$Z_{cs}(3985)^-$: A strange hidden-charm tetraquark resonance or not?

Rui Chen^{1,2,†} and Qi Huang^{3,*}

¹Center of High Energy Physics, Peking University, Beijing 100871, China

²School of Physics and State Key Laboratory of Nuclear Physics and Technology, Peking University, Beijing 100871, China

³School of Physical Sciences, University of Chinese Academy of Sciences, Beijing 100049, China

(Received 24 November 2020; accepted 21 January 2021; published 11 February 2021)

Inspired by the newly $Z_{cs}(3985)^-$ reported by the BESIII Collaboration in the K^+ recoil-mass spectrum of the $e^+e^- \rightarrow (D^{*0}D_s^-/D^0D_s^{*-})K^+$ processes, we perform a dynamical study on the $D^{(*)0}D_s^{*-}$ interactions by adopting a one-boson-exchange model and considering the coupled channel effect. After producing the phase shifts for all the discussed channels, our results exclude the newly $Z_{cs}(3985)^-$ as a $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ resonance with $I(J^P) = 1/2(1^+, 0^-, 1^-, 2^-)$.

DOI: 10.1103/PhysRevD.103.034008

I. INTRODUCTION

Since the BESIII Collaboration reported a discovery of $Z_c(3900)$ in the $e^+e^- \rightarrow \pi^+\pi^-J/\psi$ process in 2013 [1], in the past decade, a series of charged and neutral Z_c states were observed [2–7]. The discoveries of these Z_c states suddenly caught much attention for their novel properties. For example, $Z_c(3900)$ was observed in the $J/\psi\pi^\pm$ invariant mass spectrum of the $e^+e^- \rightarrow \pi^+\pi^-J/\psi$ process at $\sqrt{s} = 4.26$ GeV, and due to its charged property and decay final states, $Z_c(3900)$ should have exotic quark configurations, e.g., it is different with the conventional mesons ($q\bar{q}$) and baryons (qqq). Thus, these Z_c states bring great challenges for the conventional quark model and understanding of strong interactions.

Very recently, the BESIII Collaboration did an analysis on the processes of $e^+e^- \rightarrow (D^{*0}D_s^-/D^0D_s^{*-})K^+$ and observed a structure $Z_{cs}^-(3985)$ in the K^+ recoil-mass spectrum when $\sqrt{s} = 4.681$ GeV [8], whose pole position is

$$M = 3982.5_{-2.6}^{+1.8} \pm 2.1 \text{ MeV},$$

$$\Gamma = 12.8_{-4.4}^{+5.3} \pm 3.0 \text{ MeV}.$$

The observation of $Z_{cs}(3985)$ is widely discussed [9–12]. Similar to our above description on $Z_c(3900)$, it is easy to get that it may have four different valence quark

components [$c\bar{c}s\bar{u}$]. Thus, it is the first candidate of the charged hidden-charm tetraquark state with strangeness. Actually, as mentioned in [8], it is a partner of currently existing $Z_c(3885)$.

As shown in Fig. 1, the $Z_{cs}(3985)$ is a little above the $D^{*0}D_s^-$ and $D^0D_s^{*-}$ thresholds (right panel) compared to the positions of the $Z_c(3885)$ and $Z_c(4025)$ (left panel), but we can easily see its similarity with the $Z_c(3885)$ and $Z_c(4025)$. Since the $Z_c(3885)$ and $Z_c(4020)$ have often been assigned as hidden-charm hadronic molecules or resonances [13–16], one may propose whether the newly $Z_{cs}(3985)$ can be a strange hidden-charm tetraquark resonance. In Ref. [17], the mass of the $cq\bar{c}\bar{s}$ was predicted as 3982 MeV within the relativistic diquark-antidiquark picture. In this work, we study the $D^{(*)}\bar{D}_s^{(*)}$ interactions by using the one-boson-exchange (OBE) model to examine this possibility.

