Strong two- and three-body decays of the new excited hyperon $\Omega^*(2012)$ are discussed in the hadronic molecular approach. The $\Omega^*(2012)$ state is considered to contain the mixed $\Xi^0 \bar{K}$ and $\Omega \eta$ hadronic components. In our calculations we use a phenomenological hadronic Lagrangian describing the coupling of the bound states to its constituents and of the constituents to other hadrons occurring in the final state. Our results show that the decay widths of the two-body decay modes $\Omega^*(2012) \to \Xi \bar{K}$ lie in the few MeV region and are compatible with or dominate over the rates of the three-body modes $\Omega^*(2012) \to \Xi \pi \bar{K}$. The sum of two- and three-body decay widths is consistent with a width of the $\Omega^*(2012)$ originally measured by the Belle Collaboration. A possible scenario for the suppression of the three-body decay rate recently noticed by the Belle Collaboration is due to the dominant admixture of the $\Omega \eta$ hadronic component in the $\Omega^*(2012)$ state.

I. INTRODUCTION

Last year the Belle Collaboration reported on a new excited isosinglet hyperon $\Omega^-$ state decaying into $\Xi^0 \bar{K}^-$ and $\Xi^- K^0_s$ pairs with a mass of $2012.4 \pm 0.7 \pm 0.6$ MeV and a width of $\Gamma = 6.4^{+2.8}_{-2.0} \text{(stat)} \pm 1.6 \text{(syst)}$ MeV [1]. The spin-parity quantum numbers of the $\Omega^-$ have been favored to be $J^P = \frac{3}{2}^-$ based on two arguments: (1) the observed mass value of the $\Omega^-$ is close to the theoretical predictions for the $\frac{3}{2}^-$ states, (2) the rather narrow width of the $\Omega^-$ decaying to a $\Xi \bar{K}$ pair via a $d$ wave. Recently, the Belle Collaboration searched for the cascade three-body decay $\Omega(2012) \to \Xi \Xi(1530) \to \Xi \pi \Xi$ [2]. They did not observe any significant signal in this channel and derived upper limits for the ratios of the branchings relative to the two-body $\Xi \Xi$ decay modes. In particular, the most stringent upper limits read [2]:

\[
R_1 = \frac{B(\Omega(2012) \to \Xi(1530)^0 \Xi^- \pi^0 K^-)}{B(\Omega(2012) \to \Xi^- K^-)} < 9.3\% ,
\]

\[
R_2 = \frac{B(\Omega(2012) \to \Xi(1530)^0 \Xi^- \pi^0 K^-)}{B(\Omega(2012) \to \Xi K^-)} < 7.8\% .
\]

An excited $\Omega^*$ state with $J^P = \frac{3}{2}^-$ and a mass of 2020 MeV has been predicted in a quark model with a QCD based potential in Ref. [3]. Later, different types of quark models [4], large $N_c$ approaches [5], algebraic string-like model [6], Skyrme model [7], and lattice QCD [8] reported on an estimate for the $\Omega^*$ mass in the region from 1953 to 2120 MeV. Inclusion of sizeable five-quark Fock components and their mixing with three-quark components in the constituent quark models, considered in Refs. [9], lead to a reduction of the $\Omega^*$ mass by about 200 MeV. The dynamical generation of the $\Omega^*$ state using two ($\Omega \eta$ and $\Xi \Xi$) coupled channels in the chiral unitary approach has been proposed and developed in Refs. [10]-[12], where the mass of the $\Omega^*$ has a strong dependence on the choice of model parameters leading to 2141 MeV in Ref. [11] and 1800 MeV in Ref. [12].

