

Nonexistence of hidden-charm pentaquarks in J/ψ photoproduction

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We investigate J/ψ photoproduction off the proton, $\gamma p \rightarrow J/\psi p$, to elucidate the nonexistence of hidden-charm pentaquark signals reported by the GlueX and CLAS12 experiments. Within a coupled-channel rescattering mechanism, we employ the transition amplitudes from a previous coupled-channel analysis that dynamically generates the $P_{c\bar{c}}$ states. The kernel amplitudes for the transition to the $J/\psi N$ channel include both t -channel heavy-meson exchange and u -channel heavy-baryon exchange. We find that the rescattering contributions from the $\bar{D}^{(*)}\Sigma_c$ channels—indispensable for the formation of the $P_{c\bar{c}}$ resonances—are about one order of magnitude smaller than those from $\bar{D}^{(*)}\Lambda_c$, since $g_{\bar{D}^{(*)}N\Sigma_c}$ is roughly 5 times smaller than $g_{\bar{D}^{(*)}N\Lambda_c}$. Since the $P_{c\bar{c}}$ resonances couple to the $J/\psi N$ channel predominantly through the $\bar{D}^{(*)}\Sigma_c$ intermediate states, their suppression prevents the pentaquark signal from appearing in photoproduction. With only a single parameter controlling the overall normalization, the present work describes the GlueX and CLAS12 cross sections well. These results suggest that the null result from photoproduction need not be in conflict with the pentaquark signals observed by the LHCb Collaboration.

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I. INTRODUCTION

Since the LHCb Collaboration announced the existence of hidden-charm pentaquark baryons, a plethora of experimental and theoretical works have been performed. These states were first observed in the $J/\psi p$ invariant mass spectrum from $\Lambda_b^0 \rightarrow J/\psi p K^-$ decays [1,2]. Thus far, four of them have been established: three narrow states below the $\bar{D}^*\Sigma_c$ threshold and one broad state below the $\bar{D}\Sigma_c^*$ threshold. Notably, $P_{c\bar{c}}(4330)$ was observed in $B_s^0 \rightarrow J/\psi p \bar{p}$ decays, with no signal from the previously reported $P_{c\bar{c}}(4312)$ [3]. In contrast, the GlueX [4,5] and CLAS12 [6] Collaborations found no evidence for such states in J/ψ photoproduction off the nucleon. Given that the $P_{c\bar{c}}$ states decay almost exclusively into J/ψ and a proton, this null result is rather surprising. Several works subsequently tried to explain this absence [7–10], but they mainly reproduced

the experimental data without clarifying the underlying mechanism.

We start from a recent coupled-channel formalism for hidden-charm J/ψ and open-charm two-body processes [11], which provides fully off-shell transition amplitudes. Based on it, we show why hidden-charm pentaquarks cannot be generated in J/ψ photoproduction. There, seven peaks were found: six identified as $P_{c\bar{c}}$ resonances and one as a cusp. Four reproduce the known states $P_{c\bar{c}}(4312)$, $P_{c\bar{c}}(4380)$, $P_{c\bar{c}}(4440)$, and $P_{c\bar{c}}(4457)$, whereas two near 4.5 GeV, $P_{c\bar{c}}(4517, J^P = 3/2^-)$ and $P_{c\bar{c}}(4522, J^P = 5/2^-)$, are predicted as $\bar{D}^*\Sigma_c^*$ molecular states. Two further $J^P = 1/2^+$ resonances, at $(4401 - i35)$ and $(4533 - i17)$ MeV, lack a clean molecular interpretation and may correspond to genuine pentaquark configurations. A dynamical explanation for the missing signal was also proposed there through the $J/\psi N$ elastic channel, where the negative-parity peaks appear as dips and are washed out by the larger positive-parity (P -wave) bumps, while the resonances remain visible in the $\Lambda_b \rightarrow J/\psi p K^-$ type transitions seen by LHCb.

In this work, we aim at clarifying the theoretical origin of this suppression. To implement the transition amplitudes for J/ψ photoproduction, the u -channel diagrams must be taken into account. The crucial point is that the coupling constant at the $\bar{D}\Sigma_c N$ vertex is much smaller than that at the

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$\bar{D}\Lambda_c N$ one. In Ref. [11], the $\bar{D}\Sigma_c$ channel was shown to be essential for describing the resonances, in particular $P_{c\bar{c}}(4312)$. In J/ψ photoproduction, however, the smallness of the $\bar{D}\Sigma_c N$ coupling constant prevents the dynamical generation of the hidden-charm pentaquarks, which we demonstrate explicitly below.

The present paper is organized as follows. In Sec. II, we describe the general formalism for J/ψ photoproduction. We first construct the $\gamma N \rightarrow J/\psi N$ amplitudes from the effective Lagrangian and then combine them with the two-body hidden-charm and open-charm processes. Solving the coupled-channel Blankenbecler-Sugar-type rescattering equations, we obtain the numerical results. In Sec. III, we first present and compare the results for $\bar{D}^{(*)}\Lambda_c \rightarrow J/\psi N$ and $\bar{D}^{(*)}\Sigma_c \rightarrow J/\psi N$, where $\bar{D}^{(*)}$ generically denotes either $\bar{D}$ or $\bar{D}^*$. We then display the total cross section for J/ψ photoproduction off the proton, compare it with the GlueX and CLAS12 data, and show explicitly how the hidden-charm pentaquarks fail to emerge. The last section is devoted to a summary and conclusions.

II. FORMALISM

In this section, we explain the general formalism employed in the present work. Taking into account the rescattering effects that contain information on the dynamical generation of the hidden-charm pentaquark states (see Fig. 1), we derive the transition amplitudes for J/ψ photoproduction by solving the following coupled-channel BS-type rescattering equation:

$$T_{J/\psi p, \gamma p}(\mathbf{p}', \mathbf{p}) = V_{J/\psi p, \gamma p}(\mathbf{p}', \mathbf{p}) + \frac{1}{(2\pi)^3} \sum_k \int \frac{d^3 q}{2E_{k1}(\mathbf{q})E_{k2}(\mathbf{q})} \times V_{k, \gamma p}(\mathbf{q}, \mathbf{p}) \frac{E_k(\mathbf{q})}{s - E_k^2(\mathbf{q}) + i\epsilon} T_{J/\psi p, k}(\mathbf{p}', \mathbf{q}), \quad (1)$$

where $V_{k, \gamma p}$ denotes the two-body Feynman kernel amplitude for the intermediate state k in the rescattering equation.