In general, the resonances can be divided into two types, e.g., shape-type and Feshbach-type resonances. As we know, the potential barriers can be beneficial to generate these two types of resonances. In particular, we need to mention that the coupled channel effect plays a very important role in the generation of Feshbach-type resonances, where the mass gaps between the discussed channels contribute to the barriers.

In the following, we will consider the coupled-channel effect and $S - D$ wave mixing effect, which can provide mass gap barriers and centrifugal potential barriers, respectively. Since the spin parity of Z_{cs} is not fixed yet, the different spin-parities $I(J^P) = 1/2(1^+, 0^-, 1^-, 2^-)$ for the coupled $D^{(*)}\bar{D}_s^{(*)}$ systems will be discussed.

This paper is organized as follows. After this introduction, we present the $D^{(*)}\bar{D}_s^{(*)}$ interactions by using the OBE model in Sec. II. The corresponding numerical results are given in Sec. III. The paper ends with a summary in Sec. IV.

*Corresponding author.
huangqi@ucas.ac.cn
†chen_rui@pku.edu.cn

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

FIG. 1. A mass comparison between Z_c (Z_{cs}) states and thresholds of a charmed meson and an anticharmed (charm-strange) meson.

II. INTERACTIONS

For the coupled $D^{*0}\bar{D}_s^-/D^0\bar{D}_s^{*-}/D^{*0}\bar{D}_s^{*-}$ system, in the OBE model, there only exists contributions of σ and η exchanges. Here, we need to point out that the K/K^* -exchanges are forbidden in the cross processes $D^*\bar{D}_s \rightarrow D_s^{(*)}\bar{D}^{(*)}$. Since the descriptions of interactions between charmed mesons and Goldstone bosons are needed, we adopt the effective Lagrangians approach. The effective Lagrangians relevant to the charmed hadrons and light mesons are constructed as

$$\mathcal{L} = g_\sigma \langle H_a^{(Q)} \sigma \bar{H}_a^{(Q)} \rangle + ig \langle H_b^{(Q)} \gamma_\mu A_{ba}^\mu \gamma_5 \bar{H}_a^{(Q)} \rangle + g_\sigma \langle \bar{H}_a^{(\bar{Q})} \sigma H_a^{(\bar{Q})} \rangle + ig \langle \bar{H}_a^{(\bar{Q})} \gamma_\mu A_{ab}^\mu \gamma_5 H_b^{(\bar{Q})} \rangle, \quad (1)$$

which are based on the heavy-quark limit and chiral symmetry [18–23]. Here, $A_\mu = (\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger)/2 = i\partial_\mu \mathbb{P}/f_\pi + \dots$ is the axial current with the pion decay constant $f_\pi = 132$ MeV and $\xi = \exp(i\mathbb{P}/f_\pi)$. In addition, $g = 0.59$ is fixed by fitting the decay width of D^{*+} [24], and $g_s = 0.76$ is given in Ref. [22]. In Eq. (1), we define two multiplet fields $H^{(Q)}$ and $H^{(\bar{Q})}$, which are expressed as a linear combination of a pseudoscalar charmed (anti-charmed) meson and a vector charmed (anticharmed) meson, respectively,

$$H_a^{(Q)} = \frac{1 + \not{p}}{2} [\mathcal{P}_a^{*(Q)\mu} \gamma_\mu - \mathcal{P}_a^{(Q)} \gamma_5],$$

$$H_a^{(\bar{Q})} = [\tilde{\mathcal{P}}_a^{*\mu} \gamma_\mu - \tilde{\mathcal{P}}_a \gamma_5] \frac{1 - \not{p}}{2}. \quad (2)$$

Here, $\mathbf{v} = (1, \mathbf{0})$, $\mathcal{P}^{(*)T} = (D^{(*)+}, D^{(*)0}, D_s^{(*)+})$, and $\tilde{\mathcal{P}}^{(*)T} = (D^{(*)-}, \bar{D}^{(*)0}, D_s^{(*)-})$. The pseudoscalar matrix $\mathbb{P}$ reads as $\mathbb{P} = \text{diag}(\frac{\eta}{\sqrt{6}}, \frac{\eta}{\sqrt{6}}, -\frac{2\eta}{\sqrt{6}})$.

