

Discovering an unquenched dynamics mechanism for charmonium scattering

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In this paper, we propose a new mechanism for charmonium scattering that utilizes the internal structure of charmonium. By capturing light flavor quarks and antiquarks from the vacuum, charm, and anticharm quarks form virtual charmed mesons, which mediate an effective one-boson exchange process. This approach accurately reproduces the $di-J/\psi$ invariant mass spectrum observed by CMS and LHCb, demonstrating its validity. Our mechanism offers a comprehensive framework for understanding charmonium scattering and is applicable to the scattering problems involving all fully heavy hadrons, an area of increasing interest.

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I. INTRODUCTION

Scattering problems, as a fundamental topic in physics, offering critical insights into the interactions and internal structures of matter. Using this approach, Yukawa predicted the pion meson [1], which mediates the nuclear force, establishing a cornerstone of modern nuclear physics. In this century, advances in experimental techniques have led to the discovery of numerous new hadronic states, presenting unique opportunities to explore hadron scattering. Noteworthy achievements include the identification of exotic molecular states such as P_c states [2,3] and $T_{cc}(3875)^+$ [4,5]. These discoveries not only mark the advent of Particle 2.0 but also enhance our understanding of the nonperturbative aspects of the strong force [6–15].

Recently, the CMS Collaboration reported a high-precision measurement of the invariant mass spectrum of $di-J/\psi$ [16], confirming the previous observation of $X(6900)$ by the LHCb Collaboration [17] and revealing additional enhancements in the $di-J/\psi$ invariant mass spectrum. Furthermore, ATLAS has joined this effort, alongside LHCb and CMS [18]. This new finding surprised the community, prompting various explanations, including exotic hadronic states [19–43] and coupled-channel dynamics generation [44–53]. However, a substantial issue—charmonium scattering—has not been fundamentally resolved, despite some tentative efforts [44–53]. The highly effective one-boson (π , σ , ρ , $\dots$) exchange mechanism for describing hadron scattering fails in quantitatively characterizing charmonium scattering, as charmonium is a typical example of a fully heavy hadron. Thus, studying how charmonium undergoes scattering has become a particularly typical issue in current scattering problems and urgently requires a solution.

In this paper, we propose a new mechanism for charmonium scattering. As illustrated in Fig. 1, the charm and anticharm quarks within charmonium can capture light flavor antiquark and quark, respectively, created from the vacuum to form virtual charmed mesons. These virtual mesons provide a source that enables one-boson exchange

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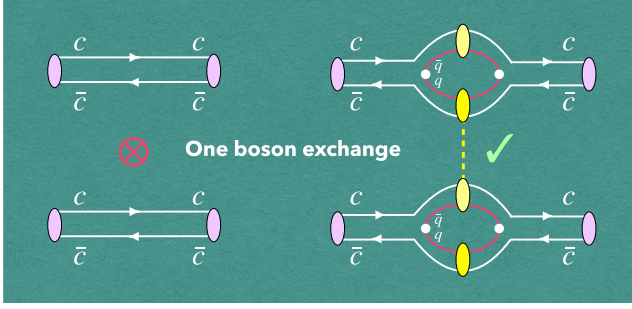


FIG. 1. A new approach enabling one-boson exchange for effective charmonium scattering. The cross symbol on the left indicates that direct light flavor one-boson exchange ($\pi, \rho, \sigma, \dots$) cannot occur.

to be effective for charmonium scattering. A direct application of this new mechanism successfully reproduces the di- J/ψ invariant mass spectrum data from CMS and LHCb, demonstrating the mechanism's validity. We should emphasize that this mechanism is not limited to charmonium scattering; it can also be applied to investigate the scattering problems involving all fully heavy hadrons, a topic that is gaining attention from the community, including lattice research groups [54,55].

