

Triple-strangeness hidden-charm pentaquarks

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We investigate the possible existence of triple-strangeness hidden-charm pentaquark states in the off-shell coupled-channel formalism. The open-charm meson-baryon $\bar{D}_s\Omega_c$, $\bar{D}_s\Omega_c^*$, $\bar{D}_s^*\Omega_c$, and $\bar{D}_s^*\Omega_c^*$ channels are considered, together with the hidden-charm $J/\psi\Omega$ channel. The two-body kernel Feynman amplitudes are constructed from an effective Lagrangian based on hidden local symmetry and heavy-quark spin symmetry. The coupled Blankenbecler-Sugar equation is solved in the partial-wave helicity basis. We observe two triple-strangeness hidden-charm pentaquark states: $P_{c\bar{c}s s s}(4787)$ and $P_{c\bar{c}s s s}(4841)$, both with $J^P = 1/2^-$. The $P_{c\bar{c}s s s}(4787)$ couples dominantly to the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels, while the $P_{c\bar{c}s s s}(4841)$ couples almost exclusively to the $\bar{D}_s^*\Omega_c^*$ channel. The total transition cross sections of $\bar{D}_s^{(*)}\Omega_c^{(*)} \rightarrow J/\psi\Omega$ indicate that the $P_{c\bar{c}s s s}(4787)$ is clearly visible in the $J/\psi\Omega$ invariant mass spectrum, whereas the $P_{c\bar{c}s s s}(4841)$ is obscured by cusp structures and background contributions.

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I. INTRODUCTION

The LHCb Collaboration reported four hidden-charm pentaquark states: $P_{c\bar{c}}(4312)$ with the decay width $\Gamma_{P_{c\bar{c}} \rightarrow J/\psi p} = (10 \pm 5) \text{ MeV}$, $P_{c\bar{c}}(4380)$ with $\Gamma_{P_{c\bar{c}} \rightarrow J/\psi p} = (210 \pm 90) \text{ MeV}$, $P_{c\bar{c}}(4440)$ with $\Gamma_{P_{c\bar{c}} \rightarrow J/\psi p} = (21^{+10}_{-11}) \text{ MeV}$, and $P_{c\bar{c}}(4457)$ with $\Gamma_{P_{c\bar{c}} \rightarrow J/\psi p} = (6.4^{+6.0}_{-2.8}) \text{ MeV}$ [1–3]. They were all discovered by the LHCb Collaboration and have not been confirmed by any other experiments. For example, the GlueX Collaboration did not observe any signal for hidden-charm pentaquarks [4]. Interestingly, these four pentaquark states lie below the corresponding meson-baryon thresholds. For instance, $P_{c\bar{c}}(4312)$ and $P_{c\bar{c}}(4380)$ are located below the $\bar{D}\Sigma_c$ and $\bar{D}\Sigma_c^*$ thresholds, respectively, while both $P_{c\bar{c}}(4440)$ and $P_{c\bar{c}}(4457)$ are positioned below the $\bar{D}^*\Sigma_c$ threshold. Subsequently, the existence of two strangeness hidden-charm pentaquark states was reported by the LHCb Collaboration: $P_{c\bar{c}s}(4338)$ [5] and $P_{c\bar{c}s}(4459)$ [6]. The Belle Collaboration appears to have confirmed the

existence of $P_{c\bar{c}s}(4459)$ but reported a slightly larger mass, $M_{P_{c\bar{c}s}} = (4471.7 \pm 4.8 \pm 0.6) \text{ MeV}/c^2$, and a decay width of $\Gamma = (21.9 \pm 13.1 \pm 2.7) \text{ MeV}$ [7]. However, it is not clear whether $P_{c\bar{c}s}(4472)$ observed by the Belle experiment can be identified as $P_{c\bar{c}s}(4459)$. For example, Clymton *et al.* [8] suggested theoretically that $P_{c\bar{c}s}(4472)$ should be distinguished from $P_{c\bar{c}s}(4459)$. $P_{c\bar{c}s}(4459)$ and $P_{c\bar{c}s}(4472)$ are found below the $\bar{D}^*\Xi_c$ threshold, whereas $P_{c\bar{c}s}(4338)$ lies just about 2 MeV above the $\bar{D}\Xi_c$ threshold. The fact that almost all observed hidden-charm pentaquark states are located below the corresponding thresholds suggests that they are likely S -wave molecular states composed of the corresponding heavy mesons and singly heavy baryons associated with those thresholds, though there are more subtle conditions required for a molecular state.

These findings on hidden-charm pentaquarks have triggered significant interest in heavy pentaquark states (see the recent reviews [9–14] and references therein). Moreover, the CMS and LHCb Collaborations have recently announced the decays $\Lambda_b^0 \rightarrow J/\psi \Xi^- K^+$ and $\Xi_b^0 \rightarrow J/\psi \Xi^- \pi^+$ [15,16], which indicates that the existence of hidden-charm pentaquark states with strangeness $S = -2$, $P_{c\bar{c}s s}$, may soon be found. Note that there have already been theoretical predictions for such double-strangeness hidden-charm pentaquarks [17–21], and even triple-strangeness ones [20,22,23].

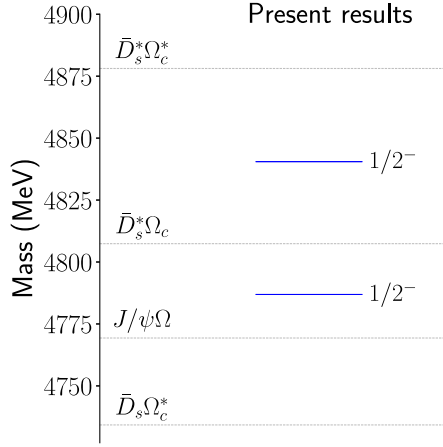
In our previous works, we have investigated the production mechanisms of hidden-charm pentaquark states with strangeness $S = 0$, $S = -1$, and $S = -2$, based on an

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FIG. 1. Predicted $P_{c\bar{c}sss}$ states from the present work.

off-shell coupled-channel approach to two-body charmed hadronic scattering [8,24,25]. We first constructed the two-body kernel Feynman amplitudes for the relevant channels by employing an effective Lagrangian that satisfies heavy-quark spin-flavor symmetry, hidden local symmetry, and chiral symmetry. We then computed the coupled-channel scattering integral equations by inserting these kernel amplitudes into them. We have seven different channels for the nonstrange hidden-charm pentaquarks, nine channels for those with $S = -1$, and 11 channels for those with $S = -2$.

The present off-shell coupled-channel formalism was successfully applied to describing various processes [26–30] in addition to the hidden-charm pentaquarks [8,24,25]. In this work, we extend the off-shell coupled-

channel formalism to examine how the triple-strangeness hidden-charm pentaquark states $P_{c\bar{c}sss}$ can be generated dynamically. To construct the transition amplitudes, we introduce five different channels: $\bar{D}_s \Omega_c$, $\bar{D}_s^* \Omega_c^*$, $J/\psi \Omega$, $\bar{D}_s^* \Omega_c$, and $\bar{D}_s \Omega_c^*$. These channels mainly consist of the lowest meson triplet (**3**) with spins 0 and 1, and the singly heavy baryon sextet (**6**) with spins 1/2 and 3/2. Figure 1 summarizes the predictions for the $P_{c\bar{c}sss}$ resonances found in the present work. Two negative-parity hidden-charm pentaquarks with $S = -3$ are obtained, as shown in Fig. 1, both with $J^P = 1/2^-$. One is located below the $\bar{D}_s^* \Omega_c^*$ threshold, and the other lies below the $\bar{D}_s^* \Omega_c$ threshold.