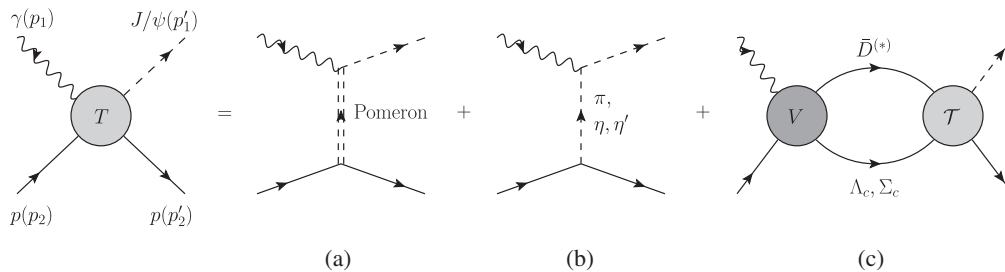


FIG. 1. Diagrams for J/ψ photoproduction off the proton ($\gamma p \rightarrow J/\psi p$): (a) Pomeron exchange, (b) the Born diagrams, and (c) the rescattering diagram.

When $k = J/\psi p$, $V_{J/\psi p, \gamma p}$ represents the $\gamma p \rightarrow J/\psi p$ Born diagrams [12]. $T_{J/\psi p, k}$ stands for the transition amplitude for the $J/\psi p \rightarrow k$ process, where k runs over the intermediate two-body open-charm states: $\bar{D}\Lambda_c$, $\bar{D}^*\Lambda_c$, $\bar{D}\Sigma_c$, $\bar{D}\Sigma_c^*$, $\bar{D}^*\Sigma_c$, and $\bar{D}^*\Sigma_c^*$. Since the elastic amplitude $T_{J/\psi p, J/\psi p}$ is much smaller than those involving open-charm intermediate channels, it is excluded in the present work. The $T_{J/\psi p, k}$ amplitudes were already constructed in Ref. [11]. Here, $\mathbf{p}$ and $\mathbf{p}'$ are the relative three-momenta of the initial and final states, respectively, while $\mathbf{q}$ denotes the three-momentum of the intermediate state in the center-of-mass (c.m.) frame. The variable s is the square of the total c.m. energy, and $E_k = E_{k1} + E_{k2}$ is the total on-mass-shell energy of the intermediate state.

We perform the partial-wave expansion of Eq. (1) in the helicity basis and obtain the partial-wave transition amplitudes as follows:

$$T_{\lambda'\lambda}^J(\mathbf{p}', \mathbf{p}) = V_{\lambda'\lambda}^J(\mathbf{p}', \mathbf{p}) + \frac{1}{(2\pi)^3} \sum_{k, \lambda_k} \int \frac{q^2 dq}{2E_{k1}E_{k2}} V_{\lambda_k \lambda}^J(\mathbf{q}, \mathbf{p}) \times \frac{E_k}{s - E_k^2 + i\epsilon} T_{\lambda' \lambda_k}^J(\mathbf{p}', \mathbf{q}), \quad (2)$$

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 channel indices are dropped to simplify the notation. 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 of the kernel amplitude $V_{\lambda'\lambda}^J$ is given by

$$V_{\lambda'\lambda}^J(\mathbf{p}', \mathbf{p}) = 2\pi \int d(\cos \theta) d_{\lambda_1 - \lambda_2, \lambda'_1 - \lambda'_2}^J(\theta) V_{\lambda'\lambda}^J(\mathbf{p}', \mathbf{p}, \theta), \quad (3)$$

where θ is the scattering angle and $d_{\lambda'\lambda}^J(\theta)$ denotes the reduced Wigner D functions. Since the propagator in Eq. (2) develops a singularity when the off-shell energy coincides with the on-shell one, we regularize the rescattering equation by isolating the singularity and treating it separately. Accordingly, the rescattering equation is decomposed into regularized and singular parts as follows:

$$\begin{aligned}
T_{\lambda'\lambda}^J(\mathbf{p}', \mathbf{p}) &= V_{\lambda'\lambda}^J(\mathbf{p}', \mathbf{p}) + \frac{1}{(2\pi)^3} \\
&\times \sum_{k,\lambda_k} \left[\int_0^\infty dq \frac{q E_k}{E_{k1} E_{k2}} \frac{\mathcal{F}(q) - \mathcal{F}(\tilde{q}_k)}{s - E_k^2} + \frac{1}{2\sqrt{s}} \right. \\
&\times \left. \left(\ln \left| \frac{\sqrt{s} - E_k^{\text{thr}}}{\sqrt{s} + E_k^{\text{thr}}} \right| - i\pi \right) \mathcal{F}(\tilde{q}_k) \right], \quad (4)
\end{aligned}$$

where

$$\mathcal{F}(q) = \frac{1}{2} q V_{\lambda_k \lambda}^J(q, \mathbf{p}) \mathcal{T}_{\lambda' \lambda_k}^J(\mathbf{p}', q). \quad (5)$$

Here, $\tilde{q}_k$ denotes the on-shell momentum defined by the condition $E_{k1}(\tilde{q}_k) + E_{k2}(\tilde{q}_k) = \sqrt{s}$. This regularization procedure is applied only when the total energy $\sqrt{s}$ exceeds the threshold energy E_k^{thr} of the k th channel.