II. APPLICATION TO REPRODUCE THE CMS AND LHCb DATA

The di- J/ψ invariant mass spectrum observed by the CMS and LHCb Collaborations exhibit significant features around the dip near 6.8 GeV and the peak near 6.9 GeV, corresponding closely to the $J/\psi\psi\psi(3686)$ and $J/\psi\psi\psi(3770)$ thresholds, respectively. In contrast, the bumps near 6.6 GeV and 7.3 GeV do not align with any thresholds involving particles from the ψ family. Therefore, to capture the main features, we consider couplings among three channels: $J/\psi J/\psi, J/\psi\psi(3686), J/\psi\psi(3770)$. As for the di- η_c and di- χ_{cJ} channels, although there may exist thresholds that close to the observed peaks, due to their much lower production cross sections [56–61], we neglect their effects in this work.

When one-boson exchange is introduced at the quark level, as depicted in Fig. 1, the scattering mechanism between two charmonia can proceed as shown in Fig. 2. Initially, the two charmonia, ψ_1 and ψ_2 , couple into pairs of virtual D mesons. Within each pair, one D meson acts as a spectator, while the other engages in an interaction mediated by the exchange of a light hadron, $\mathcal{P}$. Subsequently, the interacting virtual D mesons combine with the spectators to form the final charmonia, ψ_3 and ψ_4 .

As shown in Fig. 2, we have neglected the contributions from charmed-strange meson loops. This is primarily due to the significantly lower creation rate of $s\bar{s}$ quark pairs compared to $u\bar{u}$ or $d\bar{d}$ quark pairs, with a ratio of approximately $1/\sqrt{3}$. Additionally, the higher threshold

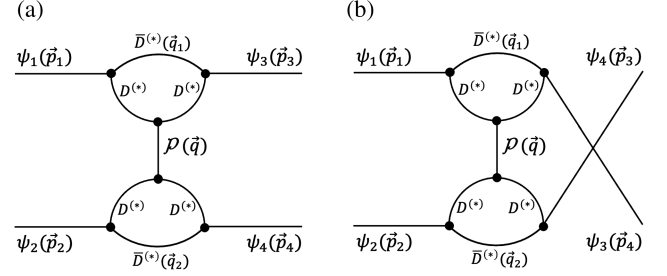


FIG. 2. The scattering mechanisms of the $\psi_1\psi_2 \rightarrow \psi_3\psi_4$ process involve D meson loops, with diagram (a) representing the t -channel contribution and diagram (b) representing the u -channel contribution. Here, ψ_i can be $J/\psi, \psi(3686)$, or $\psi(3770)$, and $\mathcal{P}$ can be η or σ . Furthermore, when $\psi_1 = \psi_2$ and $\psi_3 = \psi_4$, there is no u -channel contribution.

for producing charmed-strange meson pairs leads to further suppression on loop function. Therefore, to focus on the primary physics, we exclude contributions from charmed-strange meson loops.

Furthermore, Fig. 2 illustrates that the exchanged light hadron, $\mathcal{P}$, is limited to σ and η . This restriction arises because the final interaction between two charmonia and a light hadron still must comply with the conservation laws of the strong interaction. Given that the mass of the exchanged particle should be as small as possible, only σ and η are viable candidates, they can contribute the intermediate and long range interactions, whereas, the direct charmed mesons exchanges provide a very short range interactions due to their large masses [1], although it seems to be more efficient interactions to study the charmonia interactions. Our result demonstrates that within our model the sigma-exchange dominates over the eta-exchange, as expected.¹

We perform the entire calculation via the effective Lagrangian approach, with relevant Lagrangians being [62–66]

$$\mathcal{L}_{D^{(*)}D^{(*)}\eta} = -\frac{2g}{\sqrt{6}f_\pi} (i\epsilon_{\alpha\mu\nu\lambda} v^\alpha D^{*\mu} D_a^{*\lambda} \partial^\nu \eta + DD_\lambda^* \partial^\lambda \eta), \quad (1)$$

$$\mathcal{L}_{D^{(*)}D^{(*)}\sigma} = -2g_s DD\sigma + 2g_s D^* \cdot D^* \sigma, \quad (2)$$

$$\mathcal{L}_{D^* \bar{D}^* J/\psi} = -ig_{D^* D^* \psi} [\psi \cdot \bar{D}^* \overleftrightarrow{\partial} \cdot D^* - \psi^\mu \bar{D}^* \cdot \overleftrightarrow{\partial}^\mu D^* + \psi^\mu \bar{D}^* \cdot \overleftrightarrow{\partial} D^{*\mu}], \quad (3)$$

$$\mathcal{L}_{D^* \bar{D} J/\psi} = g_{D^* D \psi} \epsilon_{\beta\mu\alpha\tau} \partial^\beta \psi^\mu (\bar{D} \overleftrightarrow{\partial}^\tau D^{*\alpha} + \bar{D}^{*\alpha} \overleftrightarrow{\partial}^\tau D), \quad (4)$$