The understanding of the structure and decays of the $\Omega^*(2012)$ state has been of increased interest since the discovery by the Belle Collaboration. In Refs. [13]-[22] possible interpretations of the $\Omega^*$ (2012) hyperon have been critically discussed. In Ref. [13] the two-body decays $\Omega^* \to \Xi^0 K^-$ and $\Omega^* \to \Xi^- K^0_s$ have been analyzed in the chiral quark model. It was argued that the obtained numerical results are in agreement with spin-parity $\frac{3}{2}^-$ of the $\Omega^*$ state, while alternative assignments as $\frac{1}{2}^-$ and $\frac{1}{2}^+$ cannot be completely excluded. In Ref. [14] the possible structure and resulting strong decays of the $\Omega^*$ were studied using QCD sum rules. From the analysis of the mass and the strong decay properties (coupling constants and widths) it was concluded that the $\Omega^*$ hyperon is the $1P$ orbital excitation of the ground state $\Omega(1670)$ baryon with $J^P = \frac{3}{2}^-$. The same conclusion about the nature of the $\Omega^*$ state has been made from an analysis of its two-body decays in the framework of the $3P_0$ model. In Ref. [15] the $\Omega^*$ state has been considered on the basis of the $SU(3)$ flavor picture. It was found that if the $\Omega^*$ state is the $\Xi \Xi(1530)$ molecular state formed in the isospin zero channel, then its main decay mode is the tree-process mode $\Omega^* \to \Xi K \pi$. In Ref. [16] a hadronic molecular scenario for the $\Omega(2012)$ has been tested in an effective field-theoretical approach. It was found that the partial width of the three-body decay $\Omega^* \to \Xi K \pi$ is in the 2-3 MeV interval, while the partial width of the two-body decay $\Omega^* \to \Xi \bar{K}$ is in the 1-11 MeV range. Here and also in Refs. [17]-[19] the dominance or sizable contribution of the tree-body decays $\Omega^* \to \Xi K \pi$ has been based on the description of these processes by a tree-level diagram, while the two-body processes $\Omega^* \to \Xi \bar{K}$ have been...
described by a loop-diagram generated by the $\Omega^*$ constituents. In Ref. [21] the $\Omega$ baryon spectrum up to the $N = 2$ shell has been calculated based on a nonrelativistic constituent quark potential model. In Ref. [22] the analysis of the $\Omega(2012)$ strong decays has been performed using a hadronic molecular model and different spin-parity assignments for these states.

The main ideas in the application of quantum field theory to bound states using their compositeness were formulated in Refs. [23]-[26] and were widely applied to the study of hadronic molecules (HM) in Refs. [23] and [27]-[29]. In particular, in Ref. [23] the approach has been originally applied to the deuteron - the canonical example of a HM and in the latest years especially to the recently discovered exotic heavy hadrons [27]-[29] with unusual properties.

The main basic blocks and strategy of the quantum field approach for bound states [23]-[26] and specifically for HM [23], [27]-[29] are: (1) first, to derive a phenomenological Lagrangian, which is manifestly Lorenz covariant and gauge invariant, describing the interaction of the bound state with its constituents. The bound state and the constituents are described by standard local quantum field operators. The field operators of the constituents form the interpolating current with the corresponding quantum numbers of the bound state; (2) the coupling strength of a hadronic molecule to its constituents is determined by the composite-partner interaction Lagrangian of the HM with its constituents. The condition $Z_{\text{HM}} = 1 - \Pi_{\text{HM}}^* = 0$ [23]-[29]. $Z_{\text{HM}}$ is the wave function renormalization constant of the HM defining the matrix element (overlap) between physical and bare states of HM. $\Pi_{\text{HM}}^*$ is the derivative of the HM mass operator generated by the interaction Lagrangian of the HM with its constituents. The condition $Z_{\text{HM}} = 0$ means that the probability to find the HM as the bare state is always equal to zero or in other words, it is always dressed by its constituents. The compositeness condition provides an effective and selfconsistent way to describe the coupling of the HM to its constituents; (3) using an interaction Lagrangian between the HM and the constituents one can construct the $S$-matrix operator and consistently generate matrix elements for hadronic processes involving the HM (represented by corresponding Feynman diagrams). In the evaluation of the Feynman diagrams the compositeness condition enables to avoid the problem of double counting.

The main objective of the paper is to present a self-consistent study of strong two- and three-body decays of the $\Omega^*(2012)$ state in the hadronic molecular picture based on the formalism proposed and developed in Refs. [27]-[29]. In the recent paper by the Belle Collaboration [2] it is claimed that the strong three-body decay modes of the $\Omega^*(2012)$ state are suppressed in comparison with the two-body one, which raises doubts on a molecular interpretation of this state. We find that due to a possible mixture of the hadronic components $\Xi^* K$ and $\Omega\eta$ in the structure of $\Omega^*(2012)$ the three-body decays $\Omega(2012) \to \Xi\pi K$ are suppressed when the $\Omega\eta$ hadronic component dominates.

The paper is structured as follows. In Sec. II we present the details of our formalism for the treatment of the $\Omega^*(2012)$ state as a hadronic molecule. We consider the $\Omega^*(2012)$ hyperon as a mixed state of the hadronic pairs $\Xi^* K$ and $\Omega\eta$. In Sec. III we turn to the calculation of the strong two- and three-body decays of the $\Omega^*(2012)$ state: the processes $\Omega^* \to \Xi K$ and $\Omega^* \to \Xi\pi K$. We present a derivation of the corresponding matrix elements and discuss the numerical results.