In the previous study [11], the transition amplitudes $\mathcal{T}_{J/\psi N, k}$ were computed by including only the t -channel diagrams in the kernel amplitudes. To describe J/ψ photo-production quantitatively, we also need to introduce the u -channel diagrams involving heavy-baryon exchange, which provide significant contributions to the transition amplitudes. Since the $J/\psi N$ elastic scattering amplitude is

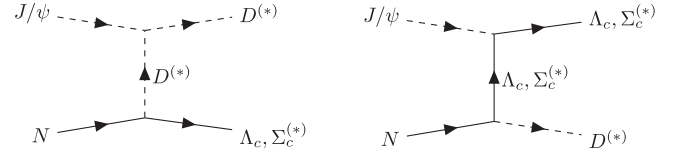


FIG. 2. The t -channel (left panel) and u -channel (right panel) diagrams.

small and the transition amplitudes to the $J/\psi N$ channel are governed by heavy-hadron exchange, the resonance properties remain almost intact despite the inclusion of the u -channel diagrams. Figure 2 illustrates the t -channel meson-exchange and u -channel baryon-exchange diagrams for the transition to the $J/\psi N$ channel. The exchanged particles considered in each reaction are summarized in Table I. A kernel amplitude associated with a given exchange diagram is written as

$$\mathcal{V} = C_I F^2(q^2) \Gamma_1 \mathcal{P}(q) \Gamma_2, \quad (6)$$

where C_I denotes the isospin factor for each exchanged particle in a given channel, as listed in Table I. The vertex functions $\Gamma_{1,2}$ are derived from the following effective Lagrangians respecting heavy-quark spin symmetry:

$$\begin{aligned}
\mathcal{L}_{DDJ/\psi} &= -ig_\psi M_D \sqrt{M_J} \left(J/\psi^\mu D^\dagger \overleftrightarrow{\partial}_\mu D \right), \\
\mathcal{L}_{D^*DJ/\psi} &= ig_\psi \sqrt{\frac{M_D M_{D^*}}{M_J}} \epsilon^{\mu\nu\alpha\beta} \partial_\mu J/\psi_\nu \left(D^\dagger \overleftrightarrow{\partial}_\alpha D_\beta^* + D_\beta^{*\dagger} \overleftrightarrow{\partial}_\alpha D \right), \\
\mathcal{L}_{D^*D^*J/\psi} &= ig_\psi M_{D^*} \sqrt{M_J} (g^{\mu\nu} g^{\alpha\beta} - g^{\mu\alpha} g^{\nu\beta} + g^{\mu\beta} g^{\nu\alpha}) \left(J/\psi_\mu D_\nu^{*\dagger} \overleftrightarrow{\partial}_\alpha D_\beta^* \right), \\
\mathcal{L}_{DN\Lambda_c} &= -ig_{I\bar{3}} \sqrt{M_D} \bar{N} \gamma_5 \Lambda_c \bar{D} + \text{H.c.}, \\
\mathcal{L}_{D^*N\Lambda_c} &= -ig_{I\bar{3}} \sqrt{M_{D^*}} \bar{N} \gamma^\mu \Lambda_c \bar{D}_\mu^* + \text{H.c.}, \\
\mathcal{L}_{DN\Sigma_c} &= ig_{I6} \sqrt{\frac{3M_D}{2}} \bar{N} \gamma_5 \boldsymbol{\tau} \cdot \boldsymbol{\Sigma}_c \bar{D} + \text{H.c.}, \\
\mathcal{L}_{D^*N\Sigma_c} &= -ig_{I6} \sqrt{\frac{M_{D^*}}{6}} \bar{N} \gamma^\mu \boldsymbol{\tau} \cdot \boldsymbol{\Sigma}_c \bar{D}_\mu^* + \text{H.c.}, \\
\mathcal{L}_{D^*N\Sigma_c^*} &= ig_{I6} \sqrt{2M_{D^*}} \bar{N} \gamma_5 \boldsymbol{\tau} \cdot \boldsymbol{\Sigma}_c^{*\mu} \bar{D}_\mu^* + \text{H.c.}, \\
\mathcal{L}_{J/\psi \Lambda_c \Lambda_c} &= -g_{J\bar{3}} \bar{\Lambda}_c \gamma^\mu \Lambda_c J/\psi_\mu, \\
\mathcal{L}_{J/\psi \Sigma_c \Sigma_c} &= -g_{J6} \bar{\boldsymbol{\Sigma}}_c \cdot \gamma^\mu \boldsymbol{\Sigma}_c J/\psi_\mu, \\
\mathcal{L}_{J/\psi \Sigma_c \Sigma_c^*} &= \frac{2g_{J6}}{\sqrt{3}} \bar{\boldsymbol{\Sigma}}_c \cdot \gamma_5 \boldsymbol{\Sigma}_c^{*\mu} J/\psi_\mu + \text{H.c.}, \\
\mathcal{L}_{J/\psi \Sigma_c^* \Sigma_c^*} &= g_{J6} \bar{\boldsymbol{\Sigma}}_{c\mu}^* \cdot \gamma^\mu \boldsymbol{\Sigma}_c^{*\mu} J/\psi_\mu. \quad (7)
\end{aligned}$$

Since no reliable theoretical information is available for the relevant coupling constants, we estimate them using SU(4) symmetry relations, even though this symmetry is strongly broken. To quantify the uncertainty arising from SU(4)

TABLE I. Isospin factors (C_I) for the exchange diagrams contributing to the transitions from the $J/\psi N$ channel to each open-charm meson-baryon channel, together with the corresponding exchanged particles.

Reactions	Exchange	C_I
$J/\psi N \rightarrow \bar{D}\Lambda_c$	$\bar{D}, \bar{D}^*, \Lambda_c$	1
$J/\psi N \rightarrow \bar{D}^*\Lambda_c$	$\bar{D}, \bar{D}^*, \Lambda_c$	1
$J/\psi N \rightarrow \bar{D}\Sigma_c$	$\bar{D}, \bar{D}^*, \Sigma_c$	$\sqrt{3}$
$J/\psi N \rightarrow \bar{D}\Sigma_c^*$	$\bar{D}^*, \Sigma_c$	$\sqrt{3}$
$J/\psi N \rightarrow \bar{D}^*\Sigma_c$	$\bar{D}, \bar{D}^*, \Sigma_c, \Sigma_c^*$	$\sqrt{3}$
$J/\psi N \rightarrow \bar{D}^*\Sigma_c^*$	$\bar{D}^*, \Sigma_c, \Sigma_c^*$	$\sqrt{3}$

symmetry, we introduce a common prefactor a by which all couplings derived from the SU(4) relations are multiplied. The couplings $g_{I\bar{3}}$ and g_{I6} are then estimated as

$$g_{I\bar{3}}\sqrt{M_D} = -\frac{3\sqrt{3}}{5}ag_{\pi NN},$$

$$g_{I6}\sqrt{\frac{3M_D}{2}} = -\frac{ag_{\pi NN}}{5}. \quad (8)$$

Similarly, $g_{J\bar{3}}$ and g_{J6} are obtained from $g_{\rho NN}$ via the SU(4) relation:

$$g_{J\bar{3}} = g_{J6} = \sqrt{2}ag_{\rho NN}. \quad (9)$$

In the present work, we adopt the coupling constants from the Nijmegen potential [13,14], namely $g_{\pi NN} = 13.2$ and $g_{\rho NN} = 2.97$.