The present work is organized as follows. In Sec. II, we describe the off-shell coupled-channel formalism in the context of the $P_{c\bar{c}sss}$ states. We first explain how the two-body kernel Feynman amplitudes with $S = -3$ are derived from an effective Lagrangian, and then solve the scattering integral equations. In Sec. III, we present and discuss the results for the triple-strangeness hidden-charm pentaquark states. Finally, we summarize the results, draw conclusions, and provide an outlook.

II. COUPLED-CHANNEL FORMALISM

We start with the definition of the scattering amplitude

$$S_{fi} = \delta_{fi} - i(2\pi)^4 \delta(P_f - P_i) T_{fi}, \quad (1)$$

where P_i and P_f denote the total four-momenta of the initial and final states, respectively. The transition amplitude T_{fi} is constructed by using the Bethe-Salpeter equation with the two-body Feynman kernel amplitudes,

$$T_{fi}(p', p; s) = \mathcal{V}_{fi}(p', p; s) + \frac{1}{(2\pi)^4} \sum_k \int d^4 q \mathcal{V}_{fk}(p', q; s) \mathcal{G}_k(q; s) T_{ki}(q, p; s), \quad (2)$$

where p and p' are the relative four-momenta of the initial and final states, and q denotes the off-shell momentum of the intermediate states in the center-of-mass (CM) frame. The variable $s \equiv P_i^2 = P_f^2$ stands for the square of the total energy and is one of the Mandelstam variables. The indices i , f , and k represent the initial, final, and intermediate states, respectively. Since we consider coupled channels relevant to the production of $P_{c\bar{c}sss}$ states, the index k runs

over all relevant intermediate channels. The coupled integral scattering equations given in Eq. (2) are illustrated schematically in Fig. 2.

To simplify the calculation of the four-dimensional integral equations, we perform a three-dimensional reduction. Among the various methods available, we employ the Blankenbecler-Sugar formalism [31,32]. The two-body propagator is given by

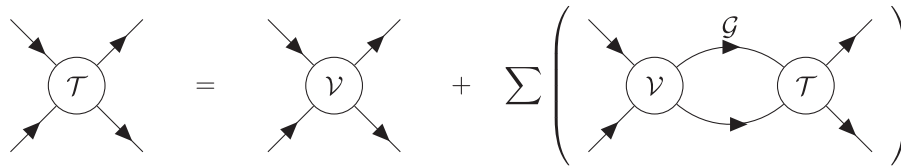


FIG. 2. Graphical representation of the coupled integral equation.

$$\mathcal{G}_k(q) = \delta\left(q_0 - \frac{E_{k1}(q) - E_{k2}(q)}{2}\right) \frac{\pi}{E_{k1}(q)E_{k2}(q)s - E_k^2(q)}, \quad (3)$$

where E_k is the total on-shell energy of the intermediate state, defined as $E_k = E_{k1} + E_{k2}$, and $\mathbf{q}$ denotes the three-

momentum of the intermediate state. Note that the spinor contributions from the meson-baryon propagator G_k have been absorbed into the matrix elements of $\mathcal{V}$ and $\mathcal{T}$. Using Eq. (3), we obtain the following coupled integral equations:

$$\mathcal{T}_{fi}(\mathbf{p}', \mathbf{p}) = \mathcal{V}_{fi}(\mathbf{p}', \mathbf{p}) + \frac{1}{(2\pi)^3} \sum_k \int \frac{d^3q}{2E_{k1}(q)E_{k2}(q)} \mathcal{V}_{fk}(\mathbf{p}', \mathbf{q}) \frac{E_k(q)}{s - E_k^2(q) + i\epsilon} \mathcal{T}_{ki}(\mathbf{q}, \mathbf{p}), \quad (4)$$

where $\mathbf{p}$ and $\mathbf{p}'$ denote the relative three-momenta of the initial and final states, respectively, in the CM frame.

To generate the $P_{c\bar{c}ss}$ states, we construct two-body coupled channels by combining the charmed meson triplet ($\mathbf{3}_0$ and $\mathbf{3}_1$) with the singly charmed baryon sextet ($\mathbf{6}_{1/2}$ and $\mathbf{6}_{3/2}$), while fixing the strangeness to $S = -3$. We also include the $J/\psi\Omega$ channel, given that both the $P_{c\bar{c}}$ and $P_{c\bar{c}s}$ states were observed in the invariant mass spectra of $J/\psi N$ and $J/\psi \Lambda$, respectively. Thus, we consider five distinct channels: $\bar{D}_s\Omega_c$, $\bar{D}_s\Omega_c^*$, $J/\psi\Omega$, $\bar{D}_s^*\Omega_c$, and $\bar{D}_s^*\Omega_c^*$. The two-body kernel Feynman amplitudes are derived from the one-meson-exchange diagrams shown in Fig. 3. We exclude s -channel pole diagrams, since our aim is to investigate how the $P_{c\bar{c}ss}$ states can be dynamically generated. We also omit the u -channel diagrams, as they involve doubly charmed baryons with masses of approximately 3.5 GeV, which makes their contributions significantly suppressed compared to the t -channel contributions.

The vertex interactions are determined by using an effective Lagrangian that satisfies heavy-quark spin symmetry, hidden local symmetry, and flavor SU(3) symmetry [33], which are given by

$$\mathcal{L}_{PPV} = -i\frac{\beta g_V}{\sqrt{2}} P_a^\dagger \overleftrightarrow{\partial}_\mu P_b \mathbb{V}_{ba}^\mu + i\frac{\beta g_V}{\sqrt{2}} P_a'^\dagger \overleftrightarrow{\partial}_\mu P_b' \mathbb{V}_{ab}^\mu, \quad (5)$$

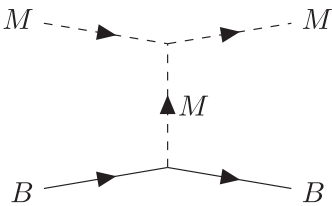


FIG. 3. t -channel meson-exchange diagrams. M and B denote a meson and a baryon involved in the process, respectively.

$$\mathcal{L}_{PP\sigma} = -2g_\sigma M P_a^\dagger \sigma P_a - 2g_\sigma M P_a'^\dagger \sigma P_a', \quad (6)$$

$$\begin{aligned} \mathcal{L}_{P^*P^*\mathbb{P}} &= -\frac{g}{f_\pi} \epsilon^{\mu\nu\alpha\beta} P_{a\nu}^{*\dagger} \overleftrightarrow{\partial}_\mu P_{b\beta}^* \partial_\alpha \mathbb{P}_{ba} \\ &\quad - \frac{g}{f_\pi} \epsilon^{\mu\nu\alpha\beta} P_{a\nu}'^{*\dagger} \overleftrightarrow{\partial}_\mu P_{b\beta}'^* \partial_\alpha \mathbb{P}_{ab}, \end{aligned} \quad (7)$$

