

Final state interaction in the $\Lambda_b \rightarrow D^+ D^- \Lambda$, $D^0 D_s^- p$, $D_s^+ D_s^- \Lambda$ reactions

Jing Song,^{1,2,3,*} Melahat Bayar^{4,3,†} and Eulogio Oset^{3,5,‡}

¹*Center for Theoretical Physics, School of Physics and Optoelectronic Engineering, Hainan University, Haikou 570228, China*

²*School of Physics, Beihang University, Beijing, 102206, China*

³*Departamento de Física Teórica and IFIC,*

Centro Mixto Universidad de Valencia-CSIC Institutos de Investigación de Paterna, 46071 Valencia, Spain

⁴*Department of Physics, Kocaeli University, 41380, Izmit, Turkey*

⁵*Department of Physics, Guangxi Normal University, Guilin 541004, China*



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We study the three-body decays $\Lambda_b \rightarrow D^+ D^- \Lambda$, $D^0 D_s^- p$, $D_s^+ D_s^- \Lambda$, considering both internal and external emission mechanisms that produce various meson-baryon final states. Using a unitarized coupled-channel approach based on the local hidden gauge formalism, we analyze final state interactions in the exotic anticharm systems with strangeness $S = -1$ and $S = -2$. Significant threshold enhancements appear in both pseudoscalar and vector meson-baryon channels, indicating strong final state interactions and the possible formation of molecular exotic states below threshold. These features remain stable under variations of model parameters. Our results provide valuable insights into heavy hadron decay dynamics and offer theoretical guidance for future experimental searches of open-charm strange exotic hadrons. These results take into account the effects of the $D^+ \Lambda$ and $D^+ D^-$ interactions, three-body rescattering, possible triangle singularities, parameter uncertainties, and large N_c behavior, making the predictions very reliable.

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I. INTRODUCTION

The discovery of exotic meson and baryon states, which challenge the traditional classification of mesons as $q\bar{q}$ and baryons as three-quark (qqq) states, has significantly advanced the field of hadron physics, particularly hadron spectroscopy. This development has motivated extensive theoretical and experimental efforts, and numerous review articles have been devoted to this topic [1–14].

The study of weak decays of heavy baryons offers valuable insights into the internal structure of hadrons and the dynamics of hadron-hadron interactions [15–27] (see review paper on this issue in Ref. [10]). In particular, the three-body decay $\Lambda_b \rightarrow D^+ D^- \Lambda$ provides an ideal platform to investigate meson-baryon interactions involving open charm mesons and strange baryons. The LHCb Collaboration has recently reported the measurement of this decay [28], making available the invariant mass distributions of the $D^+ \Lambda$, $D^- \Lambda$, and $D^+ D^-$ subsystems. These

distributions offer rich opportunities to search for exotic hadronic states dynamically generated through final state interactions (FSI). Related to the $\Lambda_b \rightarrow D^+ D^- \Lambda$ reaction one would also have $\Lambda_b \rightarrow D^0 D_s^- p$, $D_s^+ D_s^- \Lambda$. The final states in these decays contain subsystems with strangeness $S = -1$ and $S = -2$, including $D^- \Lambda$, $D_s^- p$, and $D_s^- \Lambda$. The quark content of the $D^- \Lambda$ system is $\bar{c}dsud$, which does not contain any $q\bar{q}$ couple that could annihilate, and hence is a clear exotic state which cannot be accommodated with a $3q$ system. A recent theoretical study [29], based on the interaction between heavy anticharmed mesons and octet baryons, predicts a bound state around 2888 MeV. This result is obtained using the local hidden gauge approach with a commonly adopted cutoff range of $q_{\text{max}} = 550\text{--}650$ MeV [30–32]. The predicted state lies about 20 MeV below the $D_s^- p$ threshold and about 92 MeV below the $D^- \Lambda$ threshold (strangeness $S = -1$, isospin $I = 1/2$). Bound states of these systems, using the same coupled channels, are also obtained in [33], although less bound. A discussion on these differences is given in [29]. These findings suggest that the $D^- \Lambda$ and $D_s^- p$ interactions may produce threshold enhancements or cusplike structures in their respective invariant mass distributions, which could signal the formation of near-threshold molecular states.

The theoretical framework employed here is based on the local hidden gauge approach extended to heavy flavors [34–37]. In this model, the interaction between mesons and

*Contact author: Song-Jing@buaa.edu.cn

†Contact author: melahat.bayar@kocaeli.edu.tr

‡Contact author: oset@ific.uv.es

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baryons is driven by vector meson exchange. The amplitudes for the coupled channels are unitarized by solving the Bethe-Salpeter equation [30,38–50], allowing for the generation of molecular states.

Among many others resonances, this framework has successfully explained the nature of the $D_{s0}^*(2317)$, which is widely considered a DK hadronic molecule [51–57]. Its low mass and narrow width, inconsistent with quark model expectations [58,59], are naturally described within this approach. Furthermore, the formation of hadronic molecules in weak decays has been observed or suggested in several reactions [60,61] (see review paper [10]).

In the $\Lambda_b \rightarrow D^+ D^- \Lambda$ decay, the weak transition $b \rightarrow c \bar{c} s$ can proceed through different mechanisms, including internal and external emission [62]. The hadronization of the resulting quark pairs produces various meson-baryon combinations. Internal emission processes often generate meson-meson pairs, while external emission tends to favor meson-baryon configurations. Both types of decay topologies are considered in our analysis, as they lead to different initial meson-baryon combinations and influence the strength of FSI.

The primary objective of this work is to analyze the mass distributions of the $D^- \Lambda$, $D_s^- p$, and $D_s^- \Lambda$ systems in the Λ_b decay and search for signs of near-threshold exotic structures. The results are compared with phase space distributions to highlight possible enhancements caused by strong interactions. These findings may serve as useful input for future experimental measurements. Besides the pseudoscalar meson-baryon systems, this work also explores vector meson-baryon interactions, including $D^{*-} \Lambda$, $D_s^{*-} p$, and $D_s^{*-} \Lambda$. These vector channels, show threshold effects that complement those found in pseudoscalar cases and contribute additional insight into the formation of exotic molecular states. Considering both pseudoscalar and vector mesons allows for a more comprehensive understanding of the heavy meson-baryon dynamics and the possible spectrum of exotic hadrons.

We also make a thorough discussion of possible uncertainties, coming from the consideration of the $D^- \Lambda$ and $D^+ D^-$ interaction, three body interaction effects, and possible contributions from triangle singularities. We find these effects negligible. Finally we also conduct an evaluation of the uncertainties of the results from cutoff and N_c uncertainties, finding that they are very small in the $D_s^- p$ mass distributions close to threshold normalized to same area.

In summary, we present a comprehensive study of the $\Lambda_b \rightarrow D^+ D^- \Lambda$, $D^0 D_s^- p$, $D_s^+ D_s^- \Lambda$ decays, focusing on the role of final state interactions and their potential to generate exotic hadronic molecules. By incorporating both internal and external emission mechanisms and analyzing all relevant meson-baryon combinations using a unitarized coupled-channel approach, we aim at shedding light on the structure of possible near-threshold states and the dynamics of heavy baryon decays.

II. FORMALISM

In this section, we present the theoretical framework to study the Λ_b decay, emphasizing the role of FSI that can dynamically generate hadronic molecular states. The weak decay of Λ_b is initiated by the weak transition of the b quark via external and internal W^- emission mechanisms, both of which allow the formation of meson-baryon pairs in the final state through hadronization. These pairs, through rescattering, can undergo strong interaction leading finally to the $D^- \Lambda$, $D_s^- p$ and $D_s^- \Lambda$ states.

The strong interaction between the final meson-baryon pairs is treated within a coupled-channel unitary approach, allowing for the dynamical generation of hadronic molecular states. Such interactions can lead to enhancements or resonancelike structures in the different invariant mass distributions. In the following, we detail the weak decay processes, identify the relevant hadronic channels, and discuss how the final state rescattering leads to dynamically generated resonances.

A. Weak decay mechanisms and hadronization: External and internal W^- emissions

In Fig. 1 we show the mechanisms of external and internal emission, which after hadronization will produce two mesons and one baryon for the processes that we are interested in. In the internal W^- emission mechanism Fig. 1(b), the b quark undergoes a weak transition via the exchange of a W^- boson producing $c \bar{c} s$ quarks. The resulting $c \bar{c}$ quark configuration hadronizes into two mesons. Although this Cabibbo-allowed process is color-suppressed, it can play an important role in shaping the decay amplitude through interference and final-state interaction effects.

On the other hand, the external emission mechanism, Fig. 1(a), proceeds with the b quark decaying weakly into a c quark and a W^- boson, which subsequently produces a $\bar{c} s$ pair, forming a D_s^- meson. The remaining $c(ud - du)$ quarks undergo hadronization, involving the c quark, thus producing a D, D_s meson and three extra quarks which combine to some physical baryon. This process is Cabibbo-allowed and color-favored, and generally dominates the overall decay rate.