¹With the fit parameters listed in Table I, we calculate the ratio of the $|\mathcal{M}_{J/\psi J/\psi}^I(\sqrt{s})|^2$, defined in Eq. (10), that purely caused by σ and η exchanges. The numerical showed that for both 0^{++} and 2^{++} cases, when $\sqrt{s}$ varies from 6.3 to 9.0 GeV, this ratio is always much larger than one.

$$\mathcal{L}_{D\bar{D}J/\psi} = ig_{DD\psi} \bar{D} \overleftrightarrow{\partial} D. \quad (5)$$

Here, v is the velocity of heavy quark. $g = 0.59$, $f_\pi = 132$ MeV, and $g_s = 0.76$ are the fixed coupling constants [67–71]. $g_{D^*D^*\psi}$, $g_{DD\psi}$, and $g_{D^*D\psi}$ can be related to one single parameter f_ψ as $\frac{g_{D^*D^*\psi}}{m_{D^*}} = \frac{g_{DD\psi}}{m_D} = g_{D^*D\psi} = \frac{m_\psi}{m_D f_\psi}$ due to heavy quark effective theory.

The amplitudes for the diagrams shown in Fig. 2 can be generally expressed as

$$\begin{aligned} \mathcal{M}_{(a)} &= 16\epsilon_1^\mu \epsilon_2^\nu \epsilon_3^{*\alpha} \epsilon_4^{*\beta} \frac{1}{q^2 - m_P^2 + im_P \Gamma_P} \mathcal{F}_q^2(q) \\ &\times \int \frac{d^4 q_1}{(2\pi)^4} \frac{\mathcal{V}_{\mu\alpha}}{\prod_i (p_i^2 - m_i^2 + im_i \Gamma_i)} \mathcal{F}_{q_1}(q_1) \\ &\times \int \frac{d^4 q_2}{(2\pi)^4} \frac{\mathcal{V}_{\nu\beta}}{\prod_j (p_j^2 - m_j^2 + im_j \Gamma_j)} \mathcal{F}_{q_2}(q_2), \quad (6) \end{aligned}$$

$$\begin{aligned} \mathcal{M}_{(b)} &= 16\epsilon_1^\mu \epsilon_2^\nu \epsilon_3^{*\alpha} \epsilon_4^{*\beta} \frac{1}{q^2 - m_P^2 + im_P \Gamma_P} \mathcal{F}_q^2(q) \\ &\times \int \frac{d^4 q_1}{(2\pi)^4} \frac{\mathcal{V}_{\mu\alpha}}{\prod_i (p_i^2 - m_i^2 + im_i \Gamma_i)} \mathcal{F}_{q_1}(q_1) \\ &\times \int \frac{d^4 q_2}{(2\pi)^4} \frac{\mathcal{V}_{\nu\beta}}{\prod_j (p_j^2 - m_j^2 + im_j \Gamma_j)} \mathcal{F}_{q_2}(q_2), \quad (7) \end{aligned}$$

where $\mathcal{V}_{\gamma\delta}$ denotes the Lorentz structure of the loop integral, $\mathcal{F}_q(q)$ is the form factor for the exchanged light hadron $\mathcal{P}$ to account for the off-shell effect, and $\mathcal{F}_{q_i}(q_i)$ is a form factor used to avoid ultraviolet divergence in the loop integral. We use the widely adopted monopole form for $\mathcal{F}_q(q)$ as $\mathcal{F}_q(q) = \frac{\Lambda_q^2 - m_q^2}{\Lambda_q^2 - q^2}$, where Λ_q is parametrized as $\Lambda_q = m_q + \alpha_q \Lambda_{\text{QCD}}$, with $\Lambda_{\text{QCD}} = 220$ MeV and α_q is an adjustable parameter. For $\mathcal{F}_{q_i}(q_i)$, we adopt $\mathcal{F}_{q_i}(q_i) = \exp(-\vec{q}_i^2/\Lambda_i^2)$.

