omega\Xi\Xi^*} = i\frac{g_{\omega\Xi\Xi^*}}{m_\omega}\bar{\Xi}^{*\mu}\gamma^\nu\gamma_5[\partial_\mu\omega_\nu - \partial_\nu\omega_\mu]\Xi + \text{H.c.}, \tag{14}$$

$$\mathcal{L}_{\phi\Xi\Xi^*} = i\frac{g_{\phi\Xi\Xi^*}}{m_\phi}\bar{\Xi}^{*\mu}\gamma^\nu\gamma_5[\partial_\mu\phi_\nu - \partial_\nu\phi_\mu]\Xi + \text{H.c.}, \tag{15}$$

$$\mathcal{L}_{K^*\Xi\Omega} = i\frac{g_{K^*\Xi\Omega}}{m_{K^*}}\bar{\Omega}^\mu\gamma^\nu\gamma_5[\partial_\mu K^*_\nu - \partial_\nu K^*_\mu]\Xi + \text{H.c.}, \tag{16}$$

$$\mathcal{L}_{\pi\Xi\Xi^*} = \frac{g_{\pi\Xi\Xi^*}}{m_\pi}\bar{\Xi}\partial^\mu\vec{\tau}\cdot\vec{\pi}\Xi^*_\mu + \text{H.c.} \tag{17}$$

Within the $SU(3)$ symmetry, we determine the coupling constants to be $g_{\eta\bar{K}\bar{K}^*} = \sqrt{3}g_{\pi\bar{K}\bar{K}^*} = 5.23$ [49], $g_{\phi\Xi\Xi^*} = g_{\rho\Xi\Xi^*} = g_{\omega\Xi\Xi^*} = \frac{1}{\sqrt{6}}g_{\rho\Delta N} = -2.48$ [50], and $g_{K^*\Xi\Omega} = -7.01$ [50]. The coupling constant $g_{\pi\Xi\Xi^*} = 2.24(g_{\pi\Xi\Xi^*} = 2.04)$ is evaluated using Eq. (17), and the partial decay width $\Gamma_{\Xi^*\to\Xi\pi} = \Gamma_{\Xi^*} = 9.1$ MeV$(\Gamma_{\Xi^*\to\Xi\pi} = \Gamma_{\Xi^*} = 9.9$ MeV) in the $\Xi^*$ rest frame. The masses of the particles needed in the present work are listed in Table I. The other coupling constants are taken from Refs. [46–48] and are listed in Table II.

With the above effective interaction Lagrangians and the coupling constants, we obtain the following decay amplitudes for $\Omega^*(2012)\to K^-\Xi^0$ and $\Omega^*(2012)\to\bar{K}\Xi\pi$ corresponding to the diagrams shown in Fig. 1,

$$\mathcal{M}_{\Sigma^-} = \frac{\sqrt{2}g_{11}g_{\Xi^0\Sigma^-K^-}g_{\Xi^{*0}\Sigma^-K^-}}{m_{K^-}(m_{\Sigma^-}+m_{\Xi^0})}\int\frac{d^4q}{(2\pi)^4}\bar{u}(p_1)\not{k}_2\gamma_5 \times\frac{1}{\not{q}-m_{\Sigma^-}+i\epsilon}\times\gamma_5\mathcal{P}^{\nu\eta}u^\eta_{\Omega^*}(k_0)p_2^\nu\frac{1}{k_2^2-m_{K^-}^2+i\epsilon}, \tag{18}$$

$$\mathcal{M}_{\Sigma^0} = \frac{g_{11}g_{\Xi^0\Sigma^0\bar{K}^0}g_{\Xi^{*-}\Sigma^0K^-}}{\sqrt{2}m_{K^-}(m_{\Sigma^0}+m_{\Xi^-})}\int\frac{d^4q}{(2\pi)^4}\bar{u}(p_1)\not{k}_2\gamma_5 \times\frac{1}{\not{q}-m_{\Sigma^0}+i\epsilon}\times\gamma_5\mathcal{P}^{\nu\eta}u^\eta_{\Omega^*}(k_0)p_2^\nu\frac{1}{k_2^2-m_{K^-}^2+i\epsilon}, \tag{19}$$

TABLE I. Masses of the particles needed in the present work (in units of MeV).