$$\begin{aligned} \mathcal{L}_{P^*P^*\mathbb{V}} &= i\frac{\beta g_V}{\sqrt{2}} P_{a\nu}^{*\dagger} \overleftrightarrow{\partial}_\mu P_{b\nu}^* \mathbb{V}_{ba}^\mu + i2\sqrt{2}\lambda g_V M^* P_{a\mu}^{*\dagger} P_{b\nu}^* \mathbb{V}_{ba}^{\mu\nu} \\ &\quad - i\frac{\beta g_V}{\sqrt{2}} P_{a\nu}'^{*\dagger} \overleftrightarrow{\partial}_\mu P_{b\nu}'^* \mathbb{V}_{ab}^\mu - i2\sqrt{2}\lambda g_V M^* P_{a\mu}'^{*\dagger} P_{b\nu}'^* \mathbb{V}_{ab}^{\mu\nu}, \end{aligned} \quad (8)$$

$$\mathcal{L}_{P^*P^*\sigma} = 2g_\sigma M^* P_{a\mu}^{*\dagger} \sigma P_a^{*\mu} + 2g_\sigma M^* P_{a\mu}'^{*\dagger} \sigma P_a'^{\mu}, \quad (9)$$

$$\begin{aligned} \mathcal{L}_{P^*P^*\mathbb{P}} &= -\frac{2g}{f_\pi} \sqrt{MM^*} (P_a^\dagger P_{b\mu}^* + P_{a\mu}^\dagger P_b) \partial^\mu \mathbb{P}_{ba} \\ &\quad + \frac{2g}{f_\pi} \sqrt{MM^*} (P_a'^\dagger P_{b\mu}'^* + P_{a\mu}'^\dagger P_b') \partial^\mu \mathbb{P}_{ab}, \end{aligned} \quad (10)$$

$$\begin{aligned} \mathcal{L}_{P^*P^*\mathbb{V}} &= -i\sqrt{2}\lambda g_V \epsilon^{\beta\alpha\mu\nu} (P_a^\dagger \overleftrightarrow{\partial}_\beta P_{b\alpha}^* + P_{a\alpha}^\dagger \overleftrightarrow{\partial}_\beta P_b) (\partial_\mu \mathbb{V}_\nu)_{ba} \\ &\quad - i\sqrt{2}\lambda g_V \epsilon^{\beta\alpha\mu\nu} (P_a'^\dagger \overleftrightarrow{\partial}_\beta P_{b\alpha}'^* + P_{a\alpha}'^\dagger \overleftrightarrow{\partial}_\beta P_b') (\partial_\mu \mathbb{V}_\nu)_{ab}, \end{aligned} \quad (11)$$

where $\overleftrightarrow{\partial} = \overrightarrow{\partial} - \overleftarrow{\partial}$. Here, σ represents the lowest-lying isoscalar-scalar meson, i.e., $f_0(500)$. Heavy mesons and antiheavy mesons $P^{(*)}$ and $P'^{(*)}$ are defined as

$$\begin{aligned} P &= (D^0, D^+, D_s^+), & P_\mu^* &= (D_\mu^{*0}, D_\mu^{*+}, D_{s\mu}^{*+}), \\ P' &= (\bar{D}^0, D^-, D_s^-), & P_\mu'^* &= (\bar{D}_\mu'^{*0}, D_\mu'^{-}, D_{s\mu}'^{-}), \end{aligned} \quad (12)$$

while the light pseudoscalar and vector mesons are expressed as

$$\mathbb{P} = \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix}, \quad \mathbb{V}_\mu = \begin{pmatrix} \frac{1}{\sqrt{2}}\rho_\mu^0 + \frac{1}{\sqrt{2}}\omega_\mu & \rho_\mu^+ & K_\mu^{*+} \\ \rho_\mu^- & -\frac{1}{\sqrt{2}}\rho_\mu^0 + \frac{1}{\sqrt{2}}\omega_\mu & K_\mu^{*0} \\ K_\mu^{*-} & \bar{K}_\mu^{*0} & \phi_\mu \end{pmatrix}. \quad (13)$$

Since we have already determined all the coupling constants pertinent also to this work, we will briefly mention their values. The value of g_V in Eqs. (5), (8), and (11) is obtained from the Kawarabayashi-Suzuki-Riazuddin-Fayyazuddin relation [34,35] $g_V = m_\rho/f_\pi \approx 5.8$, where $f_\pi = 132$ MeV. The value of β in Eqs. (5) and (8) is taken to be $\beta \approx 0.9$ by employing the vector-meson dominance in heavy-meson radiative decays. We fix $g = 0.59 \pm 0.07 \pm 0.01$ in Eqs. (7) and (10) by using the experimental data on the full width of D^{*+} [36]. The value of λ in Eq. (11) is given as $\lambda = -0.56$ GeV $^{-1}$, extracted from light-cone sum rules and lattice QCD. Note that the sign convention for λ is different from that in Ref. [36], since we use the same phase for heavy vector mesons as in Ref. [33]. For the σ meson in Eqs. (6) and (9), the coupling constant g_σ is determined from the analysis of the 2π transition of $D_s(1^+)$ in Ref. [37]: $g_\sigma = g_\pi/(2\sqrt{6})$ with $g_\pi = 3.73$.

Regarding the heavy baryon effective Lagrangian, we follow the framework developed by Ref. [38], which incorporates a more comprehensive Lagrangian formulation as outlined in Ref. [39]. The baryonic interaction vertices within the tree-level meson-exchange diagrams are

given by the following effective Lagrangian:

$$\mathcal{L}_{B_6 B_6 \mathbb{P}} = i \frac{g_1}{2f_\pi M_6} \bar{B}_6 \gamma_5 (\gamma^\alpha \gamma^\beta - g^{\alpha\beta}) \overleftrightarrow{\partial}_\alpha \partial_\beta \mathbb{P} B_6, \quad (14)$$

$$\mathcal{L}_{B_6 B_6 \mathbb{V}} = -i \frac{\beta_6 g_V}{2\sqrt{2}M_6} (\bar{B}_6 \overleftrightarrow{\partial}_\alpha \mathbb{V}^\alpha B_6) - \frac{i\lambda_6 g_V}{3\sqrt{2}} (\bar{B}_6 \gamma_\mu \gamma_\nu \mathbb{V}^{\mu\nu} B_6), \quad (15)$$

$$\mathcal{L}_{B_6 B_6 \sigma} = -l_6 (\bar{B}_6 \sigma B_6), \quad (16)$$

$$\mathcal{L}_{B_6^* B_6^* \mathbb{P}} = \frac{3g_1}{4f_\pi M_6^*} \epsilon^{\mu\nu\alpha\beta} (\bar{B}_6^* \overleftrightarrow{\partial}_\nu \partial_\alpha \mathbb{P} B_{6\beta}^*), \quad (17)$$

$$\mathcal{L}_{B_6^* B_6^* \mathbb{V}} = i \frac{\beta_6 g_V}{2\sqrt{2}M_6^*} (\bar{B}_6^* \overleftrightarrow{\partial}_\alpha \mathbb{V}^\alpha B_6^*) + \frac{i\lambda_6 g_V}{\sqrt{2}} (\bar{B}_6^* \mathbb{V}^{\mu\nu} B_{6\nu}^*), \quad (18)$$

$$\mathcal{L}_{B_6^* B_6^* \sigma} = l_6 (\bar{B}_6^* \sigma B_6^*), \quad (19)$$

$$\mathcal{L}_{B_6 B_6^* \mathbb{P}} = \frac{g_1}{4f_\pi} \sqrt{\frac{3}{M_6^* M_6}} \epsilon^{\mu\nu\alpha\beta} [(\bar{B}_6 \gamma_5 \gamma_\mu \overleftrightarrow{\partial}_\nu \partial_\alpha \mathbb{P} B_{6\beta}^*) + (\bar{B}_6^* \gamma_5 \gamma_\nu \overleftrightarrow{\partial}_\alpha \partial_\beta \mathbb{P} B_6)], \quad (20)$$