For both Eqs. (6) and (7), we multiply them with an overall factor of 16. It is because each D meson loop in Fig. 2 can involve not only $D^{(*)}D^{(*)}\bar{D}^{(*)}$, but also $\bar{D}^{(*)}\bar{D}^{(*)}D^{(*)}$ configurations. Furthermore, for both of these cases, there have charged and neutral situations. Thus, each loop has a total of 4 configurations, resulting in 16 configurations for each diagram in Fig. 2. We want to emphasize here that for $J/\psi J/\psi$ interactions in this model, although it seems the diagram may be suppressed by the masses of the exchanged mesons, like $1/(m_{J/\psi}^2 - 4m_D^2) \times 1/m_P^2$, this suppression can be overcome by the summations on the amplitudes with all the possible intermediate $D^{(*)}$ and $\bar{D}^{(*)}$ combinations in addition with a large but not unreasonable cutoff in form factor. In particular, when the discussed charmonia has a larger mass, the two D -mesons connected with it will be more close to their mass shells, which can reduce the suppression factor significantly.

In this paper, we consider only the S -wave behavior in the di- J/ψ spectrum. Therefore, we first perform a partial wave expansion on the amplitude to get the interaction kernel, which can be calculated as

$$\mathcal{K}_{\lambda,\lambda'}^J(\sqrt{s}, k, k') = 2\pi \int d\cos\theta d_{\lambda,\lambda'}^J(\theta) \mathcal{M}_{\lambda,\lambda'}(\sqrt{s}, \vec{k}, \vec{k}') \quad (8)$$

with $k^{(\prime)} = |\vec{k}^{(\prime)}|$ being the module of the three-momenta of initial and final particles in the center-of-mass system, λ and λ' is the helicity indices of the initial and final states, respectively, and $d_{\lambda,\lambda'}^J(\theta)$ is the Wigner- d matrix with θ being the relative angle between $\vec{k}$ and $\vec{k}'$.

Next, we substitute the partial wave interaction kernel into the coupled-channel Lippmann-Schwinger equation

$$\begin{aligned} T_{\lambda,\lambda'}^J(\sqrt{s}, k, k') &= \mathcal{K}_{\lambda,\lambda'}^J(\sqrt{s}, k, k') + \sum_{\lambda''} \int \frac{k''^2 dk''}{(2\pi)^3} \\ &\times \mathcal{K}_{\lambda,\lambda''}^J(\sqrt{s}, k, k'') G_0(\sqrt{s}, k'') \\ &\times T_{\lambda'',\lambda'}^J(\sqrt{s}, k'', k'), \quad (9) \end{aligned}$$

then, with the extracted on-shell coupled-channel T -matrices $T^J(\sqrt{s})$, the transition amplitude can be written as [47,48]

$$\begin{aligned} \mathcal{M}_{J/\psi J/\psi}^J(\sqrt{s}) \\ = \alpha e^{-\beta s} \left(1 + \sum_i r_i e^{i\phi_i} G_{0i}(\sqrt{s}) T_{i,J/\psi J/\psi}^J(\sqrt{s}) \right), \end{aligned}$$

where G_0 is the free two particle Green's function, G_{0i} means the free Green's function of i th channel, α represents an overall factor, r_i and ϕ_i describe the relative production strength and phase angle for channel i , respectively, and β is used to simulate the background. Although the introduction of r_i and ϕ_i will violate the unitarity, however, unitarity is strictly maintained only within the scattering sector—calculated via the Lippmann-Schwinger equation for the coupled channels $\{J/\psi J/\psi, J/\psi\psi(3686), J/\psi\psi(3770)\}$. When this unitary scattering amplitude is combined with the parametrized, noncalculated production amplitude to form the full observable amplitude, overall unitarity is broken, as we do not perform a fully coupled-channel calculation including the $p\bar{p}$ state.