II. HADRONIC MOLECULAR STRUCTURE OF THE $\Omega(2012)$

Following the conjecture of the Belle Collaboration [1] and assignments of most of the theoretical approaches we use the spin-parity quantum numbers $J^P = \frac{1}{2}^-$ for the $\Omega^*(2012)$ state. We consider the $\Omega^*(2012)$ as a weakly bound hadronic molecule, which involves a superposition of two hadronic components --- $\Omega(1670)\eta$ and $\Omega(1530)K$:

$$|\Omega^*(2012)) = \cos \theta \frac{|\Xi^0 K^-\rangle + |\Xi^- K^0\rangle}{\sqrt{2}} - \sin \theta |\Omega^* \eta\rangle,$$

where $\xi^* = (\Xi^0, \Xi^-)$ and $K = (K^-, \bar{K})$ are the doublets of $\Xi^*$ hyperons and $\bar{K}$ mesons, $\Phi(y^2)$ is a phenomenological correlation function describing the distribution of $(\Xi^* K)$ and $(\Omega \eta)$ constituents in the $\Omega^*$ state, $\omega_{\Xi^0} = M_{\Xi^0}/(M_{\Xi^0} + M_K)$, $\omega_{\Xi^-} = M_{\Xi^-}/(M_{\Xi^-} + M_K)$, $\omega_\eta = M_\eta/(M_{\Xi^0} + M_K)$, $\omega_\Pi = M_\Pi/(M_{\Xi^0} + M_K)$, and $\omega_\Sigma = M_\Sigma/(M_{\Xi^0} + M_K)$. To produce ultraviolet-finite Feynman diagrams, the Fourier transform of the correlation function $\Phi(y^2)$ should vanish sufficiently fast in the ultraviolet region of the Euclidean space. We use the Gaussian form for the correlation function $\hat{\Phi}(p_E^2) = \exp(-p_E^2/\Lambda_H^2)$, where $p_E$ is the Euclidean Jacobi momentum and $\Lambda_H$ is a free size parameter, which has a value of about 1 GeV.

The coupling $g_{\Omega^*}$ is determined from the compositeness condition (see Refs. [23]-[29])

$$Z_{\Omega^*} = 1 - \Sigma^T_{\Omega^*}(M_{\Omega^*}) = 0,$$
TABLE I: Quantum numbers and masses of relevant hadrons

| Hadron | I  | J^P  | Mass (MeV)  |
|--------|----|------|-------------|
| π^+   | 1  | 0^-  | 139.57061   |
| π^0   | 1  | 0^-  | 134.977     |
| K^0   | 1/2| 0^-  | 493.677     |
| K^-   | 1/2| 0^-  | 497.611     |
| η     | 0  | 0^-  | 547.862     |
| Κ^+  | 1/2| 1^-  | 891.76      |
| Κ^0  | 1/2| 1^-  | 895.55      |
| η^0   | 0  | 0^-  | 547.862     |

FIG. 1: Diagrams representing the mass operator of the \( Ω^* (2012) \) state.

where \( \Sigma_{T\Omega}^r \) is the derivative of the transverse part of the mass operator of the \( Ω^* \) state:

\[
\Sigma_{T\Omega}^{\mu\nu}(p) = g_{\perp T\Omega}^{\mu\nu}(p) + \frac{p^\mu p^\nu}{p^2} \Sigma_{\Omega\Omega}^L(p),
\]

with \( g_{\perp T\Omega}^{\mu\nu} = g^{\mu\nu} - p^\mu p^\nu / p^2 \). The corresponding Feynman diagrams contributing to the mass operator \( \Sigma_{\Omega\Omega} \), which are generated by the loops of the \( (Ξ^{-0} K^-) \), \( (Ξ^- K^0) \), and \( (Ω^- η) \) constituents, are shown in Fig. 1. Note that the compositeness condition gives a relation between the coupling constant \( g^{\mu\nu}_Ω \) and the mass \( m_{Ω^*} \). In addition we have a free parameter \( \theta \), which is the mixing angle between the \( Ξ^- K^- \) and the \( Ω^- η \) hadronic molecular components of the \( Ω^* \) state.