For the mesonic sector, we follow the same procedure and estimate the g_ψ coupling as

$$g_\psi M_D \sqrt{M_J} = \frac{\sqrt{2}}{2}ag_{\pi\pi\rho}. \quad (10)$$

The value $g_{\pi\pi\rho} = 5.97$ is taken from our previous work [15]. We find that the parameter a controls the overall magnitude of the total J/ψ photoproduction cross section without significantly affecting the hidden-charm pentaquark resonances. A value of $a = 0.47$ provides the best description of the experimental data, whereas $a = 1$ yields a cross section roughly an order of magnitude larger. Thus, the parameter a is the main source for the uncertainty of the present work.

The propagators for pseudoscalar and vector meson exchange are given by [16]

$$\mathcal{P}(q, M) = \frac{1}{q^2 - M^2},$$

$$\mathcal{P}_{\mu\nu}(q, M) = \frac{1}{q^2 - M^2} \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{M^2} \right), \quad (11)$$

whereas those for spin- $\frac{1}{2}$ and spin- $\frac{3}{2}$ baryon exchange are expressed as

$$\mathcal{P}(q, M) = \frac{\not{q} + M}{q^2 - M^2},$$

$$\mathcal{P}_{\mu\nu}(q, M) = \frac{\not{q} + M}{q^2 - M^2} \left(-g_{\mu\nu} + \frac{1}{3}\gamma_\mu\gamma_\nu + \frac{1}{3M}(\gamma_\mu q_\nu - \gamma_\nu q_\mu) + \frac{2}{3M^2}q_\mu q_\nu \right). \quad (12)$$

Here, q denotes the exchanged four-momentum and M the mass of the exchanged particle.

Since hadrons have finite sizes, a form factor is introduced at each vertex. This form factor is also necessary to ensure the unitarity of the transition amplitudes. We use the following form [17]:

$$F_n(q^2) = \left(\frac{n\Lambda^2 - M^2}{n\Lambda^2 - q^2} \right)^n, \quad (13)$$

where the positive integer n is chosen such that the power of momentum in the vertex function Γ is regularized. As $n \rightarrow \infty$, the form factor in Eq. (13) approaches a Gaussian form. Note that the energy dependence in Eq. (13) must be removed, as it leads to unphysical behavior in the kernel amplitude. Although the values of the cutoff masses Λ in Eq. (13) cannot be determined directly from experiment or theory, the associated uncertainties can be reduced by exploiting the physical properties of the hadrons involved. Since heavy hadrons are more compact than lighter ones [18,19], higher cutoff masses are appropriate for them, as the cutoff mass is inversely proportional to the size of the corresponding hadron. We therefore introduce the reduced cutoff mass $\Lambda_0 := \Lambda - M$. In Ref. [20], the cutoff mass was shown to be related to Λ_{QCD} via $\Lambda = M + \eta\Lambda_{\text{QCD}}$, where η is of order unity. This approach has been successfully applied in previous studies [11,15,21–26]. In this work, all values of the reduced cutoff mass are fixed at $\Lambda_0 = 600$ MeV. We emphasize that we do not change the values of Λ_0 for fitting. We find that varying Λ_0 only rescales the overall magnitude of the rescattering contribution, without altering its shape—the same role played by the free parameter a .

Having revised the kernel amplitudes for the transition to the $J/\psi N$ channel, we recalculate the transition amplitudes $\mathcal{T}_{J/\psi N, k}$ within the same framework as in Ref. [11]. The resulting $J/\psi N$ transition amplitudes are discussed in detail below. Throughout the calculation, isospin symmetry is assumed for the hadron masses, and this approximation is retained exclusively for the masses in the remainder of the calculation. The kernel amplitudes $V_{k, \gamma p}$ in Eq. (1) receive contributions from the three tree-level diagrams depicted in Fig. 3: (a) $D^{(*)}$ exchange in the t channel, (b) nucleon exchange in the s channel, and (c) Λ_c or Σ_c exchange in the

u channel. All possible exchange diagrams for each reaction channel are listed in Table II. The corresponding vertex functions are derived from the effective Lagrangians in Eq. (7) together with the following electromagnetic interaction Lagrangians:

$$\begin{aligned}
 \mathcal{L}_{\gamma DD} &= -ie_{\bar{D}}A_\mu(\bar{D}\partial^\mu D - \partial^\mu \bar{D}D), \\
 \mathcal{L}_{\gamma DD^*} &= ig_{\gamma DD^*}\epsilon^{\mu\alpha\beta}\partial_\mu A_\nu\left(\bar{D}\overset{\leftrightarrow}{\partial}_\alpha D_\beta^* + \bar{D}_\beta^*\overset{\leftrightarrow}{\partial}_\alpha D\right), \\
 \mathcal{L}_{\gamma D^*D^*} &= ie_{\bar{D}^*}(g^{\mu\nu}g^{\alpha\beta} - g^{\mu\alpha}g^{\nu\beta} + g^{\mu\beta}g^{\nu\alpha})A_\mu\left(\bar{D}_\nu^*\partial_\alpha D_\beta^* - \partial_\alpha \bar{D}_\nu^*D_\beta^*\right), \\
 \mathcal{L}_{\gamma NN} &= -e_N\bar{N}\left[\gamma^\mu - \frac{\kappa_N}{2M_N}\sigma^{\mu\nu}\partial_\nu\right]A_\mu N, \\
 \mathcal{L}_{\gamma\Lambda_c\Lambda_c} &= -e_{\Lambda_c}\bar{\Lambda}_c\left[\gamma^\mu - \frac{\kappa_{\Lambda_c}}{2M_{\Lambda_c}}\sigma^{\mu\nu}\partial_\nu\right]A_\mu\Lambda_c, \\
 \mathcal{L}_{\gamma\Sigma_c\Sigma_c} &= -e_{\Sigma_c}\bar{\Sigma}_c\left[\gamma^\mu - \frac{\kappa_{\Sigma_c}}{2M_{\Sigma_c}}\sigma^{\mu\nu}\partial_\nu\right]A_\mu\Sigma_c.
 \end{aligned} \tag{14}$$