$$\mathcal{L}_{B_6 B_6^* \mathbb{V}} = \frac{i\lambda_6 g_V}{\sqrt{6}} \left[\bar{B}_6 \gamma_5 \left(\gamma_\mu + \frac{i\overleftrightarrow{\partial}_\mu}{2\sqrt{M_6^* M_6}} \right) \mathbb{V}^{\mu\nu} B_{6\nu}^* + \bar{B}_6^* \gamma_5 \left(\gamma_\nu - \frac{i\overleftrightarrow{\partial}_\nu}{2\sqrt{M_6^* M_6}} \right) \mathbb{V}^{\mu\nu} B_6 \right], \quad (21)$$

with heavy baryon fields expressed as

$$B_6 = \begin{pmatrix} \Sigma_c^{++} & \frac{1}{\sqrt{2}}\Sigma_c^+ & \frac{1}{\sqrt{2}}\Xi_c'^+ \\ \frac{1}{\sqrt{2}}\Sigma_c^+ & \Sigma_c^0 & \frac{1}{\sqrt{2}}\Xi_c'^0 \\ \frac{1}{\sqrt{2}}\Xi_c'^+ & \frac{1}{\sqrt{2}}\Xi_c'^0 & \Omega_c^0 \end{pmatrix}, \quad B_6^* = \begin{pmatrix} \Sigma_c^{*++} & \frac{1}{\sqrt{2}}\Sigma_c^{*+} & \frac{1}{\sqrt{2}}\Xi_c^{*+} \\ \frac{1}{\sqrt{2}}\Sigma_c^{*+} & \Sigma_c^{*0} & \frac{1}{\sqrt{2}}\Xi_c^{*0} \\ \frac{1}{\sqrt{2}}\Xi_c^{*+} & \frac{1}{\sqrt{2}}\Xi_c^{*0} & \Omega_c^{*0} \end{pmatrix}. \quad (22)$$

The symbol B_μ denotes the spin 3/2 Rarita-Schwinger field, which is subject to the constraints

$$p^\mu B_\mu = 0 \quad \text{and} \quad \gamma^\mu B_\mu = 0. \quad (23)$$

The coupling constants for the baryonic vertices are fixed as follows [38,40]: $\beta_6 = -12/g_V$, $\lambda_6 = -3.31$ GeV $^{-1}$,

$g_1 = 0.942$, and $l_6 = 6.2$. The sign conventions we implement are in line with those given in Refs. [40,41].

We include only the $J/\psi\Omega$ channel, considering the fact that the $P_{c\bar{c}}$ and $P_{c\bar{c}s}$ can decay into the $J/\psi N$ and $J/\psi\Lambda$ states, respectively. We expect that the $P_{c\bar{c}s s s}$ states may decay into the $J/\psi\Omega$ state, if they exist experimentally. The interaction between heavy mesons and charmonium is

expressed as the following effective Lagrangian [42]:

$$\mathcal{L}_{PPJ/\psi} = -ig_\psi M \sqrt{m_J} (J/\psi^\mu P^\dagger \overleftrightarrow{\partial}_\mu P l^\dagger) + \text{H.c.}, \quad (24)$$

$$\mathcal{L}_{P^*PJ/\psi} = ig_\psi \sqrt{\frac{MM^*}{m_J}} \epsilon^{\mu\nu\alpha\beta} \partial_\mu J/\psi_\nu (P^\dagger \overleftrightarrow{\partial}_\alpha P_{\beta}^{*\dagger} + P_{\beta}^{*\dagger} \overleftrightarrow{\partial}_\alpha P l^\dagger) + \text{H.c.}, \quad (25)$$

$$\mathcal{L}_{P^*P^*J/\psi} = ig_\psi M^* \sqrt{m_J} (g^{\mu\nu} g^{\alpha\beta} - g^{\mu\alpha} g^{\nu\beta} + g^{\mu\beta} g^{\nu\alpha}) \times (J/\psi_\mu P_\nu^{*\dagger} \overleftrightarrow{\partial}_\alpha P_{\beta}^{*\dagger}) + \text{H.c.} \quad (26)$$

Since no experimental data exist for the $J/\psi \rightarrow D\bar{D}$ decay, we follow Ref. [43], where the coupling constant g_ψ was estimated to be $g_\psi = 0.679 \text{ GeV}^{-3/2}$.

As for the coupling of baryon decuplet with baryon sextet, we employ the following effective Lagrangian with heavy quark spin symmetry:

$$\mathcal{L}_{\Omega_c D_s \Omega} = g_I \sqrt{\frac{M_{D_s} \bar{\Omega}_c}{3}} \left(\gamma_\mu - \frac{i \overleftrightarrow{\partial}_\mu}{2M_{D_s}} \right) D_s \Omega^\mu + \text{H.c.}, \quad (27)$$

$$\mathcal{L}_{\Omega_c D_s^* \Omega} = g_I \sqrt{\frac{M_{D_s^*} \bar{\Omega}_c \gamma_5}{3}} \left(\gamma_\mu + \frac{i \overleftrightarrow{\partial}_\mu}{2M_{D_s^*}} \right) \gamma_\nu D_s^{*\nu} \Omega^\mu + \text{H.c.}, \quad (28)$$

$$\mathcal{L}_{\Omega_c^* D_s \Omega} = -g_I \sqrt{M_{D_s} \bar{\Omega}_{c\mu}^* \gamma_5} D_s \Omega^\mu + \text{H.c.}, \quad (29)$$

$$\mathcal{L}_{\Omega_c^* D_s \Omega} = g_I \sqrt{M_{D_s} \bar{\Omega}_{c\mu}^* \gamma_\nu} D_s^{*\nu} \Omega^\mu + \text{H.c.} \quad (30)$$

The coupling constant g_I is of the same order as the coupling constant for the baryon octet-baryon sextet-meson triplet vertex used in our previous work [25]. We therefore use the value $g_I = 1 \text{ GeV}^{-1/2}$, which renders the $J/\psi\Omega$ channel with negligible effects on the coupled-channel dynamics.

The two-body kernel Feynman amplitude for a one-meson exchange diagram can be generically expressed as

$$\mathcal{A}_{\lambda'_1 \lambda'_2, \lambda_1 \lambda_2} = \text{IS} F^2(q^2) \Gamma_{\lambda'_1 \lambda'_2}(p'_1, p'_2) \mathcal{P}(q) \Gamma_{\lambda_1 \lambda_2}(p_1, p_2), \quad (31)$$

where λ_i and p_i are, respectively, the helicity and momentum of each particle participating in the process, while q stands for the momentum transfer. The IS factor contains the SU(3) Clebsch-Gordan coefficient and isospin factor, of which the values for each exchange diagram are listed in Table I. The vertex functions Γ can be directly determined by the effective Lagrangian. We derive the propagators $\mathcal{P}(q)$ for spin-0 and spin-1 mesons as

TABLE I. Values of the IS factors for the corresponding t -channel diagrams for the given reactions.