After expanding Eqs. (1) and (2), we obtain the concrete effective Lagrangians as follows:

$$\begin{aligned} \mathcal{L}_\sigma &= -2g_\sigma \mathcal{P}_a \mathcal{P}_a^\dagger \sigma - 2g_\sigma \tilde{\mathcal{P}}_a^\dagger \tilde{\mathcal{P}}_a \sigma \\ &\quad + 2g_\sigma \mathcal{P}_a^* \cdot \mathcal{P}_a^{*\dagger} \sigma + 2g_\sigma \tilde{\mathcal{P}}_a^{*\dagger} \cdot \tilde{\mathcal{P}}_a^* \sigma, \\ \mathcal{L}_\mathbb{P} &= -\frac{2g}{f_\pi} (\mathcal{P}_b^{*\mu} \mathcal{P}_a^\dagger + \mathcal{P}_b \mathcal{P}_a^{*\mu\dagger}) \partial_\mu \mathbb{P}_{ba} \\ &\quad + \frac{2g}{f_\pi} (\tilde{\mathcal{P}}_a^{*\mu\dagger} \tilde{\mathcal{P}}_b + \tilde{\mathcal{P}}_a^\dagger \tilde{\mathcal{P}}_b^{*\mu}) \partial_\mu \mathbb{P}_{ab} \\ &\quad - i \frac{2g}{f_\pi} v^\beta \epsilon_{\beta\mu\alpha\nu} \mathcal{P}_b^{*\mu} \mathcal{P}_a^{*\nu\dagger} \partial^\alpha \mathbb{P}_{ba} \\ &\quad + i \frac{2g}{f_\pi} v^\beta \epsilon_{\beta\mu\alpha\nu} \tilde{\mathcal{P}}_a^{*\mu\dagger} \tilde{\mathcal{P}}_b^{*\nu} \partial^\alpha \mathbb{P}_{ab}. \end{aligned} \quad (3)$$

With these effective Lagrangians, we can easily deduce the corresponding OBE effective potentials, which are related to the scattering amplitudes by using a Breit approximation in a momentum space, i.e.,

$$\mathcal{V}_E^{h_1 h_2 \rightarrow h_3 h_4}(\mathbf{q}) = -\frac{\mathcal{M}(h_1 h_2 \rightarrow h_3 h_4)}{\sqrt{\prod_i 2M_i \prod_f 2M_f}}. \quad (4)$$

Here, M_i and M_f are the masses of the initial states (h_1, h_2) and final states (h_3, h_4), respectively. Note that $\mathcal{M}(h_1 h_2 \rightarrow h_3 h_4)$ is the scattering amplitude for the $h_1 h_2 \rightarrow h_3 h_4$ process. After performing a Fourier transformation, we can obtain the effective potential in the coordinate space $\mathcal{V}(\mathbf{r})$, i.e.,

$$\mathcal{V}_E^{h_1 h_2 \rightarrow h_3 h_4}(\mathbf{r}) = \int \frac{d^3 \mathbf{q}}{(2\pi)^3} e^{i\mathbf{q} \cdot \mathbf{r}} \mathcal{V}_E^{h_1 h_2 \rightarrow h_3 h_4}(\mathbf{q}) \mathcal{F}^2(q^2, m_E^2),$$

where we introduce a monopole form factor $\mathcal{F}(q^2, m_E^2) = (\Lambda^2 - m_E^2)/(\Lambda^2 - q^2)$ at every interactive vertex, which compensates the off-shell effect of the exchanged bosons. Here, Λ , m_E , and q denote the cutoff, mass, and four-momenta of the exchanged meson, respectively.