The relevant coupling constants are determined as

$$\begin{aligned}
 g_{\gamma D^0 D^{*0}} &= 0.30 \text{ GeV}^{-1}, & g_{\gamma D^- D^{*-}} &= 0.141 \text{ GeV}^{-1}, \\
 \kappa_p &= 1.79, & \kappa_{\Lambda_c} &= -0.025, \\
 \kappa_{\Sigma_c^+} &= 0.41, & \kappa_{\Sigma_c^{++}} &= 4.64,
 \end{aligned} \tag{15}$$

where $g_{\gamma DD^*}$ is extracted from the branching ratio $\text{Br}(D^* \rightarrow D\gamma)$ and the total decay width Γ_{D^*} [27,28] via the relation

$$\Gamma_{D^* \rightarrow D\gamma} = \frac{q_\gamma^3}{3\pi} (g_{\gamma DD^*})^2, \tag{16}$$

with $q_\gamma = (M_{D^*}^2 - M_D^2)/(2M_{D^*})$. The anomalous magnetic moment of the proton is taken from Ref. [27]. The anomalous magnetic moment of Y_c (denoting either Λ_c or Σ_c) is related to its magnetic moment μ_{Y_c} [29] by

$$\mu_{Y_c} = \frac{eY_c + e\kappa_{Y_c}}{2M_{Y_c}}. \tag{17}$$

The transition amplitudes are formulated to ensure gauge invariance, which is why we include the proton-pole

diagram despite its small contribution. Notably, the kernel amplitudes are governed by distinct coupling factors: $V_{\bar{D}^{(*)}\Lambda_c\gamma p}$ is proportional to $g_{\bar{D}^{(*)}N\Lambda_c}$, while $V_{\bar{D}^{(*)}\Sigma_c\gamma p}$ is proportional to $g_{\bar{D}^{(*)}N\Sigma_c}$. This distinction is a crucial feature, as it explains the relative suppression of the $\bar{D}^{(*)}\Sigma_c$ rescattering contribution.

In the present study, we exclude the $\bar{D}^{(*)}\Sigma_c^*$ intermediate states from the rescattering calculation for two primary reasons. Firstly, the coupling constant $g_{\bar{D}^{(*)}N\Sigma_c^*}$ is approximately 5 times smaller than $g_{\bar{D}^{(*)}N\Lambda_c}$, so that contributions from the $\bar{D}^{(*)}\Sigma_c^{(*)}$ intermediate states are significantly suppressed relative to those from $\bar{D}^{(*)}\Lambda_c$. Secondly, constructing gauge-invariant amplitudes that involve a spin- $\frac{3}{2}$ exchange propagator presents considerable technical difficulties. For these reasons, the contributions from the $\bar{D}^{(*)}\Sigma_c^*$ channels are omitted.

Having formulated the kernel amplitude $V_{k,\gamma p}$, we proceed to evaluate the rescattering amplitude. Since the transition amplitude for the $J/\psi N$ channel was computed in the isospin basis, the corresponding amplitudes for definite charge states are related by

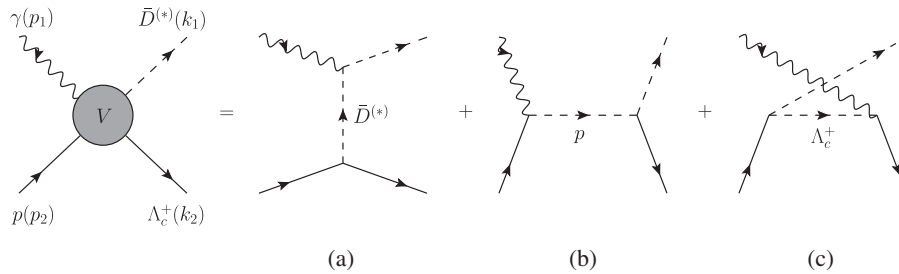


FIG. 3. Tree-level diagrams for the subprocess $\gamma p \rightarrow \bar{D}^{(*)}\Lambda_c(\Sigma_c)$ contributing to the rescattering amplitude: (a) $\bar{D}^{(*)}$ -meson exchange in the t channel, (b) nucleon exchange in the s channel, and (c) $\Lambda_c(\Sigma_c)$ baryon exchange in the u channel.

TABLE II. Exchange diagrams contributing to each charmed meson–baryon photoproduction channel off the proton.

Reaction	Exchanged particles
$\gamma p \rightarrow \bar{D}^0 \Lambda_c^+$	$p, \bar{D}^{*0}, \Lambda_c^+$
$\gamma p \rightarrow \bar{D}^{*0} \Lambda_c^+$	$p, \bar{D}^0, \Lambda_c^+$
$\gamma p \rightarrow \bar{D}^0 \Sigma_c^+$	$p, \bar{D}^{*0}, \Sigma_c^+$
$\gamma p \rightarrow D^- \Sigma_c^{++}$	$p, D^{*-}, \bar{D}^{*0}, \Sigma_c^{++}$
$\gamma p \rightarrow \bar{D}^{*0} \Sigma_c^+$	$p, \bar{D}^0, \Sigma_c^+$
$\gamma p \rightarrow D^{*-} \Sigma_c^{++}$	$p, D^{*-}, \bar{D}^0, \Sigma_c^{++}$

$$\begin{aligned} \mathcal{T}_{J/\psi p, \bar{D}^0 \Lambda_c^+} &= \mathcal{T}_{J/\psi N, \bar{D} \Lambda_c}, & \mathcal{T}_{J/\psi p, \bar{D}^0 \Sigma_c^+} &= \frac{1}{\sqrt{3}} \mathcal{T}_{J/\psi N, \bar{D} \Sigma_c}, \\ \mathcal{T}_{J/\psi p, D^- \Sigma_c^{++}} &= -\sqrt{\frac{2}{3}} \mathcal{T}_{J/\psi N, \bar{D} \Sigma_c}. \end{aligned} \quad (18)$$

For the background contribution, we employ the Pomeron-exchange amplitude of Refs. [30–32], which has been shown to describe successfully the total cross-section data for diffractive vector-meson photoproduction in the high-energy regime ($E_\gamma \geq 10$ GeV). Since this parametrization was originally developed for high-energy diffractive processes, its use in the near-threshold region, $W = 4.0\text{--}4.7$ GeV, should be regarded as a phenomenological extrapolation rather than a microscopic low-energy description. Alternative approaches to the threshold region, based on gluon generalized parton distributions (GPDs) [33] and on holographic QCD [34], have also been proposed. In the present work, the low-energy structures of primary interest are described explicitly by the open-charm coupled-channel rescattering mechanism. The total cross section for J/ψ photoproduction off the proton is then obtained by combining the rescattering and Pomeron-exchange amplitudes.