The expression for the mass operator \( \Sigma_{\Omega\Omega} \) is given by

\[
\Sigma_{\Omega\Omega}^{\mu\nu}(p) \propto \frac{\cos^2 \theta}{2} \left( \Sigma_{\Omega\Omega}(p) + \Sigma_{\Omega\Omega}(k^\mu k^\nu - k^2) \right) + \sin^2 \theta \Sigma_{\Omega\Omega}(p),
\]

\[
\Sigma_{\Omega\Omega}^{\mu\nu,H_1,H_2}(p) = \frac{g_{\perp H_1,H_2}^{\mu\nu}}{2} \int \frac{d^4 k}{(2\pi)^4} \tilde{G}^2(k + p - \omega_2) S_{H_1}^\mu(k) S_{H_2}^\nu(k),
\]

where

\[
S_{H_1}^\mu(k) = \frac{1}{M_{H_1}^2 - k^2} \left[ -g^{\mu\nu} + \frac{\gamma^{\mu\nu} k^\nu}{3} + \frac{2 k^\mu k^\nu}{3 M_{H_1}^2} + \frac{\gamma^{\mu\nu} k^\nu - \gamma^\nu k^\mu}{3 M_{H_1}} \right],
\]

\[
S_{H_2}(k) = \frac{1}{M_{H_2}^2 - k^2}
\]

are the propagators of baryon \( H_1 \) with spin \( \frac{3}{2} \) and of meson \( H_2 \) with spin 0. Here \( H_1 = Ξ^-, Ξ^0, Ω^- \) and \( H_2 = K^-, K^0, η \).
The coupling constant $g_{\Omega'}$ deduced from the compositeness condition \((9)\) reads
\[
\frac{g_{\Omega'}^2}{16\pi^2} = \frac{1}{I_{\Omega'}},
\]
where $I_{\Omega'}$ is the structure integral
\[
I_{\Omega'} = \frac{\cos^2 \theta}{2} \left( \bar{I}_{\Omega'}^{\omega K} + I_{\Omega'}^{\bar{K} \omega} \right) + \sin^2 \theta I_{\Omega'}^{\eta},
\]
\[
I_{\Omega'}^{\mu H} = \int_0^\infty \frac{d\omega_1 d\omega_2}{\Delta^3} R_{\Omega'} e^{-\omega_0}, \quad \Delta = 2 + \alpha_1 + \alpha_2
\]
and
\[
R_{\Omega'} = \left( 1 + \frac{1}{3\mu_\Delta^2} \right) \left( \alpha_2 + 2\omega_1 + 2\mu_{\Omega'} (\mu_{H_1} + \mu_{H_2} - \mu_\Omega) \right) \left( \alpha_2 + 2\omega_1 + 2\alpha_1 \omega_1^2 + 2\alpha_2 \omega_2^2 \right),
\]
\[
\omega_{\Omega'} = \frac{(\mu_{H_1} \alpha_1 - \mu_{H_2} \alpha_2)^2}{\Delta} + \left( (\mu_{H_1} + \mu_{H_2} - \mu_\Omega)^2 - \mu_\Omega^2 \right) \frac{\alpha_1 \alpha_2 + 2\alpha_1 \omega_1^2 + 2\alpha_2 \omega_2^2}{\Delta}, \quad \mu_i = \frac{M_i}{\Lambda_{\Omega'}}.
\]

The leading diagrams contributing to the strong decays of the $\Omega'$ state are shown in Fig. 2 (a) two-body decay $\Omega' \to \Xi \bar{K}$, (b) three-body decay $\Omega' \to \Xi \pi \bar{K}$. Note that in agreement with Refs. [13-19] the two-body decay $\Omega' \to \Xi \bar{K}$ proceeds via the hadronic loops ($\Xi' \bar{K}$) and ($\Xi' \eta$) involving the constituents of the $\Omega'$ state. The three-body decay $\Omega' \to \Xi \pi \bar{K}$ is described by the two-cascade tree-level diagram, where the $\Omega'$ first couples to the constituents $\Xi'$ and $\bar{K}$ and then $\Xi'$ decays via the dominant mode into $\Xi$ and $\pi$. To evaluate the diagrams in Fig. 2 we need to set the interaction Lagrangian, which includes the coupling of the $\Omega'$ with the constituents $\Xi' \bar{K}$ and $\Omega' \eta$, and two additional terms describing the couplings $\Xi' \Xi K K$ ($\Xi' \Omega' K$) and $\Xi' \Xi \pi$.