Let us start from the mechanism of Fig. 1(a) for external emission. We produce directly a D_s^- and we hadronize the cu quark to produce an extra meson baryon state. We have

$$\Lambda_b = \frac{1}{\sqrt{2}} b(ud - du) \chi_{MA}, \quad (1)$$

where χ_{MA} is the mixed antisymmetric spin wave function on the last two quarks. Through weak decay the Λ_b becomes

$$\Lambda_b \longrightarrow D_s^- \frac{1}{\sqrt{2}} c(ud - du) \chi_{MA}. \quad (2)$$

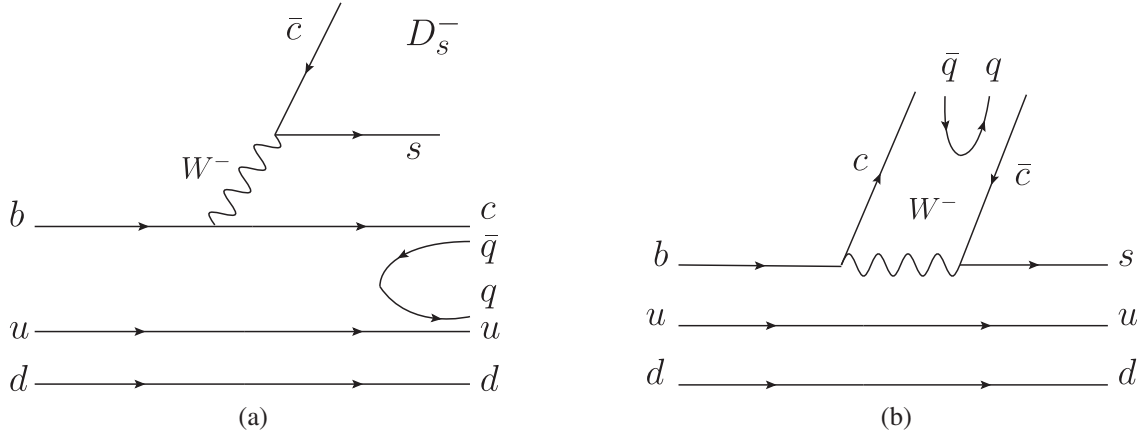


FIG. 1. The quark pairs generated through both internal and external W^- emission undergo hadronization to form the final hadrons: (a) external emission, (b) internal emission.

Next we hadronize the c quark and one of the light quarks as

$$\frac{1}{\sqrt{2}} c(ud - du) \longrightarrow \sum_i \frac{1}{\sqrt{2}} c \bar{q}_i q_i (ud - du) \longrightarrow \sum_i \frac{1}{\sqrt{2}} P_{4i} q_i (ud - du), \quad (3)$$

where P is the $q_i \bar{q}_j$ matrix written in terms of pseudoscalar mesons

$$P = \begin{pmatrix} \frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{3}} \eta + \frac{1}{\sqrt{6}} \eta' & \pi^+ & K^+ & \bar{D}^0 \\ \pi^- & -\frac{1}{\sqrt{2}} \pi^0 + \frac{1}{\sqrt{3}} \eta + \frac{1}{\sqrt{6}} \eta' & K^0 & D^- \\ K^- & \bar{K}^0 & -\frac{1}{\sqrt{3}} \eta + \sqrt{\frac{2}{3}} \eta' & D_s^- \\ D^0 & D^+ & D_s^+ & \eta_c \end{pmatrix}, \quad (4)$$

which implies the η, η' mixing of Ref. [63]. Ignoring η_c which does not play a role here, we have

$$\Lambda_b \longrightarrow D_s^- \frac{1}{\sqrt{2}} \{D^0 u(ud - du) + D^+ d(ud - du) + D_s^+ s(ud - du)\} \chi_{MA}. \quad (5)$$

We have to see which baryons come from the contribution of the 3 quarks. Recall that the baryons of the SU(3) octet are given by

$$\phi = \frac{1}{\sqrt{2}} (\phi_{MS} \chi_{MS} + \phi_{MA} \chi_{MA}). \quad (6)$$

The overlap of ϕ with the three quark obtained in Eq. (5) is

$$\left\langle \phi \left| \begin{Bmatrix} u \\ d \\ s \end{Bmatrix} (ud - du) \chi_{MA} \right\rangle = \frac{1}{\sqrt{2}} \left\langle \phi_{MA} \left| \begin{Bmatrix} u \\ d \\ s \end{Bmatrix} (ud - du) \right\rangle. \quad (7)$$

And, consistently with the structure of the Lagrangians used to describe the meson-baryon interaction, the ϕ_{MA} wave functions are given in [23]¹ as

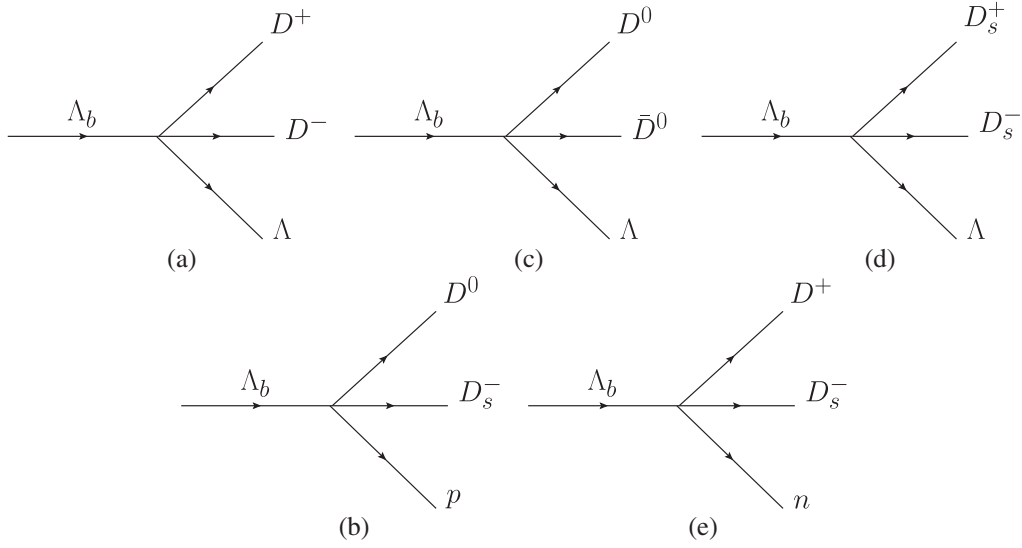
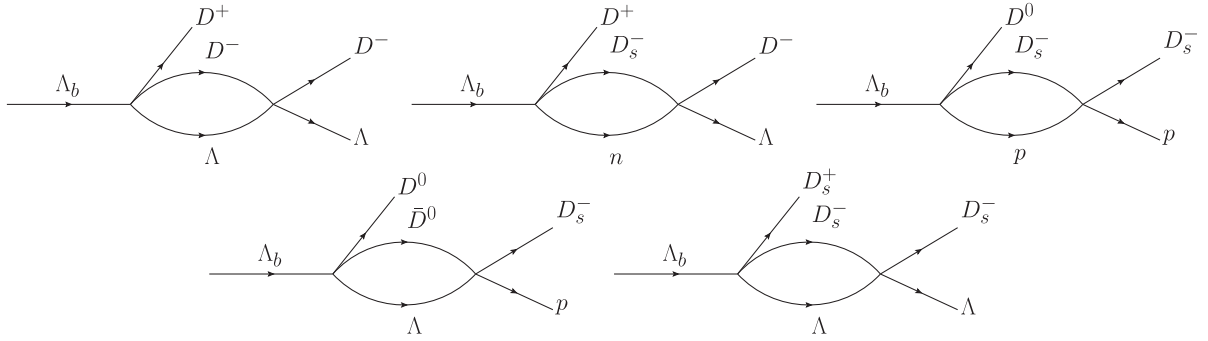
$$\begin{aligned} p &= \frac{1}{\sqrt{2}} u(ud - du), \\ n &= \frac{1}{\sqrt{2}} d(ud - du), \\ \Lambda &= \frac{1}{2\sqrt{3}} [u(ds - sd) + d(su - us) - 2s(ud - du)], \end{aligned} \quad (8)$$

Hence,

$$\Lambda_b \longrightarrow \frac{1}{\sqrt{2}} D_s^- D^0 p + \frac{1}{\sqrt{2}} D_s^- D^+ n - \frac{1}{\sqrt{3}} D_s^- D_s^+ \Lambda. \quad (9)$$

For the internal emission of Fig. 1(b), we have the hadronization of the $c\bar{c}$ pair, which proceeds as

¹The phase of the Λ, Ξ^0, Σ^+ are opposite to those of Ref. [64].

FIG. 2. Hadronic configurations from hadronization in the Λ_b decay direct production.FIG. 3. Hadronic configurations from hadronization through rescattering in the Λ_b decay.

$$\begin{aligned}
 c\bar{c} &\longrightarrow \sum_i c\bar{q}_i q_i \bar{c} \equiv \sum_i P_{4i} P_{i4} = (P^2)_{44} \\
 &= D^0 \bar{D}^0 + D^+ D^- + D_s^+ D_s^-. \quad (10)
 \end{aligned}$$

Hence

$$\begin{aligned}
 \Lambda_b &\longrightarrow \frac{1}{\sqrt{2}} (D^0 \bar{D}^0 + D^+ D^- + D_s^+ D_s^-) s (ud - du) \chi_{MA} \\
 &= -\frac{1}{\sqrt{3}} (D^0 \bar{D}^0 + D^+ D^- + D_s^+ D_s^-) \Lambda. \quad (11)
 \end{aligned}$$

From both mechanisms, we can expect the production of meson-baryon pairs such as $D^-\Lambda$, $D_s^- p$, and $D_s^- \Lambda$, either directly or through rescattering. These pairs are subject to strong interaction and can form molecular states or cusplike enhancements. The possible hadronic configurations are shown in Fig. 2.

However, we can also obtain the same final states through rescattering, as shown in Fig. 3.

We next study in detail the interactions among these configurations, starting from the most relevant one: the $D^-\Lambda$ system.

B. The $D^-\Lambda$, $D_s^- p$ coupled channel, and $D_s^- \Lambda$ state in the coupled channels unitary approach

In the final state, the $D^-\Lambda$ pair can undergo strong interactions, which can be studied using the local hidden gauge approach extended to the heavy quark sector [29,33,65,66]. The $D^-\Lambda$ system couples to the channels $D_s^- p$ and $\bar{D}\Sigma$ ($I = 1/2$).