Finally, the OBE effective potentials for the coupled $D^{*0}\bar{D}_s^-/D^0\bar{D}_s^{*-}/D^{*0}\bar{D}_s^{*-}$ systems are

$$\mathcal{V}_{\text{OBE}} = \begin{pmatrix} \langle D^{*0}\bar{D}_s^- | \mathcal{V}_\sigma | D^{*0}\bar{D}_s^- \rangle & \langle D^{*0}\bar{D}_s^- | \mathcal{V}_\eta | D^0\bar{D}_s^{*-} \rangle & \langle D^{*0}\bar{D}_s^- | \mathcal{V}_\eta | D^{*0}\bar{D}_s^{*-} \rangle \\ \langle D^0\bar{D}_s^{*-} | \mathcal{V}_\eta | D^{*0}\bar{D}_s^- \rangle & \langle D^0\bar{D}_s^{*-} | \mathcal{V}_\sigma | D^0\bar{D}_s^{*-} \rangle & \langle D^0\bar{D}_s^{*-} | \mathcal{V}_\eta | D^{*0}\bar{D}_s^{*-} \rangle \\ \langle D^{*0}\bar{D}_s^{*-} | \mathcal{V}_\eta | D^{*0}\bar{D}_s^- \rangle & \langle D^0\bar{D}_s^{*-} | \mathcal{V}_\eta | D^0\bar{D}_s^{*-} \rangle & \langle D^{*0}\bar{D}_s^{*-} | \mathcal{V}_\sigma + \mathcal{V}_\eta | D^{*0}\bar{D}_s^{*-} \rangle \end{pmatrix}, \quad (5)$$

where the subpotentials are

$$\mathcal{V}_\sigma^{D^{*0}D_s^- \rightarrow D^{*0}D_s^-} = -g_s^2 \epsilon_1 \cdot \epsilon_3^\dagger Y(\Lambda, m_\sigma, r), \quad (6)$$

$$\mathcal{V}_\sigma^{D^0D_s^{*-} \rightarrow D^0D_s^{*-}} = -g_s^2 \epsilon_2 \cdot \epsilon_4^\dagger Y(\Lambda, m_\sigma, r), \quad (7)$$

$$\mathcal{V}_\sigma^{D^{*0}D_s^{*-} \rightarrow D^{*0}D_s^{*-}} = -g_s^2 \epsilon_1 \cdot \epsilon_3^\dagger \epsilon_2 \cdot \epsilon_4^\dagger Y(\Lambda, m_\sigma, r), \quad (8)$$

$$\begin{aligned} \mathcal{V}_\eta^{D^{*0}D_s^- \rightarrow D^0D_s^{*-}} &= \frac{1}{9} \frac{g^2}{f_\pi^2} [\epsilon_1 \cdot \epsilon_4^\dagger Z(\Lambda_1, m_{\eta 1}, r) \\ &\quad + S(\hat{r}, \epsilon_1, \epsilon_4^\dagger) T(\Lambda_1, m_{\eta 1}, r)], \end{aligned} \quad (9)$$

$$\begin{aligned} \mathcal{V}_\eta^{D^{*0}D_s^- \rightarrow D^{*0}D_s^{*-}} &= \frac{1}{9} \frac{g^2}{f_\pi^2} [i(\epsilon_1 \times \epsilon_3^\dagger) \cdot \epsilon_4^\dagger Z(\Lambda_2, m_{\eta 2}, r) \\ &\quad + S(\hat{r}, i(\epsilon_1 \times \epsilon_3^\dagger), \epsilon_4^\dagger) T(\Lambda_2, m_{\eta 2}, r)], \end{aligned} \quad (10)$$

$$\begin{aligned} \mathcal{V}_\eta^{D^0D_s^{*-} \rightarrow D^{*0}D_s^{*-}} &= \frac{1}{9} \frac{g^2}{f_\pi^2} [i(\epsilon_2 \times \epsilon_3^\dagger) \cdot \epsilon_4^\dagger Z(\Lambda_3, m_{\eta 3}, r) \\ &\quad + S(\hat{r}, i(\epsilon_2 \times \epsilon_3^\dagger), \epsilon_4^\dagger) T(\Lambda_3, m_{\eta 3}, r)], \end{aligned} \quad (11)$$