Finally, the di- J/ψ spectrum can be calculated as [47,48]

$$\sigma(s) = \frac{|\vec{p}_{J/\psi}|}{8\pi s} |\mathcal{M}_{J/\psi J/\psi}^J(\sqrt{s})|^2, \quad (10)$$

with $\vec{p}_{J/\psi}$ being the three momentum of final J/ψ in the center-of-mass frame.

We primarily focus on the CMS data, utilizing the LHCb data as supplementary due to the higher precision of the CMS muon detector. The fit parameters are as follows, excluding the semifixed overall factors α and β which

control the line shape and magnitude of background, for each channel i , there are two parameters: the phase angle ϕ_i and the relative production strength r_i .

Our model includes cutoff dependent form factors, making the scattering amplitudes inherently cutoff dependent. To avoid lengthy fitting processes, we fix all cutoffs (α_q and Λ_i) to one and treat the decay constants f_ψ as free parameters. This approach absorbs the cutoff dependence into the coupling constants, adhering to the principles of renormalization. Consequently, there are three additional parameters: $f_{J/\psi}$, $f_{\psi(3686)}$, and $f_{\psi(3770)}$, resulting in a total of 11 parameters for our fit.

For the di- J/ψ invariant mass spectrum, considering only the S -wave effective potential in the scattering, there are two possible quantum numbers: 0^{++} and 2^{++} . The 1^{++} configuration is excluded because, although $1 \otimes 1 \rightarrow 1$ is possible, it cannot form a state with $C = +$ parity under S -wave conditions. Table I presents the results of the two sets of fit parameters for the 0^{++} and 2^{++} configurations. The corresponding fit results extracted pole positions, based on the parameters in Table I, are shown in Fig. 3 and Table II, respectively.

As shown in Fig. 3, our model successfully reproduces the overall features of the di- J/ψ invariant mass spectrum. This indicates that our proposed scattering mechanism between two charmonium states is reasonable, providing a new approach to understanding hadron scattering, especially for fully heavy hadrons. Furthermore, both the line shape of the black solid lines and the $\chi^2/\text{d.o.f.}$ values in Table I suggest a slight preference for the 0^{++} configuration in the di- J/ψ invariant mass spectrum, but 2^{++} configuration is also possible. Here, the $\chi^2/\text{d.o.f.}$ values are not close to 1, primarily because there are no thresholds near 6.6 and 7.3 GeV in our fit, leading to mismatches in these regions. Despite this, our primary goal of validating the new scattering mechanism is achieved through the correspondence with experimental results.

In addition, Fig. 3 reveals a twist at 6.865 GeV, which is just located at the $J/\psi\psi(3770)$ threshold. Obviously, from the diagram given in Fig. 2, it can be inferred that such phenomenon is attributed to the $\psi(3770)$, whose mass is above the $D\bar{D}$ threshold and thus can decay into a pair of real D mesons. As a result, when the scattering energy reach up to the $J/\psi\psi(3770)$ threshold, rescattering between coupled channels exceed the branch point located at the $J/\psi\psi(3770)$ threshold, which can cause a cusp at this threshold [44]. In addition, since $\psi(3770)$ can be a real particle at this time, when the two charmed mesons linked with it are $D\bar{D}$, the branch point located at $D\bar{D}$ threshold is simultaneously triggered, which may result into further enhancements and generate a tall but very narrow peak just around the $J/\psi\psi(3770)$ threshold, and consequently only slightly modifies the shape of the peak near 6.9 GeV. As for the triangle singularity, based on the algorithm we used in our previous work [72], it can occur when the intermediate

TABLE I. Fit parameters for reproducing the CMS experimental data in the 0^{++} and 2^{++} configurations. Parameters marked with an asterisk (*) are fixed during the fit.