We consider the coupled-channel meson-octet baryon interaction including coupled-channel dynamics corresponding to $|MB\rangle = \{\bar{D}_s N, \bar{D}\Lambda, \bar{D}\Sigma\}$ and $|MB\rangle = \bar{D}_s \Lambda, \bar{D}\Xi$ in $S = -1, I = 1/2$ and $S = -2, I = 0$ sectors [29], respectively, within the framework of the local hidden gauge approach [34–37]. The s -wave interaction kernel is obtained from vector-meson exchange, and the scattering amplitudes are computed by solving the Bethe-Salpeter equation in coupled channels,

$$T = [1 - VG]^{-1}V, \quad (12)$$

where V is the interaction kernel,

$$V_{ij} = C_{ij} \frac{1}{4f_\pi^2} (k^0 + k'^0). \quad (13)$$

Here, k^0 and k'^0 represent the meson energies in the initial and final states, respectively, and are given by

$$k^0 = \frac{s + m_{m_i}^2 - M_{B_i}^2}{2\sqrt{s}}, \quad k'^0 = \frac{s + m_{m_j}^2 - M_{B_j}^2}{2\sqrt{s}}, \quad (14)$$

where m_{m_i}, M_{B_i} and m_{m_j}, M_{B_j} denote the meson and baryon masses in the initial and final channels, respectively. The interaction strength between channels is characterized by the coefficients C_{ij} , which are listed in Tables I and II of Ref. [29]. The meson-baryon loop function $G_i(\sqrt{s})$ is evaluated using a cutoff regularization in momentum space

$$G_i(\sqrt{s}) = 2M_i \int_{|\vec{q}| < q_{\max}} \frac{d^3q}{(2\pi)^3} \frac{\omega_1(\vec{q}) + \omega_2(\vec{q})}{2\omega_1(\vec{q})\omega_2(\vec{q})} \times \frac{1}{s - [\omega_1(\vec{q}) + \omega_2(\vec{q})]^2 + i\epsilon}, \quad (15)$$

where the on-shell energies of the meson and baryon in the loop are given by

$$\omega_1(\vec{q}) = \sqrt{\vec{q}^2 + m_i^2}, \quad \omega_2(\vec{q}) = \sqrt{\vec{q}^2 + M_i^2}.$$

We adopt a cutoff range of $q_{\max} = 550\text{--}650$ MeV, in accordance with Refs. [30–32]. Poles of the T matrix in the complex energy plane correspond to dynamically generated resonances.

C. $D^- \Lambda$ mass distribution in the $\Lambda_b \rightarrow D^+ D^- \Lambda$ decay

Within the theoretical framework, the analysis of the $D^- \Lambda$ invariant mass distribution in the decay $\Lambda_b \rightarrow D^+ D^- \Lambda$ provides a valuable approach to investigate the dynamics of the $D^- \Lambda$ interaction. In particular, this observable is sensitive to possible resonant structures dynamically generated through final-state interaction. Such structures may manifest as enhancements or peaks

near the $D^- \Lambda$ threshold, which can be interpreted as signals of strong interaction effects, including cusp phenomena or the formation of hadronic molecular states. Notably, the $D^- \Lambda$ pair originates directly from the internal W^- emission mechanism [see Eq. (11)]. However, it can also come from external emission and rescattering. An enhancement in this region would thus reflect the underlying interaction strength and the existence of a near-threshold molecular configuration.

The differential decay width can be written as

$$\frac{d\Gamma}{dM_{\text{inv}}(D^- \Lambda)} = \frac{1}{(2\pi)^3} \frac{1}{4M_{\Lambda_b}^2} p_{D^+} \tilde{p}_{D^-} |\mathcal{M}|^2 \quad (16)$$

with

$$p_{D^+} = \frac{\lambda^{1/2}(M_{\Lambda_b}^2, m_{D^+}^2, M_{\text{inv}}^2(D^- \Lambda))}{2M_{\Lambda_b}}, \quad \tilde{p}_{D^-} = \frac{\lambda^{1/2}(M_{\text{inv}}^2(D^- \Lambda), m_{D^-}^2, m_{\Lambda}^2)}{2M_{\text{inv}}(D^- \Lambda)}, \quad (17)$$

where $\mathcal{M}$ denotes the full decay amplitude, which incorporates both the direct weak production mechanism and the contribution from the final-state interaction, as illustrated in Fig. 4.

The T matrix is evaluated using the formalism described in the previous subsection.

$$\begin{aligned} \mathcal{M} = & -\frac{1}{\sqrt{3}} \mathcal{C} - \frac{1}{\sqrt{3}} \mathcal{C} G_{D^- \Lambda}(M_{\text{inv}}(D^- \Lambda)) t_{D^- \Lambda, D^- \Lambda} \\ & \times (M_{\text{inv}}(D^- \Lambda)) \\ & + \frac{1}{\sqrt{2}} \mathcal{C} N_c G_{D_s^- n}(M_{\text{inv}}(D^- \Lambda)) t_{D_s^- n, D^- \Lambda}(M_{\text{inv}}(D^- \Lambda)), \end{aligned} \quad (18)$$

where $\mathcal{C}$ is a normalization constant that includes the common factors of the weak and hadronization amplitudes. The function $G_{D^- \Lambda}(M_{\text{inv}})$ is the loop function for the intermediate $D^- \Lambda$ channel, and $t_{D^- \Lambda, D^- \Lambda}(M_{\text{inv}})$ is the corresponding unitarized scattering amplitude evaluated at the invariant mass $M_{\text{inv}}(D^- \Lambda)$. The second term accounts for the FSI of the $D^- \Lambda$ system. The third term represents the contribution from the coupled channel

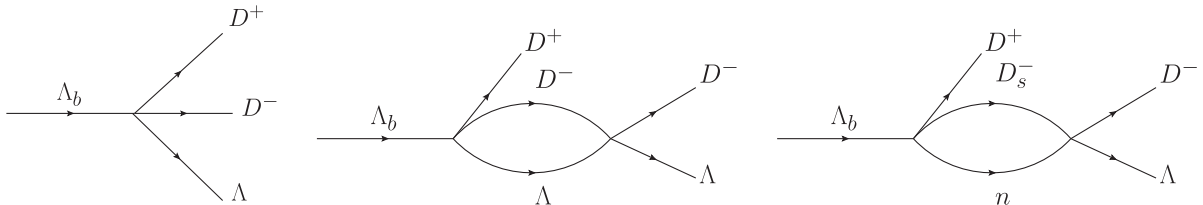
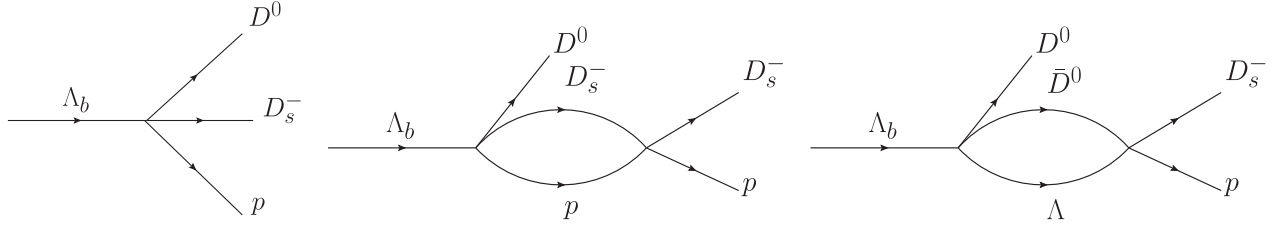


FIG. 4. Diagrammatic representation of the $D^- \Lambda$ invariant mass distribution in Λ_b decay.

FIG. 5. Diagrammatic representation of the $D_s^- p$ contribution in Λ_b decay, including both direct production and rescattering effects.

$D_s^- n \rightarrow D^- \Lambda$, where $G_{D_s^- n}$ is the loop function and $t_{D_s^- n, D^- \Lambda}$ is the transition amplitude between the two channels. The factor $N_c = 3$ is implemented to take into account the color enhancement of external emission versus internal emission. Note that while the first two diagrams in Fig. 4 involve the weak internal emission mechanism, the third diagram [third term in Eq. (18)] involves weak external emission. Hence, this term should have a large weight compared to the first two terms, enhancing the role of the coupled channels interaction.

D. $D_s^- p$ mass distribution in the $\Lambda_b \rightarrow D^0 D_s^- p$ decay

Next we consider a related reaction where the $D_s^- p$ channel is produced in the final state. As seen in Eq. (9) the $D^0 D_s^- p$ channel appears in the Λ_b decay from external emission, but it can also be reached via rescattering from a first step production of $D^0 \bar{D}^0 \Lambda$ in internal emission. The mechanisms for $\Lambda_b \rightarrow D^0 D_s^- p$ decay are depicted in Fig. 5.

The differential decay width with the $D_s^- p$ channel included can be expressed as

$$\frac{d\Gamma}{dM_{\text{inv}}(D_s^- p)} = \frac{1}{(2\pi)^3} \frac{1}{4M_{\Lambda_b}^2} p_{D^0} \tilde{p}_{D_s^-} |\mathcal{M}|^2, \quad (19)$$

with

$$p_{D^0} = \frac{\lambda^{1/2}(M_{\Lambda_b}^2, m_{D^0}^2, M_{\text{inv}}^2(D_s^- p))}{2M_{\Lambda_b}},$$

$$\tilde{p}_{D_s^-} = \frac{\lambda^{1/2}(M_{\text{inv}}^2(D_s^- p), m_{D_s^-}^2, m_p^2)}{2M_{\text{inv}}(D_s^- p)}, \quad (20)$$

where the full decay amplitude $\mathcal{M}$ now contains both the primary weak production amplitude and contributions from the final-state interactions involving the $D_s^- p$ intermediate state. Hence, the corresponding amplitudes $\mathcal{M}$ is written as,

$$\begin{aligned} \mathcal{M} = & \frac{1}{\sqrt{2}} C N_c + \frac{1}{\sqrt{2}} C N_c G_{D_s^- p}(M_{\text{inv}}(D_s^- p)) t_{D_s^- p, D_s^- p} \\ & \times (M_{\text{inv}}(D_s^- p)) \\ & - \frac{1}{\sqrt{3}} C G_{\bar{D}^0 \Lambda}(M_{\text{inv}}(D_s^- p)) t_{\bar{D}^0 \Lambda, D_s^- p}(M_{\text{inv}}(D_s^- p)), \end{aligned} \quad (21)$$

where the first term represents the tree-level contribution to the $D_s^- p$ channel, with N_c being the corresponding color enhancement factor. The second term accounts for the final state interaction of the $D_s^- p$ channel, involving the loop function $G_{D_s^- p}$ and the scattering amplitude $t_{D_s^- p, D_s^- p}$. The third term includes the effect of the coupled-channel transition from $\bar{D}^0 \Lambda$ to $D_s^- p$, with the corresponding loop function $G_{\bar{D}^0 \Lambda}$ and transition amplitude $t_{\bar{D}^0 \Lambda, D_s^- p}$.