$$\begin{aligned} \mathcal{V}_\eta^{D^{*0}D_s^{*-} \rightarrow D^{*0}D_s^{*-}} &= \frac{1}{9} \frac{g^2}{f_\pi^2} [(\epsilon_1 \times \epsilon_3^\dagger) \cdot (\epsilon_2 \times \epsilon_4^\dagger) Z(\Lambda, m_\eta, r) \\ &\quad + S(\hat{r}, \epsilon_1 \times \epsilon_3^\dagger, \epsilon_2 \times \epsilon_4^\dagger) T(\Lambda, m_\eta, r)], \end{aligned} \quad (12)$$

with

$$\begin{aligned} Y(\Lambda, m, r) &= \frac{1}{4\pi r} (e^{-mr} - e^{-\Lambda r}) - \frac{\Lambda^2 - m^2}{8\pi\Lambda} e^{-\Lambda r}, \\ T(\Lambda, m, r) &= r \frac{\partial}{\partial r} \frac{1}{r} \frac{\partial}{\partial r} Y(\Lambda, m, r), \\ Z(\Lambda, m, r) &= \nabla^2 Y(\Lambda, m, r) = \frac{1}{r^2} \frac{\partial}{\partial r} r^2 \frac{\partial}{\partial r} Y(\Lambda, m, r). \end{aligned}$$

In Eqs. (6)–(12), for conciseness, we introduce many variables, and their definitions are

$$\begin{aligned} \Lambda_1^2 &= \Lambda^2 - q_1^2, & m_{\eta 1}^2 &= m_\eta^2 - q_1^2, \\ q_1 &= \frac{(m_{D^*}^2 + m_{D_s^*}^2) - (m_{D_s}^2 + m_D^2)}{2(m_D + m_{D_s^*})}, \\ \Lambda_2^2 &= \Lambda^2 - q_2^2, & m_{\eta 2}^2 &= m_\eta^2 - q_2^2, \\ \Lambda_3^2 &= \Lambda^2 - q_3^2, & m_{\eta 3}^2 &= m_\eta^2 - q_3^2, \\ q_2 &= \frac{m_{D_s^*}^2 - m_{D_s}^2}{2(m_{D^*} + m_{D_s^*})}, & q_3 &= \frac{m_D^2 - m_{D^*}^2}{2(m_{D^*} + m_{D_s^*})}. \end{aligned}$$

Here, we also define several spin-spin interaction and tensor force operators; in the following numerical calculations, they should be sandwiched by the spin-orbit wave functions for the coupled $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ systems, i.e., $\langle {}^{2S+1}L_J' | \hat{O} | {}^{2S+1}L_J \rangle$. For the $J^P = 1^+$ case, the

spin-orbit wave functions $|{}^{2S+1}L_J\rangle$ are $D^{*0}D_s^-|{}^3S_1, {}^3D_1\rangle$, $D^0D_s^{*-}|{}^3S_1, {}^3D_1\rangle$, and $D^{*0}D_s^{*-}|{}^3S_1, {}^3D_1\rangle$. Therefore, the matrix elements for all the operators are

$$\left\{ \begin{array}{c} \epsilon_1 \cdot \epsilon_3^\dagger \\ \epsilon_1 \cdot \epsilon_4^\dagger \\ \epsilon_1 \cdot \epsilon_3^\dagger \epsilon_2 \cdot \epsilon_4^\dagger \\ \frac{i}{\sqrt{2}} (\epsilon_2 \times \epsilon_3^\dagger) \cdot \epsilon_4^\dagger \\ (\epsilon_1 \times \epsilon_3^\dagger) \cdot (\epsilon_2 \times \epsilon_4^\dagger) \end{array} \right\} \rightarrow (\mathcal{I}), \quad (13)$$

$$\left\{ \begin{array}{c} S(\hat{r}, \epsilon_1, \epsilon_4^\dagger) \\ -\sqrt{2} S(\hat{r}, i\epsilon_2 \times \epsilon_3^\dagger, \epsilon_4^\dagger) \\ S(\hat{r}, \epsilon_1 \times \epsilon_3^\dagger, \epsilon_2 \times \epsilon_4^\dagger) \end{array} \right\} \rightarrow \begin{pmatrix} 0 & -\sqrt{2} \\ -\sqrt{2} & 1 \end{pmatrix}, \quad (14)$$

where $\mathcal{I}$ stands for unit matrix.