III. NUMERICAL RESULTS

The null result for the hidden-charm pentaquark signal in J/ψ photoproduction has so far remained unexplained. Previous attempts employed the meson-exchange picture combined with Pomeron contributions, which reproduce the data reasonably well, especially when the two $\bar{D}^{(*)} \Lambda_c$ channels are included through a rescattering process that generates threshold effects at their respective thresholds [7]. These studies, however, do not include the $P_{c\bar{c}}$ states and therefore cannot explain how the signal disappears. In the present work, we investigate J/ψ photoproduction off the proton using the transition amplitudes obtained in Ref. [11], which contain dynamically generated hidden-charm pentaquark states consistent with the candidates observed in Λ_b decays by the LHCb Collaboration.

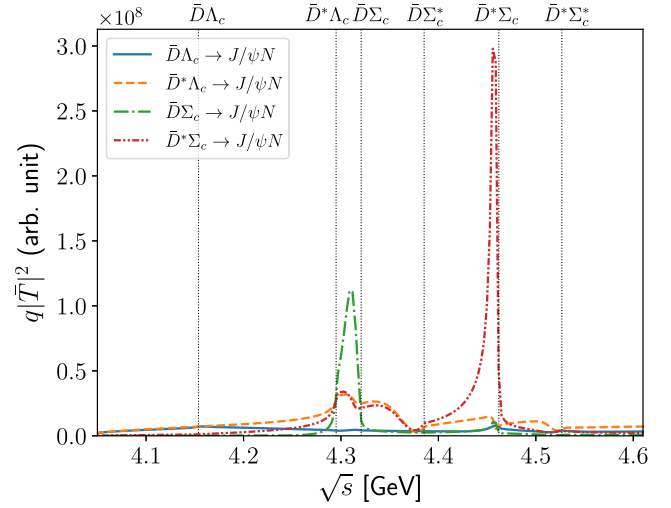


FIG. 4. Spin-averaged squared transition amplitude to the $J/\psi N$ channel, multiplied by the final-state center-of-mass momentum $q|\bar{T}|^2$, as a function of the total energy $\sqrt{s}$. The vertical dashed lines indicate the relevant meson-baryon thresholds.

Before evaluating the J/ψ photoproduction results, we discuss the $J/\psi N$ transition amplitudes. Figure 4 shows the distributions of the transition amplitudes from the $\bar{D} \Lambda_c$, $\bar{D}^* \Lambda_c$, $\bar{D} \Sigma_c$, and $\bar{D}^* \Sigma_c$ channels to the $J/\psi N$ state, respectively. We restrict ourselves to these four transitions, since the $\bar{D}^{(*)} \Sigma_c^*$ intermediate states were excluded from the rescattering contribution to the J/ψ photoproduction amplitudes, as justified in the previous section. The $P_{c\bar{c}}$ states appear prominently in the $\bar{D} \Sigma_c \rightarrow J/\psi N$ and $\bar{D}^* \Sigma_c \rightarrow J/\psi N$ transitions, with magnitudes exceeding those of the $\bar{D}^{(*)} \Lambda_c$ channels. This qualitatively implies that the resonance structure should be clearly observable unless the $\bar{D}^{(*)} \Sigma_c$ channels are significantly suppressed.

As is evident from the structure of the kernel amplitude $V_{k,\gamma p}$, the rescattering contributions from the $\bar{D}^{(*)} \Sigma_c$ intermediate states are suppressed. To demonstrate this explicitly, Fig. 5 compares the total cross section from the rescattering part, evaluated for each intermediate channel. The contributions from the $\bar{D} \Sigma_c$ and $\bar{D}^* \Sigma_c$ channels are approximately one order of magnitude smaller than those from the $\bar{D} \Lambda_c$ and $\bar{D}^* \Lambda_c$ channels. This illustrates why the $P_{c\bar{c}}$ states are suppressed in the photoproduction process: the $\bar{D}^{(*)} \Sigma_c$ channels couple only weakly to γN .

Furthermore, a close inspection of the small peak structure arising from $\bar{D}^* \Lambda_c$ rescattering reveals a weak signal of $P_{c\bar{c}}(4312)$. However, interference from the $\bar{D} \Sigma_c$ channel acts destructively, particularly in the $P_{c\bar{c}}(4312)$ peak region. Consequently, no visible peak

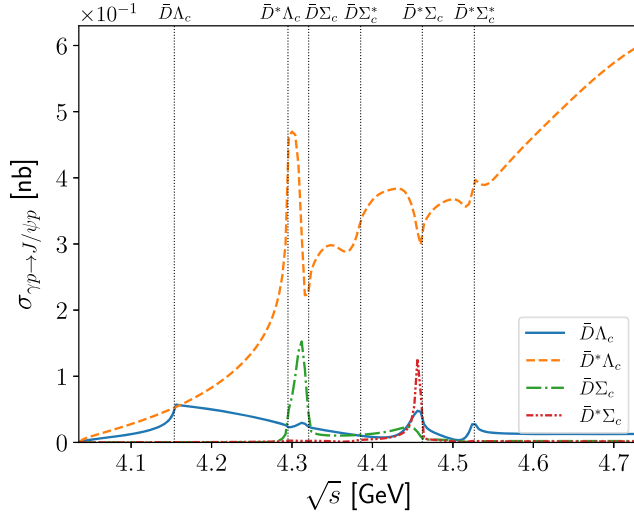


FIG. 5. Total cross section for the $\gamma p \rightarrow J/\psi p$ reaction from the rescattering contribution, evaluated separately for each intermediate state as a function of the total energy $\sqrt{s}$. The vertical dashed lines indicate the relevant meson-baryon thresholds.

remains in the total rescattering contribution to J/ψ photoproduction, as indicated by the dashed line in Fig. 6.