E. $D_s^- \Lambda$ mass distribution in the $\Lambda_b \rightarrow D_s^+ D_s^- \Lambda$ decay

The study of the $D_s^- \Lambda$ invariant mass distribution in the $\Lambda_b \rightarrow D_s^+ D_s^- \Lambda$ reaction provides an opportunity to explore coupled-channel dynamics in a more exotic sector. As in previous cases, the effect of final-state interactions is implemented through the corresponding scattering amplitudes. The direct production and rescattering processes involving the $D_s^- \Lambda$ pair are schematically illustrated in Fig. 6.

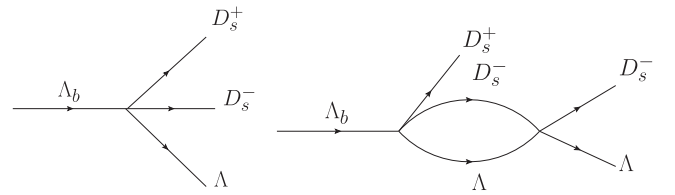
The differential decay width with respect to the $D_s^- \Lambda$ invariant mass is given by

$$\frac{d\Gamma}{dM_{\text{inv}}(D_s^- \Lambda)} = \frac{1}{(2\pi)^3} \frac{1}{4M_{\Lambda_b}^2} p_{D_s^+} \tilde{p}_{D_s^-} |\mathcal{M}|^2, \quad (22)$$

with

$$p_{D_s^+} = \frac{\lambda^{1/2}(M_{\Lambda_b}^2, m_{D_s^+}^2, M_{\text{inv}}^2(D_s^- \Lambda))}{2M_{\Lambda_b}},$$

$$\tilde{p}_{D_s^-} = \frac{\lambda^{1/2}(M_{\text{inv}}^2(D_s^- \Lambda), m_{D_s^-}^2, m_{\Lambda}^2)}{2M_{\text{inv}}(D_s^- \Lambda)}, \quad (23)$$

FIG. 6. Diagrammatic representation of the $D_s^- \Lambda$ contribution in Λ_b decay, including both direct production and rescattering effects.

where $\mathcal{M}$ denotes the full decay amplitude, which includes the direct weak decay contribution and the FSI in the $D_s^- \Lambda$ channel.

$$\mathcal{M} = -\frac{1}{\sqrt{3}}\mathcal{C} - \frac{1}{\sqrt{3}}\mathcal{C}G_{D_s^- \Lambda}(M_{\text{inv}}(D_s^- \Lambda))t_{D_s^- \Lambda, D_s^- \Lambda} \times (M_{\text{inv}}(D_s^- \Lambda)). \quad (24)$$

The first term corresponds to the tree-level amplitude, while the second term accounts for the FSI through the loop function $G_{D_s^- \Lambda}$ and the unitarized scattering amplitude $t_{D_s^- \Lambda, D_s^- \Lambda}$, both evaluated at the invariant mass $M_{\text{inv}}(D_s^- \Lambda)$.

III. RESULTS

In this section, we present the invariant mass distributions of various meson-baryon systems resulting from the three-body decay of Λ_b discussed above. These

distributions are a consequence of the near-threshold hadronic molecules and strong FSI that we obtained in our approach. We focus on subsystems such as $D^- \Lambda$, $D_s^- p$, $D_s^- \Lambda$, as well as their vector counterparts $D^{*-} \Lambda$, $D_s^{*-} p$, and $D_s^{*-} \Lambda$. This latter channels were also studied in Ref. [29] and we use the results from there. The shapes of these mass distributions, especially near threshold, once determined experimentally, can provide valuable information about the nature of the underlying interactions and the possibility of existence of dynamically generated states.

To account for theoretical uncertainties, we perform calculations using different values of the cutoff parameter q_{max} ranging from 550 to 650 MeV, consistent with earlier works [29]. For each case, we also compare the results with normalized phase space distributions to highlight dynamical effects beyond simple kinematics.

In all the calculation the normalization constant is taken as $\mathcal{C} = 1$. This means that the results are given in arbitrary normalization. However, and this is important, the relative

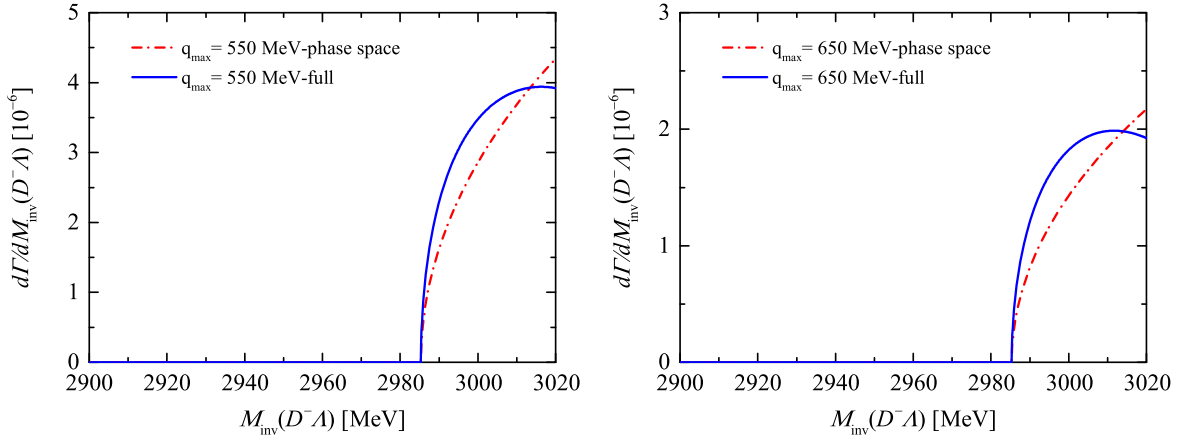


FIG. 7. Invariant mass distributions of $\frac{d\Gamma}{dM_{\text{inv}}(D^- \Lambda)}$ in the $\Lambda_b \rightarrow D^+ D^- \Lambda$ decay with $q_{\text{max}} = 550$ MeV and 650 MeV, showing phase space and full contributions. The red dot-dashed line shows the phase space, while the blue solid line represents the full contributions, respectively.

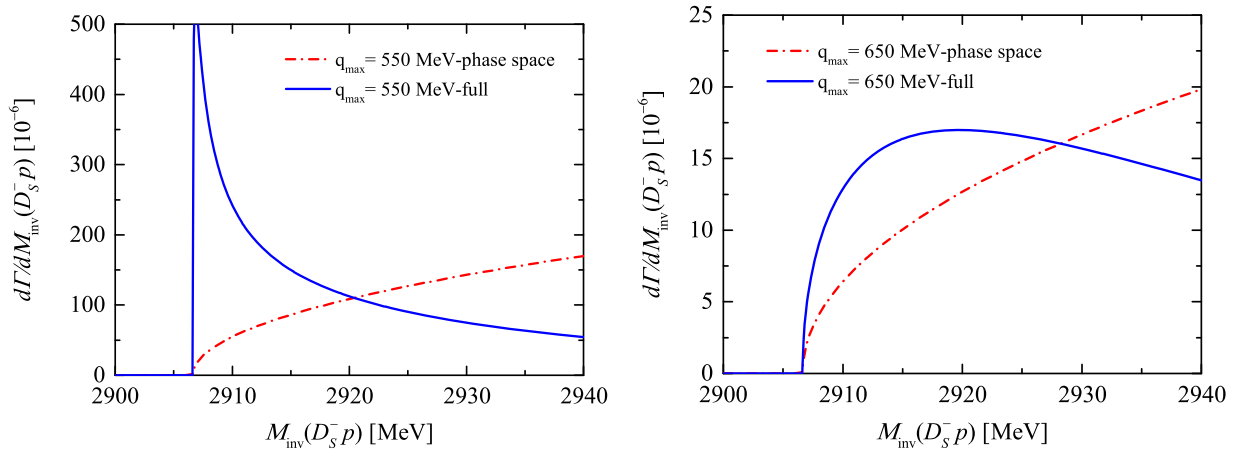


FIG. 8. Invariant mass distribution of $\frac{d\Gamma}{dM_{\text{inv}}(D_s^- p)}$ in the $\Lambda_b \rightarrow D^0 D_s^- p$ decay. Line styles and q_{max} values are the same as in Fig. 7.

TABLE I. Poles obtained for $S = -1, I = 1/2$ and $S = -2, I = 0$, unit in MeV. The thresholds of the channels are written in parentheses.