Parameter	0^{++}	2^{++}
α	617.338*	981.439*
β	0.0447*	0.0528*
$r_{J/\psi J/\psi}$	1.101	0.429
$\phi_{J/\psi J/\psi}$	2.979	3.179
$r_{J/\psi\psi(3686)}$	0.724	0.487
$\phi_{J/\psi\psi(3686)}$	2.484	2.745
$r_{J/\psi\psi(3770)}$	1.508	0.268
$\phi_{J/\psi\psi(3770)}$	0.196	0.542
$f_{J/\psi}$	0.378 GeV	0.428 GeV
$f_{\psi(3686)}$	0.237 GeV	0.116 GeV
$f_{\psi(3770)}$	0.760 GeV	0.267 GeV
$\chi^2/\text{d.o.f.}$	2.067	2.268

TABLE II. The pole positions $m - \frac{\Gamma}{2}i$ extracted from our fit parameters given in Table I for the 0^{++} and 2^{++} cases. All values are given in units of MeV. Here, the uncertainties in the pole positions come from the variations of the cutoffs within 50 MeV.

0^{++} case	2^{++} case
$(6204 \pm 24) - (23 \pm 23)i$	$(6198 \pm 7) - (7 \pm 7)i$
$(6789 \pm 17) - (51 \pm 51)i$	$(6783 \pm 9) - (8 \pm 8)i$
$(6871 \pm 11) - (25 \pm 25)i$	$(6871 \pm 10) - (38 \pm 38)i$

states in the triangle loop is $D^*\bar{D}D$, with $D^*\bar{D}$ being linked with charmonia, and D^*D connected to the light exchanged boson. However, at this time, the kinematical limitation needs the charmonia should be an intermediate virtual particle with energy in a range from 3.872 to 3.875 GeV, and the exchanged light boson should have energy less than 0.14 GeV. Apparently, such an effect is highly possible to be suppressed, since either the charmonia or the exchanged light boson is deviated from the center value of mass. Thus, the presence of such a strong threshold effect is a crucial feature of our model. More precise future experimental measurements around 6.9 GeV in the future would greatly aid in validating our model.

Moreover, the wide peak around 6.6 GeV can be reproduced by interference effects without introducing additional channels. The sharp resonancelike structure in $\sigma_{J/\psi J/\psi}$ in both configurations suggests a potential $J/\psi J/\psi$ bound state, as it is not reflected in the overall spectrum due to phase space limitations. This conclusion aligns with Refs. [48–50]. Importantly, if seeing Table II and Fig. 3 together, the dip near 6.8 GeV, particularly evident in the $J/\psi\psi(3686)$ contribution in the 2^{++} fit, indicates the existence of a state similar to $\Upsilon(10753)$, which also shows as a dip in $e^+e^- \rightarrow b\bar{b}$ cross section [73]. This pole around 6.8 GeV is also identified in Refs. [48,50,52,53]. Thus,