III. NUMERICAL RESULTS

In the above OBE effective potentials, there is a cutoff parameter needed to be fixed. According to the quantitative description of the deuteron properties and the NN scattering data (see review in Ref. [25]), the value of cutoff Λ is taken around 1 to 2 GeV. Here, we will adopt this empirical value to give our conclusion.

In Fig. 2, we present the OBE effective potentials for the S wave $D^{*0}D_s^-$ system with several typical cutoff values. Here, we find that the interaction from the σ exchange is weak attractive, and it becomes stronger with the increasing of cutoff Λ . It is obvious that this interaction is not strong enough to form a bound state in the current cutoff choices.

As mentioned above, a repulsive barrier plays an important role in generating a resonance. However, as we can see from Fig. 2, there does not exist any barrier in the attractive interaction from the σ exchange, thus, the S wave $D^{*0}D_s^-$ cannot bind as a scattering state. So, we further introduce the $S-D$ wave mixing since it may be

FIG. 2. OBE effective potential for the S wave $D^{*0}D_s^-$ system with $\Lambda = 1.0, 1.5, 2.0$ GeV and its centrifugal force $l(l+1)/2mr^2$ in D wave $l = 2$.

helpful to generate a $D^{*0}D_s^-$ scattering state with $I(J^P) = 1/2(1^+)$ because the centrifugal force $l(l+1)/2mr^2$ from the D wave $l=2$ can provide a repulsive potential barrier as shown in Fig. 2. However, compared to the weak attractive σ exchange interaction, the repulsive centrifugal force is much stronger. Therefore, we can give an ambitious estimation that there cannot exist a $D^{*0}D_s^-$ scattering state with $I(J^P) = 1/2(1^+)$ when we only consider the single $D^{*0}D_s^-$ systems.

In our calculation, we take the same range of the cutoff value $1.0 \leq \Lambda \leq 2.0$ GeV to produce the phase shifts for the coupled $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ systems. As is well known, when the phase shift satisfies $\delta = (n + 1/2)\pi$ with $n = 0, 1, 2, \dots$, the cross section $\sigma(E)$ takes the maximum $\sigma(E_0)$. It corresponds to a resonance, and its mass and width are

$$M = E_0, \quad \Gamma = 2/\left(\frac{d\delta(E)}{dE}\right)_{E_0}, \quad (15)$$

respectively.

Finally, we cannot find any possible resonances for the S wave single $D^{*0}D_s^-$ and $D^0D_s^{*-}$ when cutoff Λ is taken from 1.0 to 2.0 GeV. If the $S-D$ wave mixing is considered, unfortunately, then we cannot find any resonant evidences for the single $D^{*0}D_s^-$ and $D^0D_s^{*-}$ systems either. Therefore, there does not exist any possible $D^{*0}D_s^-$ and $D^0D_s^{*-}$ resonance candidates from the single channel analysis. Here, we can conclude that the newly Z_{cs} is excluded as a $D^{*0}D_s^-$ or $D^0D_s^{*-}$ shape-type resonance.

In the following, we further consider the coupled channel effect to check whether the newly $Z_{cs}(3985)$ can be a possible Feshbach-type resonance from the $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ interactions. In Fig. 3, we present the OBE effective potentials for the $D^{*0}D_s^- \rightarrow D^0D_s^{*-}$, $D^{*0}D_s^-(D^0D_s^{*-}) \rightarrow D^{*0}D_s^{*-}$ and $D^{*0}D_s^{*-} \rightarrow D^{*0}D_s^{*-}$ processes with $\Lambda = 1.5$ GeV. Here, the OBE effective potentials from $D^{*0}D_s^- \rightarrow D^{*0}D_s^{*-}$ scattering process have the

same numerical values with those in the $D^0D_s^{*-} \rightarrow D^{*0}D_s^{*-}$ process. Compared to the σ exchange, the η exchange interactions are much stronger, which may play an important role in the discussed coupled systems.