Reactions	Exchange particles	IS
$\bar{D}_s \Omega_c \rightarrow \bar{D}_s \Omega_c$	ϕ	1
	σ	1
$\bar{D}_s \Omega_c \rightarrow \bar{D}_s \Omega_c^*$	ϕ	1
$\bar{D}_s \Omega_c \rightarrow J/\psi \Omega$	$\bar{D}_s, \bar{D}_s^*$	1
$\bar{D}_s \Omega_c \rightarrow \bar{D}_s^* \Omega_c$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
$\bar{D}_s \Omega_c \rightarrow \bar{D}_s^* \Omega_c^*$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
$\bar{D}_s \Omega_c^* \rightarrow \bar{D}_s \Omega_c^*$	ϕ	1
	σ	1
$\bar{D}_s \Omega_c^* \rightarrow J/\psi \Omega$	$\bar{D}_s, \bar{D}_s^*$	1
$\bar{D}_s \Omega_c^* \rightarrow \bar{D}_s^* \Omega_c$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
$\bar{D}_s \Omega_c^* \rightarrow \bar{D}_s^* \Omega_c^*$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
$J/\psi \Omega \rightarrow \bar{D}_s^* \Omega_c$	$\bar{D}_s, \bar{D}_s^*$	1
$J/\psi \Omega \rightarrow \bar{D}_s^* \Omega_c^*$	$\bar{D}_s, \bar{D}_s^*$	1
$\bar{D}_s^* \Omega_c \rightarrow \bar{D}_s^* \Omega_c$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
	σ	1
$\bar{D}_s^* \Omega_c \rightarrow \bar{D}_s^* \Omega_c^*$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
$\bar{D}_s^* \Omega_c^* \rightarrow \bar{D}_s^* \Omega_c^*$	η	$\frac{1}{3}\sqrt{2}$
	ϕ	1
	σ	1

$$\mathcal{P}(q) = \frac{1}{q^2 - m^2}, \quad \mathcal{P}_{\mu\nu}(q) = \frac{1}{q^2 - m^2} \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{m^2} \right). \quad (32)$$

Since we take the finite value of the heavy-quark mass, we use the same form for the heavy mesons as for the light mesons.

We introduce a form factor at each vertex, since hadrons have finite sizes. We employ the following form of the form factor [44]

$$F(q^2) = \left(\frac{n\Lambda^2 - m^2}{n\Lambda^2 - q^2} \right)^n, \quad (33)$$

where the value of n varies depending on the momentum power present in the vertex function Γ . This parametrization has an essential merit, since we do not need to modify Λ when n changes. Note that as n approaches infinity, Eq. (33) becomes a Gaussian form. While it is difficult to fix the cutoff masses Λ in Eq. (33) due to the lack of experimental information, we make use of the physical property of hadron sizes to reduce the uncertainties associated with the cutoff masses. Investigations of the electromagnetic form factors of singly heavy baryons

[45,46] have shown that heavy hadrons exhibit more compact sizes than their lighter counterparts. In other words, the cutoff mass Λ is proportional to the inverse of the radius of the corresponding particle, $\Lambda \sim \langle r \rangle^{-1}$. This indicates that higher cutoff masses are appropriate for heavy hadrons compared to light hadrons. Following this argument, we define the reduced cutoff mass as $\Lambda_0 := \Lambda - m$, where m denotes the mass of the exchange particle. In the absence of experimental data on the $P_{c\bar{c}ss}$, we adopt a common value of 700 MeV for all reduced cutoff parameters Λ_0 . Moreover, Cheng *et al.* proposed that the cutoff mass can be related to Λ_{QCD} [47], i.e., $\Lambda = m + \eta\Lambda_{\text{QCD}}$, where η is of order unity. It depends on the exchanged particle as well as on the external particles involved in the strong interaction vertex. The value of η tends to be larger as the mass of the exchange particle increases. In addition, the corresponding cutoff masses increase with the mass of the external particles. Although this tendency cannot be proven mathematically, it

has been shown to be plausible phenomenologically in previous works [8,24–30].

Since the strong interaction is invariant under parity transformation, the number of transition amplitudes can be reduced by employing parity invariance,

$$\mathcal{A}_{-\lambda'_1-\lambda'_2,-\lambda_1-\lambda_2} = \eta(\eta')^{-1} \mathcal{A}_{\lambda'_1\lambda'_2,\lambda_1\lambda_2}, \quad (34)$$

where η and η' are defined as

$$\eta \equiv \eta_1\eta_2(-1)^{J-s_1-s_2}, \quad \eta' \equiv \eta'_1\eta'_2(-1)^{J-s'_1-s'_2}. \quad (35)$$

Here, η_i and s_i denote the intrinsic parity and spin of the i th particle, respectively, while J represents the total angular momentum. The relation in Eq. (34) significantly reduces the computational time required for numerical calculations.

To classify the spin-parity J^P of the $P_{c\bar{c}ss}$ states, we perform the partial-wave decomposition of the $\mathcal{V}$ and $\mathcal{T}$ matrices,

$$\mathcal{T}_{\lambda'\lambda}^{J(fi)}(\mathbf{p}', \mathbf{p}) = \mathcal{V}_{\lambda'\lambda}^{J(fi)}(\mathbf{p}', \mathbf{p}) + \frac{1}{(2\pi)^3} \sum_{k,\lambda_k} \int \frac{q^2 dq}{2E_{k1}E_{k2}} \mathcal{V}_{\lambda'\lambda_k}^{J(fk)}(\mathbf{p}', \mathbf{q}) \frac{E_k}{s - E_k^2 + i\epsilon} \mathcal{T}_{\lambda_k\lambda}^{J(ki)}(\mathbf{q}, \mathbf{p}), \quad (36)$$

where the helicities of the final, initial, and intermediate states are denoted by $\lambda' = \{\lambda'_1, \lambda'_2\}$, $\lambda = \{\lambda_1, \lambda_2\}$, and $\lambda_k = \{\lambda_{k1}, \lambda_{k2}\}$, respectively. The variables $\mathbf{p}'$, $\mathbf{p}$, and $\mathbf{q}$ represent the magnitudes of the corresponding three-momenta $\mathbf{p}'$, $\mathbf{p}$, and $\mathbf{q}$, respectively.

The partial-wave expansion for the kernel amplitudes $\mathcal{V}_{\lambda'\lambda}^{J(fi)}$ is given by

$$\mathcal{V}_{\lambda'\lambda}^{J(fi)}(\mathbf{p}', \mathbf{p}) = 2\pi \int d(\cos\theta) d_{\lambda_1-\lambda_2, \lambda'_1-\lambda'_2}^J(\theta) \mathcal{V}_{\lambda'\lambda}^{fi}(\mathbf{p}', \mathbf{p}, \theta), \quad (37)$$

where θ is the scattering angle and $d_{\lambda\lambda'}^J(\theta)$ denotes the reduced Wigner D functions.

Singularities arise from the two-body propagator $\mathcal{G}$ in the coupled integral equation in Eq. (36), which can be isolated to handle. So, the integral equation is decomposed into the regularized one and the singular part,

$$\mathcal{T}_{\lambda'\lambda}^{fi}(\mathbf{p}', \mathbf{p}) = \mathcal{V}_{\lambda'\lambda}^{fi}(\mathbf{p}', \mathbf{p}) + \frac{1}{(2\pi)^3} \sum_{k,\lambda_k} \left[\int_0^\infty dq \frac{qE_k}{E_{k1}E_{k2}} \frac{\mathcal{F}(q) - \mathcal{F}(\tilde{q}_k)}{s - E_k^2} + \frac{1}{2\sqrt{s}} \left(\ln \left| \frac{\sqrt{s} - E_k^{\text{thr}}}{\sqrt{s} + E_k^{\text{thr}}} \right| - i\pi \right) \mathcal{F}(\tilde{q}_k) \right], \quad (38)$$

where

$$\mathcal{F}(q) = \frac{1}{2} q \mathcal{V}_{\lambda'\lambda_k}^{fk}(\mathbf{p}', \mathbf{q}) \mathcal{T}_{\lambda_k\lambda}^{ki}(\mathbf{q}, \mathbf{p}), \quad (39)$$

and $\tilde{q}_k$ denotes the momentum $\mathbf{q}$ when $E_{k1} + E_{k2} = \sqrt{s}$. This regularization procedure is exclusively implemented when the total energy $\sqrt{s}$ is greater than the threshold energy of the k th channel E_k^{thr} . Notably, the form factors in the kernel amplitudes $\mathcal{V}$ guarantee the unitarity of the transition amplitudes in the high-momentum domain.