| $\pi^\pm$ | $\pi^0$ | $\eta$ | $\rho$ | $\omega$ | $\phi$ | $K^\pm$ |
|---|---|---|---|---|---|---|
| 139.57 | 134.98 | 547.86 | 775.26 | 782.65 | 1019.46 | 493.68 |
| $K^0$ | $K^*$ | $\Xi^0$ | $\Xi^-$ | $\Xi^0(1530)$ | $\Xi^-(1530)$ | |
| 497.61 | 891.76 | 1314.86 | 1321.71 | 1531.80 | 1535.00 | |

TABLE II. Values of the effective meson-baryon and meson-meson coupling constants.

| $g_{KK\phi}$ | $g_{KK\rho}$ | $g_{KK\omega}$ | $g_{\Xi^*\Sigma K}$ | $g_{\Xi^*\Lambda K}$ | $g_{\Xi\Sigma K}$ | $g_{\Xi\Lambda K}$ |
|---|---|---|---|---|---|---|
| −3.02 | −3.02 | −3.02 | 3.22 | 5.58 | −13.26 | 3.37 |

$$\mathcal{M}_{\Lambda} = \frac{g_{11} g_{\Xi^0\Lambda\bar{K}^0} g_{\Xi^{*-}\Lambda K^-}}{\sqrt{2} m_{K^-}(m_\Lambda + m_{\Xi^-})} \int \frac{d^4 q}{(2\pi)^4} \bar{u}(p_1) \not{k}_2 \gamma_5 \frac{1}{\not{q} - m_\Lambda + i\epsilon} \times \gamma_5 \mathcal{P}^{\nu\eta} u^{\eta}_{\Omega^*}(k_0) p_2^{\nu} \frac{1}{k_2^2 - m_{K^-}^2 + i\epsilon}, \tag{20}$$

$$\mathcal{M}_{\rho^0/\omega/\phi} = \frac{-i g_{11} g_{\Xi^*\Xi\rho/\omega/\phi} g_{\rho/\omega/\phi\bar{K}\bar{K}}}{\sqrt{2} m_{\rho/\omega}} \int \frac{d^4 q'}{(2\pi)^4} \bar{u}(p_1)\gamma^{\mu}\gamma_5 \times \mathcal{P}^{\nu\eta} u^{\eta}_{\Omega^*}(k_0)(q'_{\nu} g_{\mu\alpha} - q'_{\mu} g_{\nu\alpha})(k_1^{\beta} - p_2^{\beta}) \times \frac{1}{k_2^2 - m_{K^-}^2 + i\epsilon} \frac{-g^{\alpha\beta} + q'^{\alpha} q'^{\beta}/m^2_{\rho/\omega}}{q'^2 - m^2_{\rho/\omega/\phi} + i m_{\rho/\omega/\phi}\Gamma_{\rho/\omega/\phi}}, \tag{21}$$

$$\mathcal{M}_{\rho^-} = \frac{-i\sqrt{2} g_{11} g_{\Xi^*\Xi\rho} g_{\rho\bar{K}\bar{K}}}{m_\rho} \int \frac{d^4 q'}{(2\pi)^4} \bar{u}(p_1)\gamma^{\mu}\gamma_5 \times \mathcal{P}^{\nu\eta} u^{\eta}_{\Omega^*}(k_0)(q'_{\nu} g_{\mu\alpha} - q'_{\mu} g_{\nu\alpha})(k_1^{\beta} - p_2^{\beta}) \times \frac{1}{k_2^2 - m_{K^-}^2 + i\epsilon} \frac{-g^{\alpha\beta} + q'^{\alpha} q'^{\beta}/m^2_{\rho}}{q'^2 - m^2_{\rho} + i m_{\rho}\Gamma_{\rho}}, \tag{22}$$

$$\mathcal{M}_{K^{*-}} = \frac{-i g_{22} g_{\eta K^- K^{*-}} g_{\Omega^- K^{*-}\Xi^0}}{m_{K^{*-}}} \int \frac{d^4 q'}{(2\pi)^4} \bar{u}(p_1)\gamma^{\mu}\gamma_5 \times \mathcal{P}^{\nu\eta} u^{\eta}_{\Omega^*}(k_0)(q'_{\nu} g_{\mu\alpha} - q'_{\mu} g_{\nu\alpha})(k_1^{\beta} - p_2^{\beta}) \times \frac{1}{k_2^2 - m_{\eta}^2 + i m_\eta \Gamma_\eta} \frac{-g^{\alpha\beta} + q'^{\alpha} q'^{\beta}/m^2_{K^*}}{q'^2 - m^2_{K^*} + i m_{K^*}\Gamma_{K^*}}, \tag{23}$$