Having discussed the rescattering mechanism, we now examine the total cross section for J/ψ photoproduction off the proton as a function of energy, comparing the results with the experimental data in Fig. 6. The present model is constructed primarily from the Pomeron-exchange and rescattering contributions, since the Born term has a nearly negligible effect. The results indicate that the sizable cross section for J/ψ production originates from the rescattering effect, which facilitates processes that would otherwise be OZI-suppressed. The model reproduces the experimental data from both the GlueX and CLAS12 Collaborations remarkably well. As discussed above, the present model supports the nonexistence of hidden-charm pentaquarks in J/ψ photoproduction, owing to the suppression of the $\bar{D}^{(*)}\Sigma_c$ rescattering contribution, as mentioned already. The shaded area in Fig. 6 shows the uncertainty from varying the parameter a in the range 0.4–0.5. As noted above, a affects only the overall magnitude of the rescattering contribution. Even within this band, a $P_{c\bar{c}}(4312)$ signal cannot be identified.

While the GlueX experiment shows no signal for the hidden-charm pentaquark states, it exhibits cusp structures at the $\bar{D}^*\Lambda_c$, $\bar{D}\Sigma_c$, and $\bar{D}^*\Sigma_c^*$ thresholds [5,7]. These cusp structures arise mainly from the coupled-channel rescattering effects. In particular, we find that they are enhanced by the inclusion of the baryon-exchange

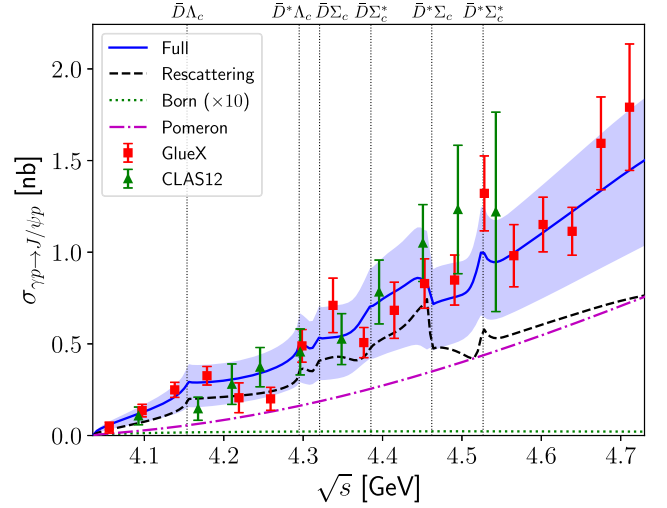


FIG. 6. Total cross section as a function of the total energy. Experimental data are taken from Refs. [5] (squares) and [6] (triangles). The error bars represent the statistical and systematic uncertainties added in quadrature. The vertical dashed lines indicate the relevant meson-baryon thresholds. The shaded area shows the uncertainty from varying the parameter a in the range 0.4–0.5.

diagrams, which were not considered in our previous work [11]. Baryon exchange is therefore essential in describing the cusp structures. As shown in Fig. 6, the present results for the cusp structures are consistent with the experimental data.

Here we give a brief comment on the effect of including the $\bar{D}^{(*)}\Sigma_c^*$ channels in the total cross section shown in Fig. 6. We estimate the order of magnitude of their contribution to be the same as, or smaller than, that of the $\bar{D}^{(*)}\Sigma_c$ channels already included. This means their effect would be visible mainly in the energy regions where the corresponding resonances are located. The $\bar{D}\Sigma_c^*$ contribution would appear in the region around 4.37 GeV, where the present coupled-channel analysis [11] dynamically generates the $P_{c\bar{c}}(4380)$ state. The $\bar{D}^*\Sigma_c^*$ one would be concentrated just below its threshold, near which the same analysis places the relevant states.

In Fig. 7, we present the numerical results for the differential cross section $d\sigma/dt$ as a function of $-t$ in the range of the average photon energy $8.93 \text{ GeV} \leq \langle E_\gamma \rangle \leq 10.82 \text{ GeV}$. The data are reasonably well described over the full range of $-t$ shown, keeping in mind that the experimental uncertainties also grow with $-t$. The uncertainty band from varying a , shown as the shaded area in Fig. 7, mainly affects the high $-t$ region. This is expected, since in the low $-t$ region the dominant contribution is the Pomeron-exchange term, which by itself explains the data in that region.