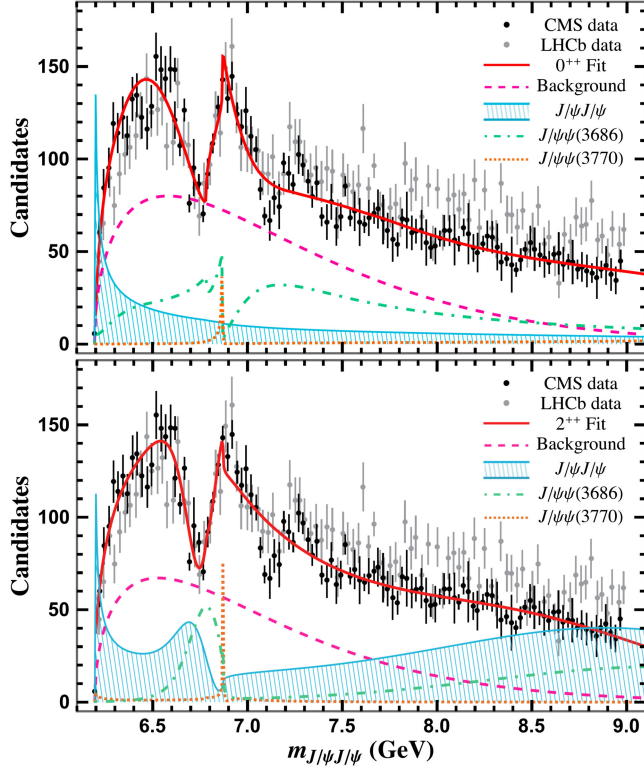


FIG. 3. The fit results on experimental data are presented, with the top panel for 0^{++} configuration and the bottom panel for 2^{++} configuration. In each panel, the dark and light gray error bars represent the CMS data and LHCb data respectively, and the red solid line indicates the fit result. The magenta dashed line, light blue line filled with shadows, light green dot-dashed line, and orange dotted line show the contributions of background, $J/\psi J/\psi$ channel, $J/\psi\psi(3686)$ channel, and $J/\psi\psi(3770)$ channel, respectively.

detailed future experimental analyses of this dip could further test our model.

Overall, it indicates that for either 0^{++} or 2^{++} configurations, there always exists a states below the $J/\psi J/\psi$ threshold. In addition, $X(6600)$ is not a real physical state but an effect of interference, while the dip around 6.8 GeV correspond to a pole mainly generated by the $J/\psi\psi(3686)$ channel. The peak around 6.9 GeV is dominantly generated by a physical state emerging from the $J/\psi\psi(3770)$ channel, but the strong threshold cusp effect there will also make an explicit twist on the shape of the peak. Due to the reason that our mechanism integrates nearly all the conclusions from previous studies into a unified framework, this initial success in understanding the scattering mechanism between two charmonium states motivates further study and broader application of our model. This work represents the first step in that direction.

III. SUMMARY

Recently, the CMS Collaboration reported a high-precision measurement of the di- J/ψ invariant mass spectrum [16], which confirmed the previous observation

of $X(6900)$ by the LHCb Collaboration [17] and revealed additional enhancements. The ATLAS Collaboration has also contributed to this effort, reinforcing these findings [18]. These observations have spurred a variety of theoretical interpretations, including the existence of exotic hadronic states and dynamics generated by coupled channels [19–43]. However, a critical issue remains unresolved: the nature of charmonium scattering. Traditional one-boson exchange mechanisms, highly effective for describing hadron scattering, fall short in quantitatively characterizing charmonium scattering. This shortcoming is particularly pronounced in fully heavy hadron systems.

To solve this puzzle, in this letter, we introduce a new mechanism for charmonium scattering. Our approach leverages the internal structure of charmonium, where charm and anticharm quarks can capture light flavor antiquarks and quarks from the vacuum, forming virtual charmed mesons. These virtual mesons act as mediators, enabling an effective one-boson exchange process for charmonium scattering.

Our proposed mechanism successfully reproduces the di- J/ψ invariant mass spectrum data from CMS and LHCb [16,17], validating its effectiveness. A significant advantage of this approach is its broad applicability. It is not limited to charmonium scattering but can be extended to investigate the scattering of all fully heavy hadrons, an area of increasing interest within the physics community, including lattice QCD research groups [54,55]. With our new mechanism, we can easily predict that there should also exist several states in the di- χ_{cJ} system, which are all very closed to their corresponding thresholds. Furthermore, due to the heavy quark symmetry, our mechanism should also exist in $\Upsilon\Upsilon$, $J/\psi\Upsilon$, $\chi_{bJ}\chi_{bJ'}$, $\chi_{cJ}\chi_{bJ'}$ channels, which can result into the generations of near threshold states. Nevertheless, this approach addresses the limitations of traditional one-boson exchange mechanisms, which are inadequate for fully heavy hadron systems due to the absence of light quark exchanges. It lays the groundwork for a deeper understanding of scattering mechanisms in fully heavy hadron systems, with potential implications for broader applications in particle physics. Thus, we emphasize the importance of further experimental investigations, particularly high-precision measurements around key energy thresholds, to refine and validate our model.

Note added. We noticed that the CMS collaboration has measured the spin-parity of all-charm tetraquarks [74]. The result showed that the quantum numbers of these states prefer $J^{PC} = 2^{++}$, which corroborates the conclusions of this work.

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