$$\mathcal{M}_{\bar{K}(k_2)\Xi(p_3)\pi(p_4)} = f_I \frac{i g_{\pi\Xi\Xi^*} g_{\Omega^*\bar{K}\Xi^*}}{m_\pi} \bar{u}(p_3) p_3^{\nu} \mathcal{P}^{\nu\eta} u_{\eta}(k_0), \tag{24}$$

where

$$\mathcal{P}^{\nu\eta} = \frac{\not{k}_1 + m_{\Xi^*}}{k_1^2 - m_{\Xi^*}^2 + i m_{\Xi^*}\Gamma_{\Xi^*}} \left[ -g^{\nu\eta} + \frac{1}{3}\gamma_{\nu}\gamma_{\eta} + \frac{2}{3m_{\Xi^*}^2} k_{1\nu} k_{1\eta} + \frac{1}{3m_{\Xi^*}}(\gamma^{\nu} k_{1\eta} - \gamma^{\eta} k_{1\nu}) \right], \tag{25}$$

with $q = k_2 - p_1 = p_2 - k_1$ and $q' = k_2 - p_2 = p_1 - k_1$. We take $\Gamma_{\Xi^{*0}} = 9.1$ MeV, $\Gamma_{\Xi^{*-}} = 9.9$ MeV, $\Gamma_\rho = 149.1$ MeV, $\Gamma_\omega = 8.5$ MeV, $\Gamma_\eta = 1.3$ MeV, $\Gamma_\phi = 4.2$ MeV, and $\Gamma_{K^*} = 50.3$ MeV. For the three-body final states $K^-\Xi^0\pi^0$ and $\bar{K}^0\Xi^-\pi^0$, the isospin factors are $f_I = 1$ and $-1$, respectively. The isospin factor is $f_I = \sqrt{2}$ for all the other three-body final states. To take into account the finite size of hadrons, for each vertex, we multiply a form factor $F(k_1^2)$ of the following form [19]

$$F(k_1^2) = \frac{\Lambda^4}{(m^2 - k_1^2)^2 + \Lambda^4}, \tag{26}$$

where $m$ is the mass of the exchanged particle and $k_1$ is its momentum, with the cutoff $\Lambda$ varying from 0.8 to 2.0 GeV.

In order to avoid ultraviolet divergence in the triangle diagrams, we take the three-momentum truncation method to compute the amplitudes. In the rest frame of the $\Omega^*$, the relevant momenta are defined as follows:

$$k_0 = (M, 0, 0, 0), \qquad p_1 = (E_\Xi, 0, 0, p_{cm}), \tag{27}$$

$$p_2 = (E_K, 0, 0, -p_{cm}), \quad q = (q_0, q_1 \sin\theta, 0, q_1\cos\theta) = q', \tag{28}$$

and we can rewrite the $\int d^4 q^{(\prime)} = \int dq_0 q_1^2 dq_1 d\cos\theta d\phi$ with $q_0 \epsilon(-\infty, \infty)$, $q_1 \epsilon(0, \Lambda)$, $\cos\theta \epsilon(-1, 1)$, and $\phi \epsilon(0, 2\pi)$, where $\Lambda$ is a free parameter and is also taken to vary from 0.8 to 2.0 GeV. Here, we have introduced a cutoff $\Lambda$ in the three-momentum integration.[1]