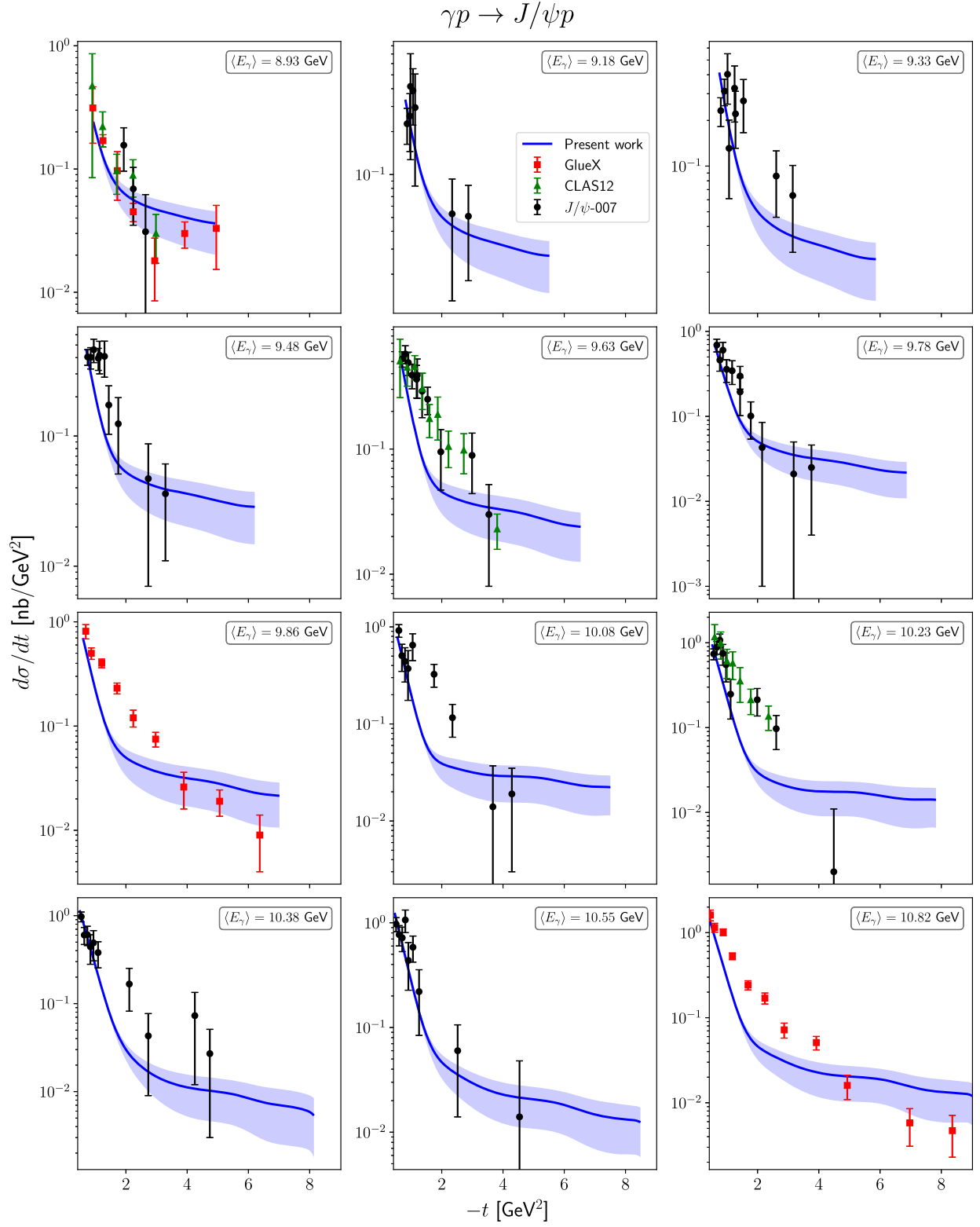


FIG. 7. Differential cross section as a function of the exchange momentum transfer $-t$. Experimental data are taken from Refs. [5] (squares), [6] (triangles), and [35] (circles). The error bars represent the statistical and systematic uncertainties added in quadrature. The shaded area shows the uncertainty from varying the parameter a in the range 0.4–0.5.

IV. SUMMARY AND CONCLUSIONS

In this work, we have investigated J/ψ photoproduction off the proton, $\gamma p \rightarrow J/\psi p$, to clarify why no hidden-charm pentaquark signals appear in the GlueX and CLAS12 experiments. The present framework combines a Pomeron-exchange background with a coupled-channel rescattering mechanism. The transition amplitudes $\mathcal{T}_{J/\psi N, k}$ in the rescattering kernel were taken from our coupled-channel analysis, which dynamically generates hidden-charm pentaquark states consistent with the $P_{c\bar{c}}$ observed by the LHCb Collaboration. We have revised the kernel amplitudes for the transition to the $J/\psi N$ channel so as to include both t -channel heavy-meson exchange and u -channel heavy-baryon exchange, derived from effective Lagrangians. The coupling constants were estimated from SU(4) symmetry relations, and the resulting uncertainty was absorbed into a single common prefactor a , which controls the overall magnitude of the cross section. The reduced cutoff mass was kept at $\Lambda_0 = 600$ MeV throughout, and the value $a = 0.47$ was the only quantity adjusted to the data.

In this study, we provide a theoretical mechanism for the suppression of the hidden-charm pentaquark signal in J/ψ photoproduction. The rescattering contributions from the $\bar{D}^{(*)}\Sigma_c$ intermediate channels are about one order of magnitude smaller than those from $\bar{D}^{(*)}\Lambda_c$. This difference originates from the coupling constants in the photoproduction kernel: $V_{\bar{D}^{(*)}\Sigma_c, \gamma p}$ is governed by $g_{\bar{D}^{(*)}N\Sigma_c}$, which is about one-fifth of $g_{\bar{D}^{(*)}N\Lambda_c}$. Since the $P_{c\bar{c}}$ resonances couple to the $J/\psi N$ channel predominantly through the $\bar{D}^{(*)}\Sigma_c$ intermediate states, the strong suppression of these channels prevents the pentaquark signal from appearing in J/ψ photoproduction off the proton. The weak signal of $P_{c\bar{c}}(4312)$ that would otherwise emerge through $\bar{D}^*\Lambda_c$ rescattering is further diminished by destructive interference from the $\bar{D}\Sigma_c$ channel.

While the pentaquark signal is suppressed, the GlueX data exhibit cusp structures at the meson-baryon thresholds, which arise mainly from the coupled-channel rescattering effects. We have found that these cusps are enhanced by the inclusion of the u -channel baryon-exchange diagrams, not

considered in our previous work, so that baryon exchange is essential for describing them. With these ingredients, the total cross section agrees well with the GlueX and CLAS12 data over a wide energy range, and the differential cross section as a function of $-t$ is well reproduced in the small $-t$ region. These results confirm that the sizable J/ψ photoproduction cross section is driven primarily by the rescattering effect through open-charm meson-baryon loops.

In conclusion, we have provided a natural explanation for the nonexistence of hidden-charm pentaquark signals in J/ψ photoproduction: the $\bar{D}^{(*)}\Sigma_c$ rescattering channels, which are indispensable for the formation of the $P_{c\bar{c}}$ states, couple only weakly to the γN initial state because $g_{\bar{D}^{(*)}N\Sigma_c}$ is much smaller than $g_{\bar{D}^{(*)}N\Lambda_c}$. This contrasts with Λ_b decays, where the $\bar{D}\Sigma_c$ channel acts as a filter that selects the pentaquark states. The present results therefore suggest that the null result from GlueX and CLAS12 need not be in conflict with the pentaquark signals observed by the LHCb Collaboration and that the absence of a pentaquark signal in photoproduction is not by itself evidence against the existence of hidden-charm pentaquark baryons.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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