![FIG. 2: Leading diagrams contributing to the strong decays of the $\Omega'$ state: (a) two-body decay $\Omega' \to \Xi \bar{K}$, (b) three-body decay $\Omega' \to \Xi \pi \bar{K}$.

The $\Xi' \Xi \pi$ interaction is described by the phenomenological Lagrangian
\[
\mathcal{L}_{\Xi' \Xi \pi} = \frac{g_{\Xi' \Xi \pi}}{M_{\Xi'}} \bar{\Xi'} \gamma^\mu \pi \Xi + \text{H.c.},
\]
where the dimensionless coupling $g_{\Xi' \Xi \pi}$ is fixed from data on the $\Xi' \to \Xi \pi$ decay width. In particular, the corresponding two-body decay width reads
\[
\Gamma(\Xi' \to \Xi \pi) = \frac{g_{\Xi' \Xi \pi}^2}{64 M_{\Xi'}^2} 2 \lambda^{3/2} (M_{\Xi'}^2, M_{\Xi'}^2, M_{\Xi}^2) \left( (M_{\Xi'} + M_{\Xi})^2 - M_{\pi}^2 \right).
\]

Using the measured central values of
\[
\Gamma_{\text{total}}(\Xi^0) = 2 \Gamma(\Xi^0 \to \Xi^- \pi^+) + \Gamma(\Xi^0 \to \Xi^0 \pi^0) = 9.14 \text{ MeV}, \quad \Gamma_{\text{total}}(\Xi^+) = 2 \Gamma(\Xi^+ \to \Xi^0 \pi^+) + \Gamma(\Xi^+ \to \Xi^+ \pi^0) = 9.9 \text{ MeV}
\]
one gets $g_{\Xi' \Xi \pi} = 6.79$ from the $\Xi^0$ set and $g_{\Xi' \Xi \pi} = 6.71$ from $\Xi^+$. In the following we will use the averaged value $g_{\Xi' \Xi \pi} = (6.79 + 6.71)/2 = 6.75$.

The couplings $\Xi' \Xi K K$ and $\Xi' \Omega' K$ are generated by a phenomenological Lagrangian:
\[
\mathcal{L}_{\Xi' DB} = g_{\Xi' DB} \gamma^\mu \pi^{ij} \Omega^{\mu j} \epsilon^{ilm} + \text{H.c.},
\]
where $B_{mk}$ and $D_{ijk}$ are the octet and decuplet baryon fields. $\Gamma_\mu$ is the chiral connection which in absence of external vector and axial fields is defined as

$$\Gamma_\mu = \frac{1}{2} [u^\dagger, \partial_\mu u] = \frac{1}{4F_2^2}[\hat{\Phi}, \partial_\mu \hat{\Phi}] + O(\hat{\Phi}^4),$$

where $\hat{\Phi} = \sum_{i=1}^{8} \phi_i \lambda_i$ is the octet matrix of pseudoscalar mesons. The dimensionless coupling $g_{TDB}$ can be expressed through the

FIG. 3: Two-body decay widths for $\Omega^* \rightarrow \Xi K$ at $\Lambda = 1$ GeV (left panel) and $\Lambda = 1.5$ GeV (right panel).

FIG. 4: Three-body decay widths for $\Omega^* \rightarrow \Xi\pi K$ at $\Lambda = 1$ GeV (left panel) and $\Lambda = 1.5$ GeV (right panel).

FIG. 5: Ratios $R_1$ and $R_2$ in comparison with the Belle upper limits of Eq. (1) at $\Lambda = 1$ GeV (left panel) and $\Lambda = 1.5$ GeV (right panel).
fundamental constants of hadron physics: $\pi NN$ coupling $g_{\pi NN} \approx 13.4$, $\pi N\Delta$ coupling $f_{\pi N\Delta} \approx 2$, pion mass $M_\pi = 0.13957$ GeV, and nucleon mass $M_N = 0.93827$ GeV:

$$g_{\Gamma DB} = \frac{f_{\pi N\Delta}}{g_{\pi NN}} \frac{M_N}{M_\pi} \approx 1. \quad (17)$$

From the Lagrangian (15) we need the leading-order term in the chiral expansion, i.e. the term quadratic in the pseudoscalar coupling $g$. The invariant Mandelstam variables $s$ are correlated with the assumption that the $\Omega^*$ two-body decay rates of the hadrons -- $\Omega^*$, $\eta$, and $\pi$ -- are related to the momenta $p_1$, $p_2$, and $p_3$ of the hadrons, with $\Omega^*$ being the hadronic molecular structure. The prediction given here can hopefully support a possible structure interpretation of the $\Omega^*(2012)$.