$q_{\max}$	$S = -1, I = 1/2 (D^- \Lambda, D_s^- p)$	$q_{\max}$	$S = -2, I = 0 (D_s^- \Lambda)$
550	2906 (2980.52, 2906.62)	560	3083 (3084.03)
650	2880	650	3046
750	2822	750	2959
$q_{\max}$	$S = -1, I = 1/2 (D^{*-} \Lambda, D_s^{*-} p)$	$q_{\max}$	$S = -2, I = 0 (D_s^{*-} \Lambda)$
550	3049 (3122.53, 3050.47)	550	3227 (3227.88)
650	3020	650	3184
750	2975	750	3.111

weights of all the mass distributions that we present are a prediction of the theory.

A. Mass distributions of $D^- \Lambda$ and $D_s^- p$

Figures 7 and 8 show the invariant mass distributions $\frac{d\Gamma}{dM_{\text{inv}}(D^- \Lambda)}$ and $\frac{d\Gamma}{dM_{\text{inv}}(D_s^- p)}$ as functions of their respective invariant masses. We also plot the phase space contribution normalized to the same area in the region chosen of about 40 MeV above threshold. Both distributions for the full mechanism exhibit strong enhancements near threshold, which point to strong final state interactions and the existence of molecular states of our approach. To understand these results we show in Table I the results obtained in Ref. [29] for the $D^- \Lambda, D_s^- p$ channel and in $D_s^- \Lambda$.

These results correspond to the case of octet baryons, and consider strangeness $S = -1$ with isospin $I = 1/2$. From Table I, we can see that for $q_{\max} = 550$ MeV, we find a barely bound state with respect to the $D_s^- p$ threshold (2906.62 MeV). This is the reason for the type of enhancement found in Fig. 8, with respect to phase space. When we use $q_{\max} = 650$ MeV, the state is bound by 26 MeV, and the enhancement seen in Fig. 8 is not so big. Even then there is a clear difference of the shape of the mass distribution with respect to phase space.

Note that with respect to $D^- \Lambda$, the bound states are further away, and hence, as seen in Fig. 7, the enhancement of the mass distribution close to the $D^- \Lambda$ threshold, with respect to phase space is not significant.

Our results show which is the decay mode most favorable to see the predicted bound states of these exotic systems, which is the $\Lambda_b \rightarrow D^0 D_s^- p$ reaction. While the exact position and magnitude of the peaks exhibit some sensitivity to the cutoff value, the presence of strong interaction effects near threshold remains well established. These findings encourage further experimental investigation in this invariant mass region, particularly with higher-statistics data.

B. Mass distribution of $D_s^- \Lambda$

Figure 9 shows the invariant mass distribution $\frac{d\Gamma}{dM_{\text{inv}}(D_s^- \Lambda)}$. This channel belongs to the $S = -2, I = 0$ sector and provides a unique opportunity to explore doubly strange hadronic molecules, which have not been extensively studied.

We explore the theoretical uncertainty by varying the cutoff parameter $q_{\max}$ between 550 and 650 MeV, consistent with previous works using the local hidden gauge approach [30–32].

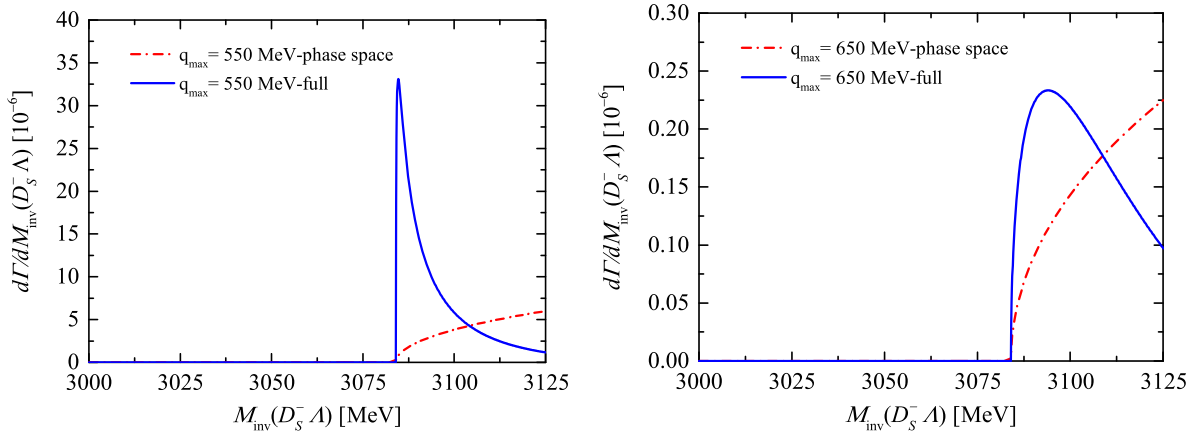


FIG. 9. Invariant mass distribution of $\frac{d\Gamma}{dM_{\text{inv}}(D_s^- \Lambda)}$. Line styles and $q_{\max}$ values are the same as in Fig. 7.

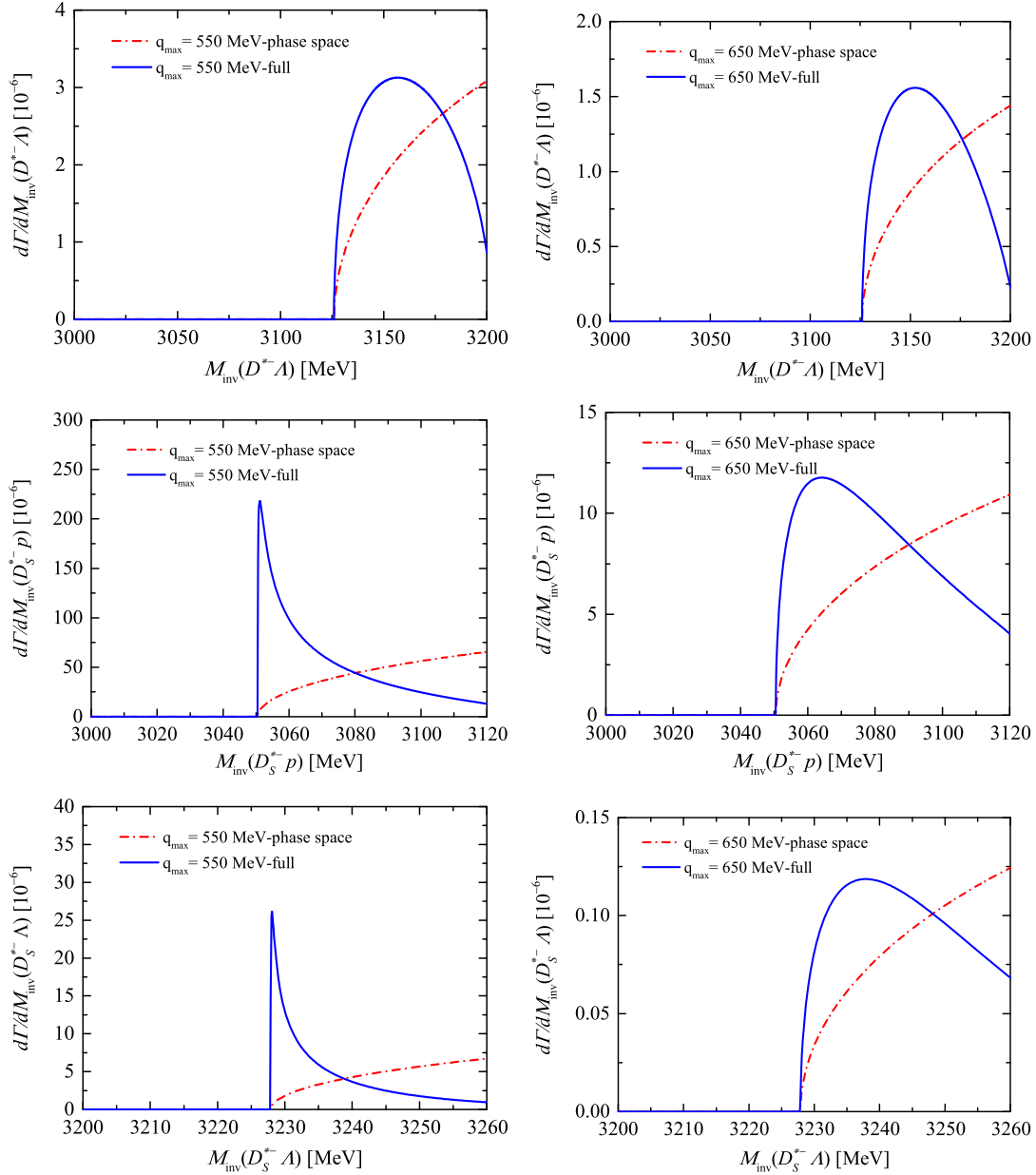


FIG. 10. Invariant mass distributions of $D^{*-}\Lambda$ in the $\Lambda_b \rightarrow D^{*+} D^{*-} \Lambda$ decay, $D_s^{*-} p$ in the $\Lambda_b \rightarrow D_s^{*0} D_s^{*-} p$ decay, and $D_s^{*-} \Lambda$ in the $\Lambda_b \rightarrow D_s^{*+} D_s^{*-} \Lambda$ decay. Line styles and $q_{\max}$ values are the same as in Fig. 7.

A visible enhancement near threshold is observed for all values of $q_{\max}$, as $q_{\max}$ decreases, the enhancement becomes slightly more pronounced and shifts subtly toward lower invariant masses, similar to the behavior seen in the $S = -1$ channels. This is a clear consequence of the pole getting closer to the threshold for smaller value of $q_{\max}$.

The uncertainties in the position of the poles revert on uncertainties of the mass distributions. Yet, the strong theoretical support for these bound states and the results shown here indicate that the measurements of these mass distributions will be determinant to clarify this issue and find evidence for the existence of these exotic bound states.

C. Mass distributions of $D^{*-}\Lambda$, $D_s^{*-} p$, and $D_s^{*-} \Lambda$

In Fig. 10, we show the invariant mass distributions $\frac{d\Gamma}{dM_{\text{inv}}}$ for the vector meson-baryon systems: $D^{*-}\Lambda$, $D_s^{*-} p$, and $D_s^{*-} \Lambda$. These channels extend our investigation into vector meson interactions in the same strangeness sectors as the pseudoscalar cases.