To perform the numerical evaluation of the $\mathcal{T}$ matrix in Eq. (38), we consider the $\mathcal{V}$ matrix in the helicity basis and

represent it in momentum space, with momenta discretized using Gaussian quadrature. The $\mathcal{T}$ matrix is then obtained via the Haftel-Tabakin matrix inversion method [48],

$$\mathcal{T} = (1 - \mathcal{V}\tilde{\mathcal{G}})^{-1} \mathcal{V}. \quad (40)$$

The $\mathcal{T}$ matrix in the helicity basis does not possess definite parity. To determine the spin-parity classification of the $P_{c\bar{c}ss}$ states, we transform the transition amplitudes into partial-wave amplitudes with well-defined parity,

$$\mathcal{T}_{\lambda'\lambda}^{J\pm} = \frac{1}{2} [\mathcal{T}_{\lambda'\lambda}^J \pm \eta_1\eta_2(-1)^{s_1+s_2+\frac{1}{2}} \mathcal{T}_{\lambda',-\lambda}^J], \quad (41)$$

where $\mathcal{T}^{J\pm}$ denotes the partial-wave transition amplitude with total angular momentum J and parity $(-1)^{J\pm 1/2}$. The factor $1/2$ ensures that no additional scaling is needed when reverting to the partial-wave component,

$$\mathcal{T}_{\lambda'\lambda}^J = \mathcal{T}_{\lambda'\lambda}^{J+} + \mathcal{T}_{\lambda'\lambda}^{J-}. \quad (42)$$

To analyze the dynamically generated $P_{c\bar{c}sss}$ resonances, we reformulate the $\mathcal{T}$ matrix in the IJL particle basis [49]. The relations between the $\mathcal{T}$ -matrix elements in these two bases are given by

$$\mathcal{T}_{L'L}^{JS'S} = \frac{\sqrt{(2L+1)(2L'+1)}}{2J+1} \sum_{\lambda'_1\lambda'_2\lambda_1\lambda_2} (L'0S'\lambda'|J\lambda')(s'_1\lambda'_1s'_2 - \lambda'_2|S'\lambda')(L0S\lambda|J\lambda)(s_1\lambda_1s_2 - \lambda_2|S\lambda)\mathcal{T}_{\lambda'_1\lambda'_2,\lambda_1\lambda_2}^J. \quad (43)$$

In the present analysis, we focus on the diagonal components of $\mathcal{T}_L^{JS}$, as these provide the dominant contributions to the production of triple-strangeness hidden-charm pentaquarks.

III. RESULTS AND DISCUSSIONS

In Fig. 4, we show the numerical results for the elastic partial-wave cross sections for the given spin-parity classifications $J^P = 1/2^-, 3/2^-, 5/2^-$, expressed as

$$\sigma^{J\pm} = \frac{2J+1}{(2s_1+1)(2s_2+1)} \sum_{\lambda'\lambda} p' |\mathcal{T}_{\lambda'\lambda}^{J\pm}|^2, \quad (44)$$

where s_1 and s_2 denote the spins of the incoming meson and baryon, respectively. In the upper left panel of Fig. 4, a peak clearly emerges in the $\bar{D}_s^*\Omega_c$ channel, which can be identified as the first triple-strangeness hidden-charm pentaquark resonance, $P_{c\bar{c}sss}(4787)$ with $J^P = 1/2^-$. This peak also appears in the $\bar{D}_s^*\Omega_c^*$ channel. On the other hand, we

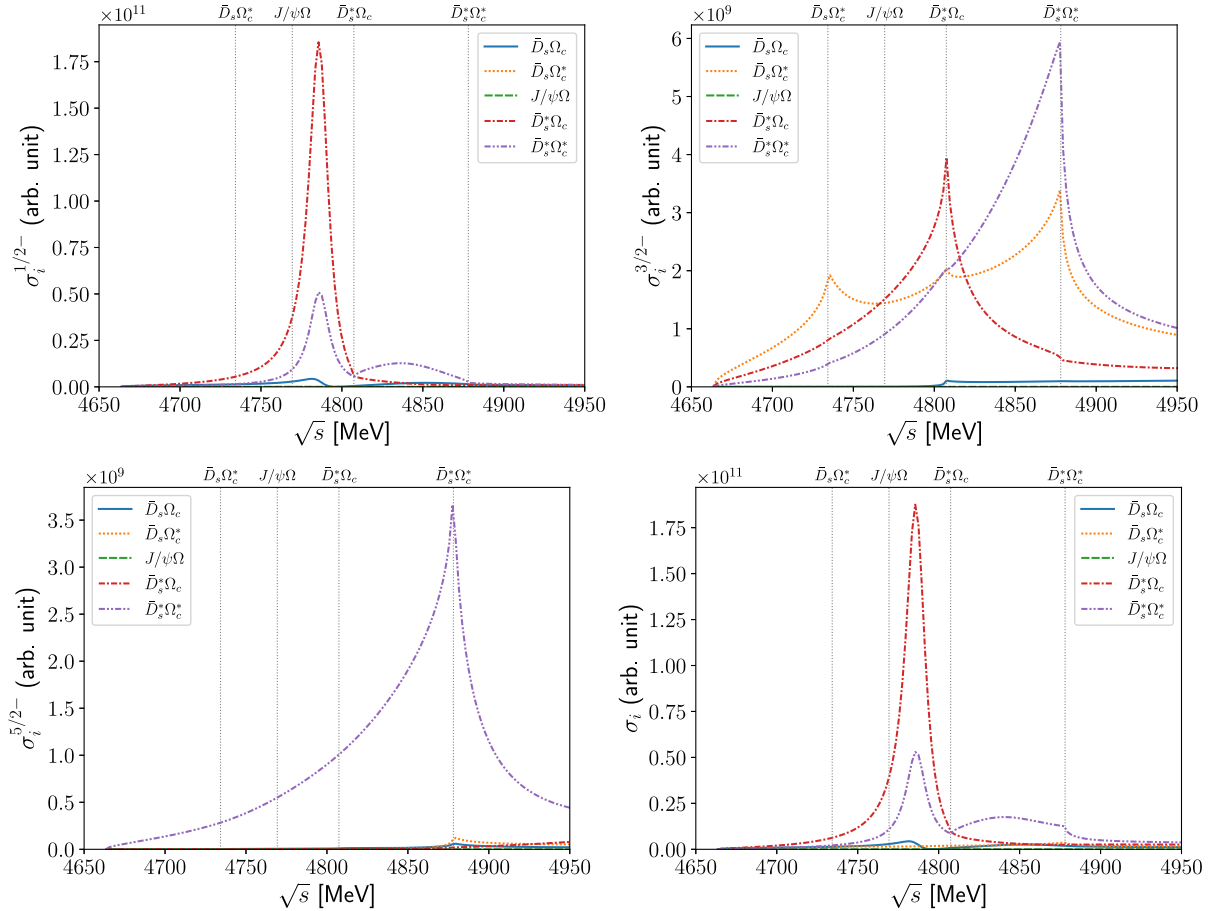


FIG. 4. Elastic partial-wave cross sections for $J^P = 1/2^-, 3/2^-, 5/2^-$ (upper left, upper right, and lower left panels) and total cross section (lower right panel) versus total energy. The subscript i labels the two-particle states shown in the legend.