The partial decay width of the $\Omega^* \to \bar{K}\Xi$ and $\Omega^* \to \bar{K}\Xi\pi$ in its rest frame are given by

$$d\Gamma_{\Omega^*\to\bar{K}\Xi} = \frac{1}{32\pi^2} \overline{|\mathcal{M}|^2} \frac{p}{M^2} d\Omega, \tag{29}$$

$$d\Gamma_{\Omega^*\to\bar{K}\Xi\pi} = \frac{1}{(2\pi)^5} \frac{1}{16M^2} \overline{|\mathcal{M}|^2} |\vec{p}_3^{\,*}| \times |\vec{k}_2| dm_{\pi\Xi} d\Omega^*_{p_3} d\Omega_{k_2}, \tag{30}$$

where $M$ is the mass of the $\Omega(2012)$, while $p$ is the module of the $\Xi$ (or $\bar{K}$) three-momentum in the rest frame of the $\Omega(2012)$. The $(|\vec{p}_3^{\,*}|, \Omega^*_{p_3})$ is the momentum and angle of the particle $\Xi$ in the rest frame of $\Xi$ and $\pi$, and $\Omega_{k_2}$ is the angle of the $\bar{K}$ in the rest frame of the decaying particle. The $m_{\pi\Xi}$ is the invariant mass for $\pi$ and $\Xi$ and $m_\pi + m_\Xi \le m_{\pi\Xi} \le M - m_{\bar{K}}$. The averaged squared amplitude is then

$$\overline{|\mathcal{M}|^2} = \frac{1}{4} \sum_{s_{\Omega^*}} \sum_{s_\Xi} |\mathcal{M}|^2, \tag{31}$$

where

[1] We have checked that using the transition form factors such as those of Refs. [25–44] to regulate the amplitudes yields qualitatively the same results as those presented in the present work. We noticed that Ref. [19] adopted a three-dimensional version of the transition factors.

$$\mathcal{M}_{\Omega^*\to\bar{K}\Xi} = \mathcal{M}_{\Sigma/\Lambda} + \mathcal{M}_{\rho/\omega/\phi} + \mathcal{M}_{K^*}, \tag{32}$$

$$|\mathcal{M}|^2_{\Omega^*\to\bar{K}\Xi\pi} = |\mathcal{M}|^2_{K^-\Xi^0\pi^0} + |\mathcal{M}|^2_{K^-\Xi^-\pi^+} + |\mathcal{M}|^2_{\bar{K}^0\Xi^0\pi^-} + |\mathcal{M}|^2_{\bar{K}^0\Xi^-\pi^0}. \tag{33}$$

## III. NUMERICAL RESULTS AND DISCUSSION

As explained in the previous section, the triangle diagrams are regularized with a sharp momentum cutoff $\Lambda$, which is taken to be the same as those appearing in the form factor, $F(k_1^2)$. Because the triangle diagrams are ultraviolet divergent, our two-body decay width will depend on the value of the cutoff. Therefore, it is important to check whether one can obtain a decay width consistent with the experimental data using a reasonable value for the cutoff.

In Fig. 2, we show the total decay width of $\Omega(2012) \to \bar{K}\Xi$ as a function of the cutoff parameter $\Lambda$. Note that the $\Omega(2012) \to \bar{K}\pi\Xi$ three-body decay does not depend on the cutoff parameter $\Lambda$, but it depends weakly on the parameter $\Lambda_0$ as in Ref. [19]. We can see that the $\bar{K}\Xi(1530)$ component provides the dominant contribution to the partial decay width of the $\bar{K}\Xi$ two-body channel. The $\eta\Omega$ contribution to the $\bar{K}\Xi$ two-body channel is very small. However, the interference between them is still sizable and increases with the cutoff parameter $\Lambda$.

In Refs. [17,19], the three-body decay was emphasized, while our result shows that two-body $\bar{K}\Xi$ decay width can become larger than the $\bar{K}\pi\Xi$ three-body decay width when $\Lambda$ is larger than 1.65 GeV. We note that the hyperon exchange and $\eta\Omega$ component contribution are not considered in Ref. [19]. More specifically, our three-body partial decay width, ∼3.0 MeV, is close to the estimate of Ref. [19] but smaller than that of Ref. [17], about 10 MeV. We note that our total decay width and that of Ref. [18] agree with the experimental data. The difference is that Ref. [18] assumes that the $\Omega^*(2012)$ is a pure $\Xi(1530)\bar{K}$ molecule and invokes some power counting arguments to calculate its two-body decay width. Indeed, our study shows that the contribution from the $\eta\Omega$ component is small. Note that the molecule picture is different from the $qqq$ picture of the chiral quark model [14] and light cone QCD sum rules [15,16].