The $D^{*-}\Lambda$ and $D_s^{*-} p$ distributions, both belonging to the $S = -1$, $I = 1/2$ sector, exhibit noticeable threshold enhancements, as a consequence of the existence of bound states in this sector. The $D_s^{*-} \Lambda$ mass distribution, with $S = -2$, $I = 0$, also shows an enhancement near the threshold.

We again perform an uncertainty analysis by varying $q_{\max}$ from 550 to 650 MeV. The corresponding results are displayed in the figures and show that the qualitative features of the threshold behavior are stable. The phase space distributions, normalized to the same area, are included for comparison and demonstrate the clear impact of final state interactions on the line shapes.

These findings suggest that vector meson-baryon systems deserve further attention in the search for exotic states, complementing the results found in pseudoscalar channels.

IV. FURTHER CONSIDERATIONS

To finalize the work, we discuss several issues that could lead to further uncertainties in our results in order to better control the accuracy of the conclusions.

A. Further interactions: Three body effects and triangle singularities

We take the $\Lambda_b \rightarrow D^+ D^- \Lambda$ as an example

- (1) Interaction of the $D^+ \Lambda$ and $D^+ D^-$ pairs

Our analysis in Fig. 4 only took into account the $D^- \Lambda$ interaction, since we are interested in the bound state of this system. In the reaction, we would also have $D^+ D^-$ and $D^+ \Lambda$ rescattering. These mechanisms are not negligible when considering the whole $\Lambda_b \rightarrow D^+ D^- \Lambda$ reaction in the full ranges of invariant masses. However, as we show below, they have a negligible effect on the $D^- \Lambda$ mass distribution around the $D^- \Lambda$ threshold. To see that, we look at the Dalitz plot in Fig. 11 for the M_{12}, M_{23} distributions, assuming $D^+(1)$, $\Lambda(2)$, $D^-(3)$.

The $D^+ \Lambda$ interaction is not negligible, and together with other coupled channels it leads to the production of Ξ_c excited states. This has been studied in detail in [67]. There the ΛD , together with the $\Xi_c \pi$, $\Xi'_c \pi$, $\Lambda_c \bar{K}$, $\Sigma_c \bar{K}$, ΛD , $\Xi_c \eta$, ΣD , $\Xi'_c \eta$, $\Omega_c K$, and ΞD_s channels, was studied within

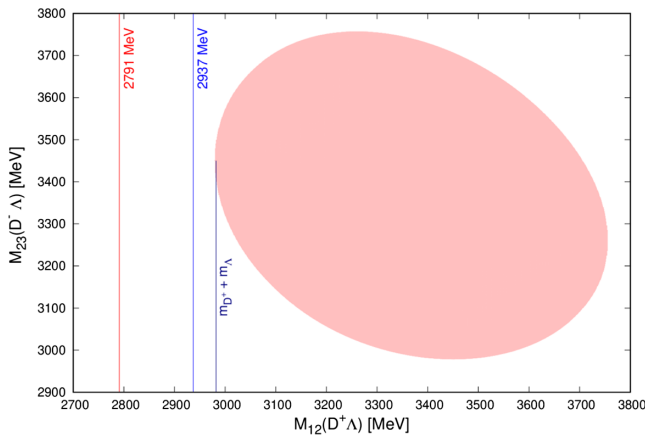


FIG. 11. Dalitz plot for M_{12}, M_{23} with $D^+(1)$, $\Lambda(2)$, and $D^-(3)$.

the local hidden gauge approach, and states were found that could be associated with some Ξ_c experimental states found in [68–70]. One state was found at 2791 MeV and the other one at 2937 MeV, with widths smaller than 15 MeV. The first state was dominated by the ΣD channel and the second one by ΛD . Hence, we shall see how they could influence the $D^- \Lambda$ mass distribution.

The two mentioned states are below the $D^+ \Lambda$ threshold, as we observe in Fig. 11. The states are relatively far away from the $D^+ \Lambda$ threshold, and if anything they could influence the $D^+ \Lambda$ mass distribution close to its threshold. However, at the $D^+ \Lambda$ threshold, $M(D^- \Lambda)_{\max}$ and $M(D^- \Lambda)_{\min}$ can be calculated using the formula of the PDG [68].

$$(M_{23})_{\max} = (E_2^* + E_3^*)^2 - \left(\sqrt{E_2^{*2} - m_2^2} - \sqrt{E_3^{*2} - m_3^2} \right)^2,$$

$$(M_{23})_{\min} = (E_2^* + E_3^*)^2 - \left(\sqrt{E_2^{*2} - m_2^2} + \sqrt{E_3^{*2} - m_3^2} \right)^2,$$

$$E_2^* = \frac{M_{12}^2 - m_1^2 + m_2^2}{2M_{12}},$$

$$E_3^* = \frac{M^2 - M_{12}^2 - m_3^2}{2M_{12}}$$

with M the D^+ mass, where E_2^*, E_3^* are the energies of particles 2, 3 in the rest frame of particles 1, 2. Clearly, if $M_{12} = m_1 + m_2$, then $E_2^* = m_2$ and $(M_{23})_{\max} = (M_{23})_{\min}$, as one can see in Fig. 11 in the extreme left side of the Dalitz plot. We find

$$(M_{23})_{\max} = (M_{23})_{\min} = M(D^- \Lambda) = 3450 \text{ MeV}$$

However, the $D^- \Lambda$ threshold is at 2983 MeV, which is 467 MeV apart from that energy. It is then clear that the effect of these resonances in the $D^- \Lambda$ mass distribution close to the $D^- \Lambda$ threshold should be negligible.

Next, we perform a similar discussion to examine the possible effect of the $D^+ D^-$ interaction in the $D^- \Lambda$ mass distribution close to threshold. For this, we take now $D^+(1)$, $D^-(2)$, $\Lambda(3)$ and show the Dalitz plot in Fig. 12.

The $D^+ D^-$ interaction is also attractive and it leads to a bound state around 3720 MeV according to [71,72]. This bound state has also been corroborated by lattice QCD calculations [73], and there are claims of experimental evidence by comparing the theoretical results of [74,75] with data from the $\gamma\gamma \rightarrow D\bar{D}$ reaction [76,77]. Once again, this state would affect the $D^+ D^-$ mass distribution close to threshold, but following the same argument as before, the threshold mass of $D^+ D^-$ corresponds to

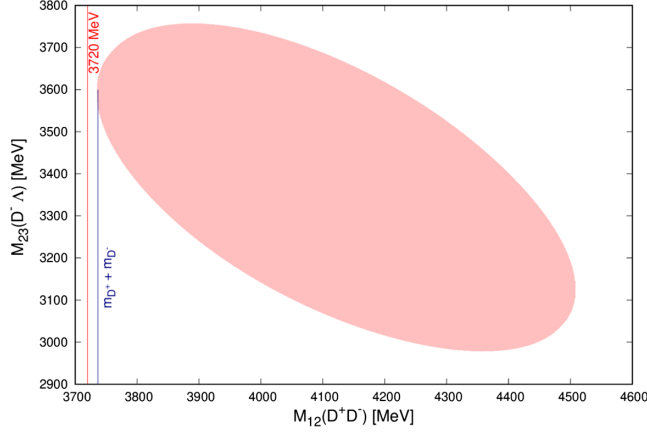


FIG. 12. Dalitz plot for M_{12} , M_{23} with $D^+(1)$, $D^-(2)$, and $\Lambda(3)$.

$$(M(D^-\Lambda))_{\max} = (M(D^-\Lambda))_{\min} = 3594 \text{ MeV}$$

Once again, this mass is very far away from the $D^-\Lambda$ threshold (611 MeV) to have any influence in the $D^-\Lambda$ mass distribution close to threshold.

(2) Three body effects

We considered the $D^-\Lambda$ interaction close to threshold. In Fig. 4 the D^+ was considered a spectator. Yet, once the $D^-\Lambda$ has a scattering, the D^+D^- would undergo rescattering, as shown in Fig. 13(a), or the $D^+\Lambda$ could rescatter as shown in Fig. 13(b).

Once again, we make use of the Dalitz plot in Fig. 14 and take the particles as $\Lambda(1)$, $D^-(2)$, $D^+(3)$. At the $D^-\Lambda$ threshold $M(D^+D^-)$ is given by

$$(M(D^+D^-))_{\max} = (M(D^+D^-))_{\min} = 4358 \text{ MeV}$$

This energy is very far away from the 3720 MeV, to be influenced by the D^+D^- bound state.

Similarly we plot in Fig. 15 the Dalitz plot for the particles $D^-(1)$, $\Lambda(2)$, $D^+(3)$

At the threshold of the $D^-\Lambda$ we have now

$$(M(D^+\Lambda))_{\max} = (M(D^+\Lambda))_{\min} = 3450 \text{ MeV}$$

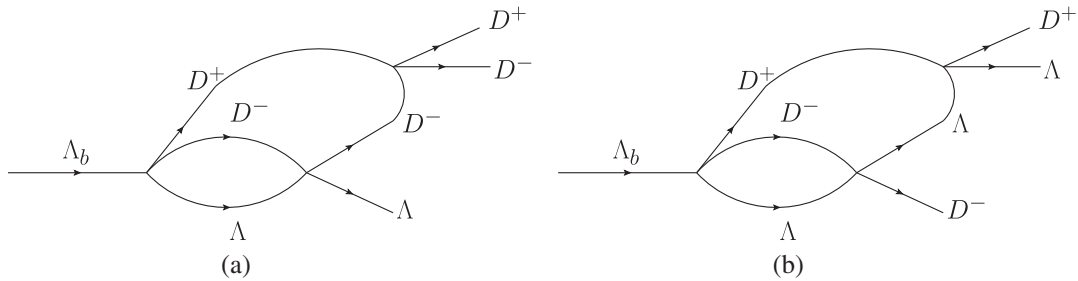


FIG. 13. Three body rescattering; (a) of D^+D^- after $D^-\Lambda$ and (b) of $D^+\Lambda$ after $D^-\Lambda$.