observe a tiny effect from the $\bar{D}_s\Omega_c$ channel, and we do not see any signals in the $\bar{D}_s\Omega_c^*$ and $J/\psi\Omega$ channels. Note that the broad resonance structure is found in the $\bar{D}_s^*\Omega_c^*$ channel, which is related to the second resonance $P_{c\bar{c}sss}(4841)$ with $J^P = 1/2^-$. Interestingly, the $P_{c\bar{c}sss}(4841)$ state does not appear in other channels. We will discuss the reason later when presenting the results for the coupling strengths to the relevant channels. The right upper panel of Fig. 4 draws the results for the $\sigma_i^{3/2-}$. While we do not observe any resonances, we see three cusp structures at the thresholds corresponding to the $\bar{D}_s\Omega_c$, $\bar{D}_s^*\Omega_c$, and $\bar{D}_s^*\Omega_c^*$ channels. It is interesting to see that the $\bar{D}_s\Omega_c$ channel exhibits the three cusp structures even at the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ thresholds. In the case of the $J^P = 5/2^-$ partial wave, which is depicted in the lower left panel, the single cusp structure is revealed only at the $\bar{D}_s^*\Omega_c^*$ threshold. This is natural, since it is only possible to have the total spin $5/2$ from the $\bar{D}_s^*\Omega_c^*$ channel in S wave. Finally, we plot the total elastic cross section including the relevant partial waves with both negative and positive parities. We clearly observe the $P_{c\bar{c}sss}(4787)$ state in the $\bar{D}_s^*\Omega_c$ channel, and both the $P_{c\bar{c}sss}(4787)$ and $P_{c\bar{c}sss}(4841)$ in the $\bar{D}_s^*\Omega_c^*$ channel. In contrast to the cases of the $P_{c\bar{c}}$, $P_{c\bar{c}s}$, and $P_{c\bar{c}ss}$, we do not find any $P_{c\bar{c}sss}$ resonances with positive parity (P wave).

Figure 5 shows the contour plot of the modulus squared for the $\bar{D}_s^*\Omega_c^*$ transition amplitude, $|T_{\bar{D}_s^*\Omega_c^*}|^2$ in the complex total energy plane. We can easily find the pole positions for the two triple-strangeness hidden-charm pentaquarks states. The first resonance, $P_{c\bar{c}sss}(4787)$, is located at $(4786.9 - i7.2)$ MeV below the $\bar{D}_s^*\Omega_c$ threshold. It indicates that the width of $P_{c\bar{c}sss}(4787)$ is $\Gamma = 14.4$ MeV. The $P_{c\bar{c}sss}(4841)$ state is positioned at $(4840.5 - i43.1)$ MeV below the $\bar{D}_s^*\Omega_c^*$ threshold. It has a larger width compared to the first resonance, i.e., $\Gamma = 86.2$ MeV.

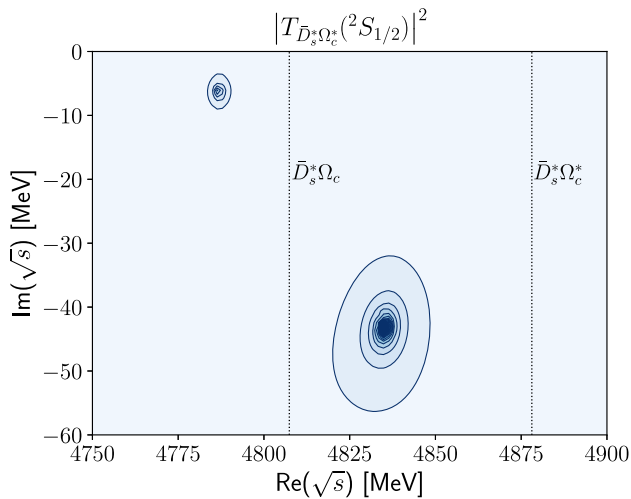


FIG. 5. Contour plot of the modulus squared for the $\bar{D}_s^*\Omega_c^*$ transition amplitude in the complex total energy plane. Two poles are observed, which correspond to the $P_{c\bar{c}sss}(4787)$ (upper left pole) and $P_{c\bar{c}sss}(4841)$ (lower right pole).

The coupling strength can be extracted from the residue of the partial-wave transition amplitude, which is expressed as

$$\mathcal{R}_{ab} = \lim_{s \rightarrow s_R} (s - s_R) \mathcal{T}_{ab} / 4\pi = g_a g_b. \quad (45)$$

We want emphasize that the definition of the coupling strength in Eq. (45) does not allow one to extract its absolute sign. To determine the relative signs, thus, we consider the real part of the coupling to the lowest threshold channel to be positive.

Table II lists the values of the coupling strengths for the $P_{c\bar{c}sss}(4787)$ and $P_{c\bar{c}sss}(4841)$ to all channels involved, in units of GeV. As shown in the upper left panel of Fig. 4, the hidden-charm pentaquark $P_{c\bar{c}sss}(4787)$ state appears in the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels. This indicates that these two channels are the most dominant ones. As expected, the coupling strength $g_{\bar{D}_s^*\Omega_c}(^2S_{1/2})$ dominates over all other coupling strengths, followed by $g_{\bar{D}_s^*\Omega_c^*}(^2S_{1/2})$. The mixing with the D wave is negligibly small. The coupling strength to the $\bar{D}_s\Omega_c$ channel is marginal compared with those to the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels, as already illustrated in the upper left panel of Fig. 4.

The $P_{c\bar{c}sss}(4841)$ is only found in the $\bar{D}_s^*\Omega_c^*$ channel. The corresponding coupling strength is $g_{\bar{D}_s^*\Omega_c^*}(^2S_{1/2}) = (20.41 + i3.50)$ GeV, the largest among all channels. Although it is not possible to conclude whether these two triple-strangeness hidden-charm resonances are genuine pentaquarks or molecular states solely from the coupling strengths, they do provide partial insight into the nature of these resonances. Examining the results in Table II, the $P_{c\bar{c}sss}(4787)$ is best interpreted as a mixed state of the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels, whereas the second resonance, $P_{c\bar{c}sss}(4841)$, is most likely a $\bar{D}_s^*\Omega_c^*$ molecular state.

Since several theoretical predictions for the pentaquark states with $S = -3$ already exist, it is worthwhile to compare

TABLE II. Coupling strengths of the two $P_{c\bar{c}sss}$'s with $J^P = 1/2^-$ to all channels involved.