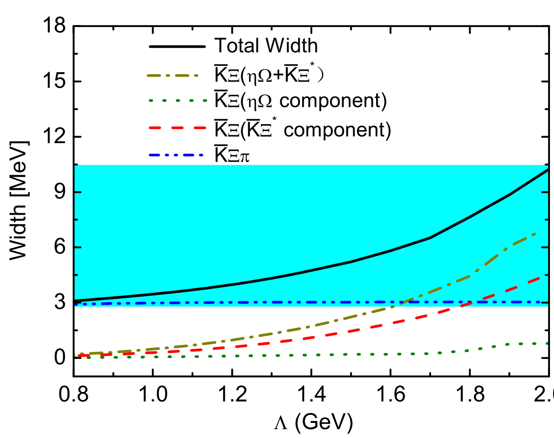


FIG. 2. Total decay width of the $\Omega(2012)$ as a function of the parameter $\Lambda$. The cyan error bands correspond to the experimental decay width [1].

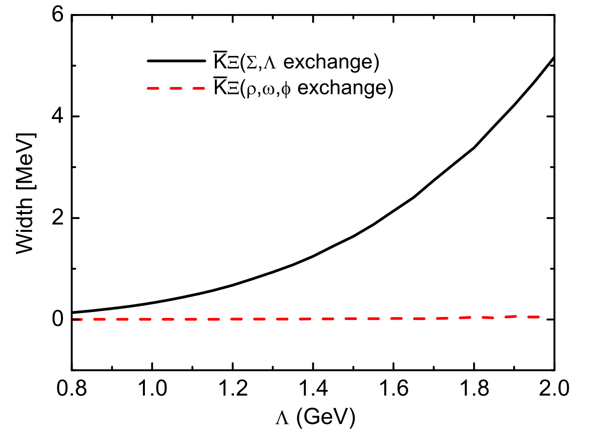


FIG. 3. Decomposed contributions to the decay width of the $\Omega(2012)$ into $K\Xi$ as a function of the parameter $\Lambda$

The contribution of the $\bar{K}\Xi(1530)$ component includes two parts: (i) $\Sigma$ and $\Lambda$ exchanges [Figs. 1(a) and 1(b)] and (ii) $\rho$, $\phi$, and $\omega$ meson exchanges [Figs. 1(c) and 1(d)]. The relative importance of these two mechanisms to the $\Omega^* \to \bar{K}\Xi$ decay is shown in Fig. 3. One can see that the contribution from the $\Sigma$ and $\Lambda$ exchanges is larger than those from $\rho$, $\phi$, and $\omega$ meson exchanges for the cutoff range studied.

## IV. SUMMARY

In summary, we studied the $\bar{K}\Xi$ decay of the newly observed $\Omega^*(2012)$ assuming that the $\Omega(2012)$ is a dynamically generated state with spin parity $3/2^-$ from the coupled channel interactions of $\bar{K}\Xi(1530)$ and $\eta\Omega$ in $S$ wave. Taking $\alpha_\mu = -2.5$ and $\mu = 0.7$ GeV, we obtained a pole at $M = 2012.7$ MeV and associated it to the newly observed $\Omega^*(2012)$. With the coupling constants between the $\Omega^*(2012)$ and its components calculated from the residue at the pole position $g_{\Omega^*\bar{K}\Xi^*} = 1.65$ and $g_{\Omega^*\eta\Omega} = 2.80$, we obtained the partial decay widths of the $\bar{K}\Xi$ final state through triangle diagrams in an effective Lagrangian approach. In such a picture, the decay $\Omega^*(2012) \to \bar{K}\Xi$ occurs by exchanging $\Sigma$, $\Lambda$ hadrons and $\rho$, $\phi$, $\omega$, and $K^*$ vector mesons. The contribution to the three-body decay $\Omega(2012) \to \bar{K}\Xi\pi$ only comes from the $\bar{K}\Xi(1530)$ component.

We showed that the calculated total decay width of the $\Omega^*(2012)$ is in fair agreement with the experimental data, thus, supporting the assignment of its spin parity as $3/2^-$.

In addition, we showed that the $\eta\Omega$ channel plays a less relevant role.

The present work should be viewed as a natural extension of the works of Refs. [12,13], where the chiral unitary approach was employed to dynamically generate such an exited $\Omega^*$. The present work showed indeed that the chiral unitary approach can provide a satisfactory explanation of not only the mass but also the decay width of the $\Omega(2012)$, consistent with the experimental data.

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