Once again this energy is very far away from 2791 MeV and 2937 MeV to be influenced by the resonances appearing below the $D^+\Lambda$ threshold with these energies. What we have seen is that by choosing to work close to the $D^-\Lambda$ threshold, where we expect to see effects of the $D^-\Lambda$ bound state, we can safely disregard the D^+D^- and $D^+\Lambda$ interaction and possible three body effects.

(3) Possible triangle singularities

Triangle singularities (TS) [78–80] do sometimes show up, producing peaks which could be misidentified with some resonance. To have a TS in our case, need a Λ_b decay into two particles and one of them should also decay, with one of the products reinteracting with the second particle of the Λ_b decay to produce the desired final state. In order to magnify the chances to obtain a triangle singularity we choose the $\Lambda_b \rightarrow \Lambda_c D_s^-$ decay with a branching ratio of 1.1% (second largest hadronic decay in the PDG [68]). Then we get a triangle diagram as shown in Fig. 16.

The triangle diagram leads to a TS when the three intermediate particles can be placed simultaneously on shell, they are collinear, and the process can occur at the classical level (Norton-Coleman theorem [81]). However, the Λ_c^+ , with a mass of 2286 MeV cannot decay to $D_s^+\Lambda$ with a combined mass of 3083 MeV. This large difference of masses makes the triangle diagram of Fig. 16 to be very far away from a TS and we can disregard it. We could think of having a Λ_c state with higher mass. The highest mass reported in the PDG has a mass of 2940 MeV, still insufficient to decay to $D_s^+\Lambda$. We must go to an intermediate Λ_c^+ with a mass of 3500 MeV to reach a TS for the diagram of Fig. 16. This can be tested with Eq. (18) of [80], which offers a simple and practical way to see the fulfillment of Norton-Coleman conditions. It is unclear whether a Λ_c^+ with this mass exists, or if the Λ_b decays to this Λ_c^+ and D_s^- , and also if the Λ_c^+ decays to $D_s^+\Lambda_c$ and with which rates. However, this is not all. The fact that we have the elastic $D_s^-\Lambda \rightarrow D_s^-\Lambda$ transition in

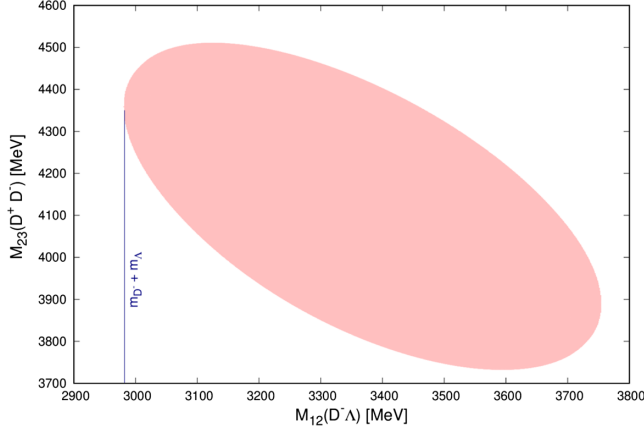


FIG. 14. Dalitz plot for M_{12} versus M_{23} with $\Lambda(1)$, $D^-(2)$, and $D^+(3)$.

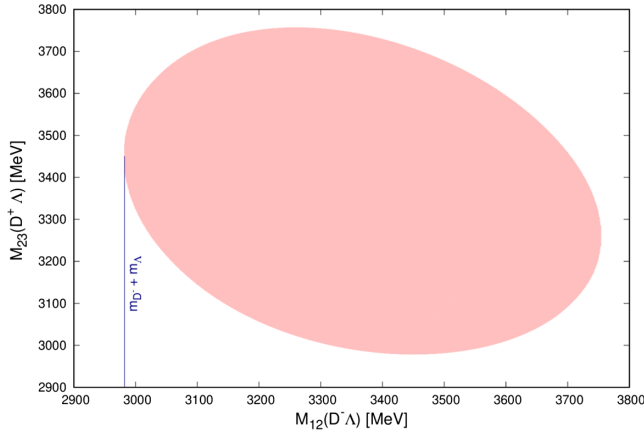


FIG. 15. Dalitz plot for M_{12} , M_{23} with $D^-(1)$, $\Lambda(2)$, $D^+(3)$.

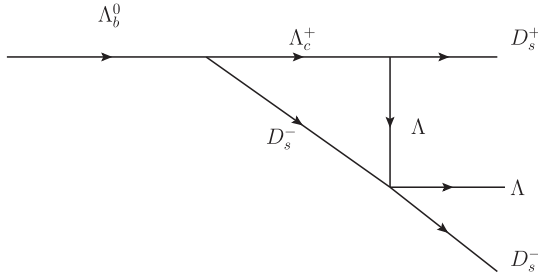


FIG. 16. Triangle mechanism for $\Lambda_b \rightarrow D_s^- \Lambda D_s^+$.

the loop makes one face the Schmid theorem [82], that states, that even if we found some combination of particles to make the TS conditions fulfilled, the TS appearing from the triangle diagram is reabsorbed into the tree level diagram (same diagram of Fig. 16 without $D_s^- \Lambda$ rescattering) through a phase change in that amplitude and there are no visible effects of the

TS in the mass distributions (see detailed discussions and implications of Schmid theorem in [83]).

The former discussion shows in a particular case how difficult is it to find a TS contributing to a desired process in a particular kinematical region. A review of TS is given in [84], showing that TS appear in some very particular reactions and particular kinematics, the signals are always small, and no case is known where it competes with other mechanisms that produce strong signals for real resonances.²

B. SU(3) breaking considerations

The Lagrangians that we have used in [29] are SU(3) symmetric. However upon unitarization in coupled channels, using physical masses, SU(3) symmetry is broken. One example where this is studied in detail is given in [41], where the interaction of the KN , $\pi\Sigma$, $\pi\Lambda$, $\eta\Sigma$, $\eta\Lambda$, $K\Xi$ channels is studied within the chiral unitary approach. The chiral Lagrangian used is SU(3) symmetric, which makes the transition potential fulfill this symmetry. When the Bethe-Salpeter equation in coupled channel is used to generate scattering amplitudes, using equal masses for the mesons and equal masses for the baryons, there is still SU(3) symmetry in the results and one generates a singlet bound state and two degenerate octet states. Then one breaks gradually the SU(3) symmetry approaching the physical masses of the mesons and the baryons, and two branches appear for the octet with $I = 0$, two branches for $I = 1$, and the singlet also evolves. One of the octet becomes the $\Lambda(1420)$ and the former singlet evolves to the $\Lambda(1380)$, the two new states of the old $\Lambda(1405)$. The interesting thing is that the original SU(3) octet state gives now rise to the $\Lambda(1420)$ and $\Lambda(1670)$ states simply because of the SU(3) breaking in the unitarization due to different masses of physical particles in the same multiplets.

Since the unitarization, which has to do with transition probabilities, requires the use of physical masses to be realistic, one considers those results realistic and not those obtained by imposing SU(3) symmetry also in the unitarization. This said, we take advantage here to complete the work of Ref. [29] and show the couplings that we obtain for the case when SU(3) symmetry is also used in the unitarization. We show the results in Table II for the case of the $S = -1$ sector, with $I = 1/2$, and one value of $q_{\max} = 600$ MeV, by taking an average meson and baryon masses

²A case apparently violating this rule seems to be the claim of [84] suggesting a contribution from a TS to eventually being responsible for the peak attributed to the pentaquark peaks $P_c(4440)$, $P_c(4457)$, which can be produced by the $\bar{D}^* \Sigma_c$ interaction with coupled channels, as shown in Ref. [85] where the states were predicted. However, this is not the case as shown in [80], where it was shown that the process suggested involved a transition amplitude in P -wave at threshold and was negligible.

TABLE II. Poles (in MeV) couplings of the $S = -1, I = 1/2$ state with isospin $I = 1/2$. results in parentheses are for the full SU(3) symmetric case.

$q_{\max}$	Pole		$\bar{D}_s N$	$\bar{D} \Lambda$	$\bar{D} \Sigma$	$\sum_i X_i$
600	2898(2948)	g_i	2.51(3.69)	1.64(1.51)	3.59(4.52)	
		$g_i G_i$	-17.89(-16.89)	-7.23(-6.89)	-11.52(-20.68)	
		X_i	0.77(0.36)	0.06(0.06)	0.14(0.54)	~ 1

$$\bar{M}_m = \frac{1}{2}(m_{\bar{D}} + m_{\bar{D}_s}), \quad (25)$$

$$\bar{M}_B = \frac{1}{3}(m_N + m_{\Lambda} + m_{\Sigma}). \quad (26)$$

Similar results are obtained in other cases, which we find unnecessary to detail, because, as shown before, only the results with unitarization with physical masses, which are shown in Ref. [29] are meaningful.