	$P_{c\bar{c}sss}(4787)$	$P_{c\bar{c}sss}(4841)$
$\sqrt{s_R}$ (MeV)	4786.9 - i7.2	4840.5 - i43.1
$g_{\bar{D}_s\Omega_c}(^2S_{1/2})$	$3.40 + i2.26$	$6.71 - i3.58$
$g_{\bar{D}_s\Omega_c^*}(^4D_{1/2})$	$0.40 + i0.11$	$1.71 - i1.22$
$g_{\bar{D}_s^*\Omega_c}(^2S_{1/2})$	$-13.23 - i2.24$	$1.61 + i5.79$
$g_{\bar{D}_s^*\Omega_c}(^4D_{1/2})$	$-0.03 - i0.02$	$-1.13 - i0.43$
$g_{J/\psi\Omega}(^2S_{1/2})$	$-0.03 - i0.00$	$0.02 + i0.01$
$g_{J/\psi\Omega}(^4D_{1/2})$	$-0.00 + i0.00$	$0.00 + i0.00$
$g_{J/\psi\Omega}(^6D_{1/2})$	$0.00 - i0.00$	$0.00 + i0.01$
$g_{\bar{D}_s^*\Omega_c^*}(^2S_{1/2})$	$-10.19 + i2.86$	$20.41 + i3.50$
$g_{\bar{D}_s^*\Omega_c^*}(^4D_{1/2})$	$0.00 - i0.03$	$-0.07 - i0.07$
$g_{\bar{D}_s^*\Omega_c^*}(^6D_{1/2})$	$0.29 + i0.01$	$0.04 + i0.18$

them with the present results. Meng *et al.* [22] studied the $P_{c\bar{c}sss}$ states as compact pentaquarks within a quark potential model. They found seven pentaquark states with spin-parity $J^P = 1/2^-$. The mass of the lowest $1/2^-$ state in Ref. [22] is 4855 MeV, which is close to the second resonance $P_{c\bar{c}sss}(4841)$ observed in the present work. In addition, they predicted nine $P_{c\bar{c}sss}$ states with $J^P = 3/2^-$ and 15 different pentaquarks with $J^P = 5/2^-$. Even more pentaquark states with positive parity were also predicted in Ref. [22]. Wang *et al.* [23] constructed the hadronic potentials based on the one-boson-exchange model and four different coupled channels: $\bar{D}_s\Omega_c$, $\bar{D}_s\Omega_c^*$, $\bar{D}_s^*\Omega_c$, and $\bar{D}_s^*\Omega_c^*$. They then solved the coupled-channel Schrödinger equation in the presence of these hadronic potentials. They suggested two triple-strangeness hidden-charm pentaquark states with $J^P = 3/2^-$ and $5/2^-$ as bound states of the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels, rather than as resonances. To describe the decays of these pentaquark states, they additionally employed the quark interchange model. Roca *et al.* [20] also investigated the $P_{c\bar{c}sss}$ states. However, they concluded that the meson-baryon interaction they employed was insufficient to generate a bound or resonance state. In the present work, on the other hand, we have further considered the $\bar{D}_s\Omega_c$, $\bar{D}_s^*\Omega_c$, and $\bar{D}_s^*\Omega_c^*$ channels, which are absent in Ref. [20]. We have also included the scalar-isoscalar meson-exchange diagrams. It is essential to take these contributions into account in order to generate the $P_{c\bar{c}sss}$ states.

The $P_{c\bar{c}}$ and $P_{c\bar{c}s}$ were experimentally found to decay into the $J/\psi N$ and $J/\psi \Lambda$, respectively [1–3,5,6]. So, we anticipate that the hidden-charm pentaquark states with triple-strangeness may decay into J/ψ and Ω , if they are observed experimentally. Therefore, it is worthwhile to evaluate the transition cross sections from the $\bar{D}_s^{(*)}\Omega_c^{(*)}$ channels to the $J/\psi\Omega$ channel. Figure 6 shows the

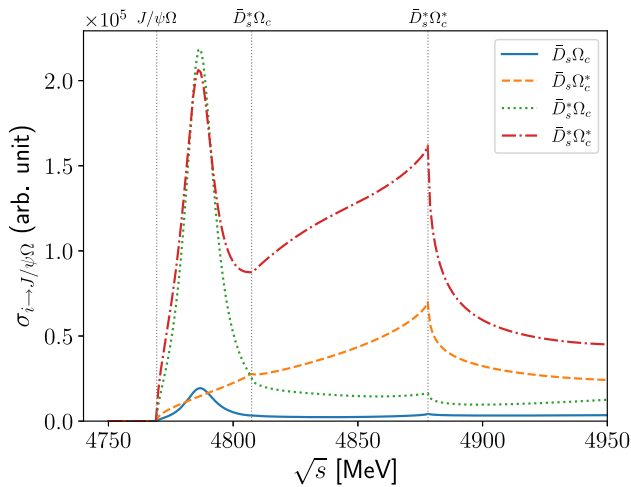


FIG. 6. Total cross section as a function of energy for the transitions from the $\bar{D}_s^{(*)}\Omega_c^{(*)}$ channels to the $J/\psi\Omega$ channel. The subscript i labels the two-particle states shown in the legend.

numerical results for the total transition cross sections of $\bar{D}_s^{(*)}\Omega_c^{(*)} \rightarrow J/\psi\Omega$. We can evidently observe the first resonance $P_{c\bar{c}sss}(4787)$ in the $J/\psi\Omega$ invariant mass spectrum. Interestingly, we can also see the resonance structure in the $\bar{D}_s\Omega_c$ channel, though its strength is weaker than in the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels. Since the $\bar{D}_s\Omega_c$ channel yields only the D -wave contribution to the $\bar{D}_s\Omega_c \rightarrow J/\psi\Omega$ interaction, we do not find any signals for the resonance. On the other hand, the second resonance $P_{c\bar{c}sss}(4841)$ is screened by the cusp structures and background contributions. So, we expect that only the first resonance could be observed experimentally.

IV. SUMMARY AND CONCLUSIONS

We have investigated the possible existence of the triple-strangeness hidden-charm pentaquark states in the off-shell coupled-channel formalism. We have considered the open-charm meson-baryon channels $\bar{D}_s\Omega_c$, $\bar{D}_s\Omega_c^*$, $\bar{D}_s^*\Omega_c$, and $\bar{D}_s^*\Omega_c^*$, as well as the hidden-charm channels $J/\psi\Omega$. We have constructed the two-body kernel Feynman amplitudes using the effective Lagrangian based on the hidden local symmetry, constrained by heavy-quark spin symmetry. We have solved the coupled Blankenbecler-Sugar equation in the partial-wave helicity basis and have searched for poles in the complex energy plane to identify possible resonance states.

We have found two triple-strangeness hidden-charm pentaquark states with $J^P = 1/2^-$: $P_{c\bar{c}sss}(4787)$ and $P_{c\bar{c}sss}(4841)$. The $P_{c\bar{c}sss}(4787)$ couples dominantly to the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels, while the $P_{c\bar{c}sss}(4841)$ couples almost exclusively to the $\bar{D}_s^*\Omega_c$ channel. We have also provided the coupling strengths of these states to all relevant channels. Examining the results for the coupling strengths, the $P_{c\bar{c}sss}(4787)$ is best interpreted as a mixed state of the $\bar{D}_s^*\Omega_c$ and $\bar{D}_s^*\Omega_c^*$ channels, whereas the second resonance, $P_{c\bar{c}sss}(4841)$, is most likely a $\bar{D}_s^*\Omega_c^*$ molecular state. Finally, we have evaluated the total transition cross sections of $\bar{D}_s^{(*)}\Omega_c^{(*)} \rightarrow J/\psi\Omega$. The first resonance $P_{c\bar{c}sss}(4787)$ was clearly observed in the $J/\psi\Omega$ invariant mass spectrum, whereas the second one $P_{c\bar{c}sss}(4841)$ was screened by the cusp structures and background contributions. Therefore, we expect that only the first resonance can be observed experimentally.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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