The invariant Mandelstam variables $s_1 = (p_1 + p_2)^2 = (p - p_1)^2$ and $s_2 = (p_2 + p_3)^2 = (p - p_1)^2$ are related to the momenta $p_1$, $p_2$, and $p_3$ of the hadrons -- $\Omega^*$, $K$, $\Xi$, and $\pi$, respectively. The variable $s_1$ has the upper/lower limits $s_1^+$ with:

$$s_1^+ = M_k^2 + M_\Xi^2 + \frac{1}{2s}(s - s_2 - M_\Xi^2)(s_2 + M_\Xi^2 - M_k^2) \pm \frac{1}{2s_2} \lambda^{1/2}(s, s_2, M_k^2) \lambda^{1/2}(s_2, M_\Xi^2, M_k^2). \quad (20)$$

where $\lambda(x,y,z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ is the Källen kinematical triangle function. The $\Omega^* \Xi K$ coupling $g_{\Omega^* \Xi K}$ is evaluated from the diagram in Fig. 2.

In Figs. 3 and 4, we show our predictions for the respective charged combinations of two- and three-body decay widths of the $\Omega^*(2012)$ and their sums. Results are indicated for the values of $\Lambda = 1$ and 1.5 GeV, respectively, and for the mixing angle $\theta$ varied from 0 to 90 grad. In Fig. 5, we plot our predictions for the ratios $R_1$ and $R_2$ of Eq. (11) of the three- and two-body decays and compare them with the upper limits derived by the Belle Collaboration in Ref. [2]. One can see that an increase in the admixture of the hadronic component $\Omega R$ leads to a suppression of the three-body decay rates. For the mixing angle bigger 70° (73°) for $\Lambda = 1$ GeV and bigger 55° (61°) for $\Lambda = 1.5$ GeV our ratios are consistent with Belle results for $R_1$ and $R_2$. To get a handle on the parameters $\theta$ and $\Lambda$ further data on the $\Omega^*(2012)$ decays are needed. The predictions presented here are strongly correlated with the assumption that the $\Omega^*(2012)$ has a hadronic molecular structure.

Finally, we discuss our main numerical results. Varying the model scale parameter $\Lambda$ from 1 to 1.5 GeV we find that the magnitude of the two-body decay rates slightly increase, while the three-body rates decrease. The relative contribution of two to three-body decays is governed by the mixing angle $\theta$ -- mixing of the $\Xi K$ and $\Omega R$ hadronic components. An increase of $\theta$ leads to a suppression of the three-body decay rates in comparison with the two-body ones. For the total two- and three-body decay rates we obtain the following numerical results:

$$\Gamma(\Omega^* \to \Xi K) = 2.9 \pm 1.6 \text{ MeV}, \quad \Gamma(\Omega^* \to \Xi \pi K) = 1.5 \pm 1.5 \text{ MeV} \quad (21)$$

for $\Lambda = 1$ GeV and $\theta \in [0, \pi/2]$.

$$\Gamma(\Omega^* \to \Xi K) = 6 \pm 3.4 \text{ MeV}, \quad \Gamma(\Omega^* \to \Xi \pi K) = 1.4 \pm 1.4 \text{ MeV} \quad (22)$$

for $\Lambda = 1.5$ GeV and $\theta \in [0, \pi/2]$. The upper limits of Eq. (11) set by the Belle Collaboration [2] for the ratios of three- and two-body decay rates of the $\Omega(2012)$ are fulfilled for following values of the $\theta$ angle:

1) $\Lambda = 1$ GeV

$$R_1 < 9.3\% \quad \text{at } \theta \geq 70^0, \quad R_2 < 7.8\% \quad \text{at } \theta \geq 73^0. \quad (23)$$

2) $\Lambda = 1.5$ GeV

$$R_1 < 9.3\% \quad \text{at } \theta \geq 55^0, \quad R_2 < 7.8\% \quad \text{at } \theta \geq 61^0. \quad (24)$$

Hence a sizable $\Omega R$ hadronic component in the $\Omega^*(2012)$ leads to a suppression of the $\Xi \pi K$ mode relative to the $\Xi K$ channel. To get a further handle on the parameters $\theta$ and $\Lambda$ in the context of the hadronic structure precise data on the $\Omega^*(2012)$ decays are needed. The prediction given here can hopefully support a possible structure interpretation of the $\Omega^*(2012)$. 