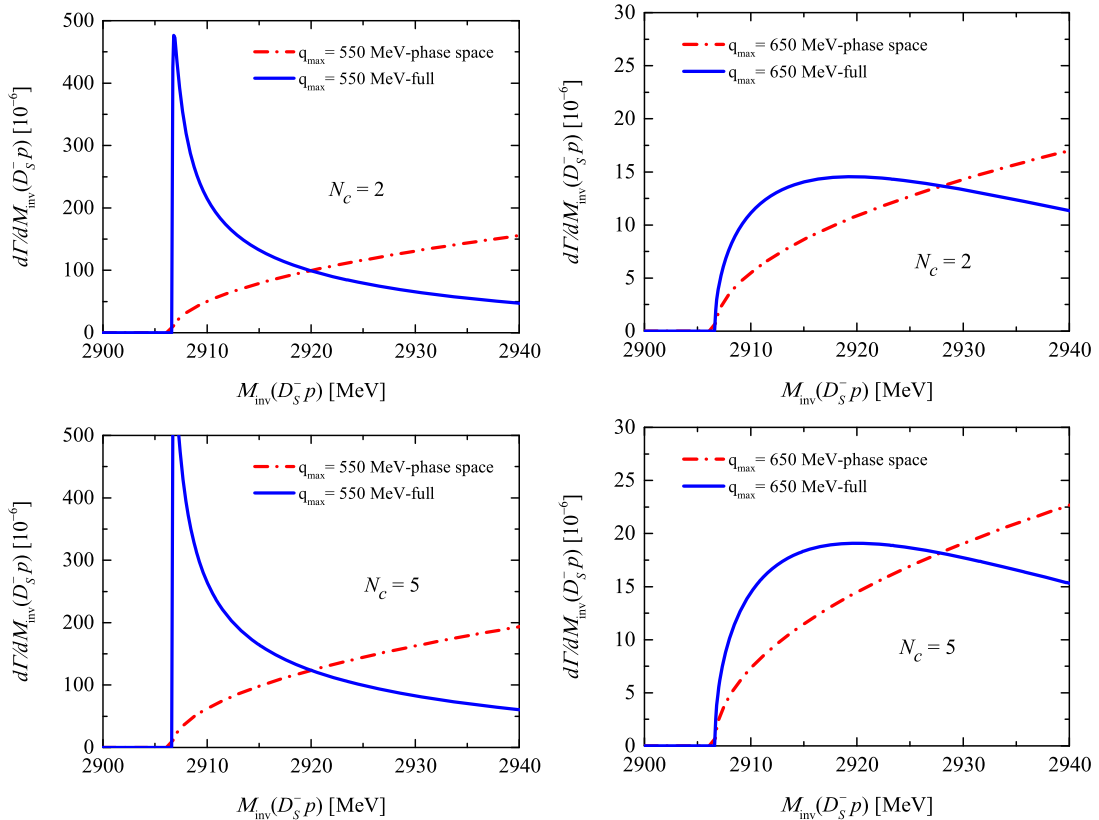
C. Dependence of results on N_c

We take here again one case to show how the results depend on N_c . The results presented so far are calculated for $N_c = 3$. Since the dynamics can be more complicated, we show here the results in Fig. 17 for different values of $N_c = 2, 5$.

In order to only change the contribution of internal emission, which goes as $1/N_c$ with respect to external emission, we keep CN_c constant in Eq. (18) when varying N_c . Comparing Fig. 17 with Fig. 8, we see that the changes are moderate, even for relatively large changes in N_c . But most important, what we observe is that compared to the phase space normalized to the same area, the changes are extremely small.

D. Further results with different $q_{\max}$

We have chosen the values of $q_{\max} = 550$ MeV and $q_{\max} = 650$ MeV, based on previous studies on meson-baryon interaction in [30], the pentaquarks with strangeness in [32] and the Ω_c state in the study of Ω_c states in [31]. We can extend the range a bit by taking into account that in the study of Ξ_c states that we have mentioned above [67], a value of $q_{\max} = 700$ MeV was used to reproduce the

FIG. 17. Results of the $D_s p$ mass distribution, by varying $N_c = 2, 5$.

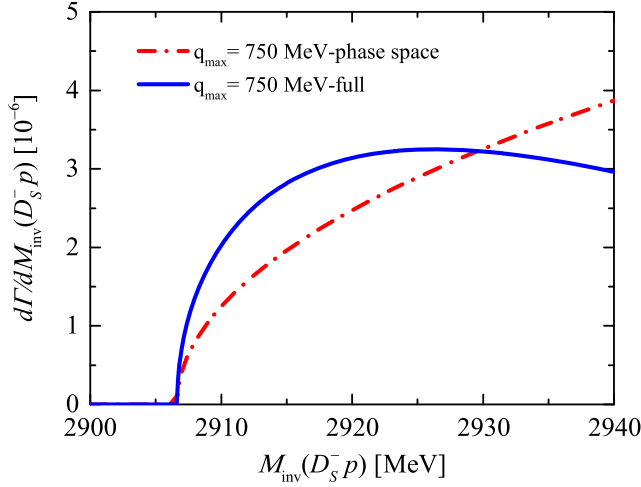


FIG. 18. The $D_s p$ mass distribution in $q_{\max} = 750$ MeV ($N_c = 3$).

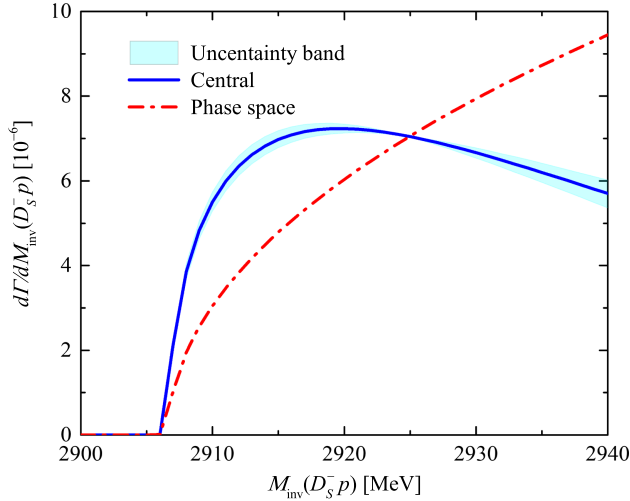


FIG. 19. Uncertainty band of $D_s p$ mass distribution, by varying N_c and $q_{\max}$.

experimental states. We then extend our results using an extra value of $q_{\max} = 750$ MeV. We consider that the range from 550 to 750 MeV is wide enough to account for possible uncertainties in our framework. The results for $q_{\max} = 750$ MeV are shown in Table I. We do not repeat all the results for all the cases but concentrate on Fig. 17, for the $D_s p$ mass distribution. The new results are shown in Fig. 18.

As we can see, the differences between the predicted shape and the phase space are still sizeable.

E. Combined uncertainties

In view of uncertainties in N_c and $q_{\max}$, we perform here an evaluation of combined uncertainties, for which we follow a standard statistical method, taking random

numbers of $N_c \in [2 - 5]$ and $q_{\max} \in [550 - 750]$ MeV. This follows the philosophy of the resampling method [86,87]. To have the band of uncertainties we eliminate 16% of the highest mass distributions and 16% of the lowest one, to have 68% confidence level. However, since we are interested in the diversion with respect to phase space, we normalize all the curves to the area of phase space, and we concentrate on Fig. 17 again, showing the uncertainty band in Fig. 19.

It is most welcome that once the mass distributions are normalized to the same area, the shapes of the different mass distributions are remarkably similar, and this tells us that, within the relatively large uncertainties explored in the present work, the shapes predicted for the mass distributions are very different from the one of phase space, which can only encourage the performance of the experiments to confirm the prediction of those new exotic states.

V. CONCLUSION

In this work, we have carried out a detailed study of the three-body decay processes $\Lambda_b \rightarrow D^+ D^- \Lambda$, $\Lambda_b \rightarrow D^0 D_s^- p$, and $\Lambda_b \rightarrow D_s^+ D_s^- \Lambda$, focusing on the invariant mass distributions of various meson-baryon subsystems. By employing a unitarized coupled-channel approach based on the local hidden gauge formalism, we have investigated both pseudoscalar and vector meson-baryon interactions in strangeness sectors $S = -1$ and $S = -2$, with particular emphasis on the effects of final state interactions on the mass distributions.

Our analysis reveals enhancements near the thresholds of the $D^- \Lambda$, $D_s^- p$, and $D_s^- \Lambda$ channels, as a consequence of the strong final state interactions and the existence of molecular exotic states. These threshold effects remain evident despite variations of the cutoff parameter $q_{\max}$, confirming the stability of the predicted structures.

In addition to pseudoscalar meson-baryon systems, we also investigate vector meson-baryon interactions such as $\Lambda_b \rightarrow D^{*+} D^{*-} \Lambda$, $D^{*0} D_s^{*-} p$, and $D_s^{*+} D_s^{*-} \Lambda$. These vector channels, sharing the same strangeness sectors, exhibit threshold behaviors that complement and enrich the search for exotic molecular states. This extended study provides a more complete picture of heavy meson-baryon dynamics and highlights the importance of considering both pseudoscalar and vector mesons in understanding the spectrum of possible exotic hadrons.

The comparison with normalized phase space distributions clearly highlights the dynamical origin of the observed mass distribution features. The study of these decay channels offers valuable insights into the interactions of multistrange meson-baryon systems and could uncover exotic states beyond the conventional hadron framework. These findings contribute to a deeper understanding of the hadronic dynamics in Λ_b decays and provide valuable theoretical input that can be tested in ongoing and future experimental investigations, particularly in the search for

exotic hadronic molecules involving open-charm mesons and strange baryons.

It is worth emphasizing that the possible influences of $D^+ \Lambda$ and $D^+ D^-$ interactions on the near-threshold mass distribution of $D^- \Lambda$ are found negligible. Three-body rescattering effects are also found to have minimal impact near the $D^- \Lambda$ threshold. The analysis of potential triangle singularities shows no significant effect on the reactions studied. Furthermore, uncertainties arising from model parameters, such as the cutoff and N_c , do not appreciably change the mass distributions when they are normalized to the same area, showing a substantial difference from phase space expectations. These considerations further reinforce the reliability of the results presented in this work.

Overall, this study underscores the importance of considering both pseudoscalar and vector meson-baryon interactions to achieve a comprehensive picture of the heavy hadron spectrum and the mechanisms that give rise to exotic hadronic states near thresholds.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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