In Fig. 4, we take three groups of typical cutoff values $\Lambda = 1.0, 1.5, 2.0$ GeV to plot the energy E dependence of the phase shift δ for the coupled $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ systems with $I(J^P) = 1/2(1^+)$. As shown in Fig. 4, with the increasing of the cutoff Λ , the OBE interactions become much stronger, which leads to the absolute values of the corresponding phase shifts for all channels to become much larger. Also, small peaks appear at the $D^{*0}D_s^-$ threshold in the phase shifts for all the $D^{*0}D_s^-$ and $D^0D_s^{*-}$ channels, which are interpreted as cusps. Most importantly, in the cutoff range of $1.0 \leq \Lambda \leq 2.0$ GeV, phase shifts for all discussed channels, except the $D^{*0}D_s^-|{}^3D_1\rangle$, rise up, but all their values are less than $\pi/2$, which indicates that no resonance exists. Thus, we can further conclude that the newly $Z_{cs}(3985)$ cannot be interpreted as a $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ Feshbach-type resonance.

So, according to our study, we can exclude the newly $Z_{cs}(3985)$ as the $I(J^P) = 1/2(1^+)$ $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ shape or Feshbach resonance in the OBE model.

Apart from $S-D$ wave mixing, we also test if the newly Z_{cs} can be a P wave $D^{(*)}\bar{D}_s^{(*)}$ resonance. In this situation, the spin-orbit wave functions are

$$J^P = 0^- : D^{*0}D_s^-|{}^3P_0\rangle, D^0D_s^{*-}|{}^3P_0\rangle, D^{*0}D_s^{*-}|{}^3P_0\rangle,$$

$$J^P = 1^- : D^{*0}D_s^-|{}^3P_1\rangle, D^0D_s^{*-}|{}^3P_1\rangle, D^{*0}D_s^{*-}|{}^1P_1, {}^3P_1, {}^5P_1\rangle,$$

$$J^P = 2^- : D^{*0}D_s^-|{}^3P_2\rangle, D^0D_s^{*-}|{}^3P_2\rangle, D^{*0}D_s^{*-}|{}^3P_2, {}^5P_2\rangle,$$

and the corresponding matrix elements for the spin-spin interaction and tensor force operators are summarized in Eqs. (16) and (17). For the P wave case, when we vary the cutoff value from 1.0 to 2.0 GeV, we still do not find any resonances. Therefore, we further find that the newly Z_{cs} cannot be a $D^{(*)}\bar{D}_s^{(*)}$ resonance with 0^- , 1^- , or 2^- :

FIG. 3. OBE effective potentials for the $D^{*0}D_s^- \rightarrow D^0D_s^{*-}$, $D^{*0}D_s^-(D^0D_s^{*-}) \rightarrow D^{*0}D_s^{*-}$, and $D^{*0}D_s^{*-} \rightarrow D^{*0}D_s^{*-}$ processes with $\Lambda = 1.5$ GeV.

FIG. 4. Phase shifts for the $D^{*0}D_s^- - D^0D_s^{*-} - D^{*0}D_s^{*-}$ coupled systems with $\Lambda = 1.0, 1.5$, and 2.0 GeV.

$$\left\{ \begin{array}{l} \epsilon_1 \cdot \epsilon_3^\dagger \\ \epsilon_2 \cdot \epsilon_4^\dagger \\ \epsilon_1 \cdot \epsilon_3^\dagger \epsilon_2 \cdot \epsilon_4^\dagger \end{array} \right\} \rightarrow (\mathcal{I})_{J^P=0^-, 1^-, 2^-}, \quad (16)$$

$$\begin{aligned} \frac{i}{\sqrt{2}}(\epsilon_2 \times \epsilon_3^\dagger) \cdot \epsilon_4^\dagger &\rightarrow \{(1)_{0^-}, (0, 1, 0)_{1^-}, (1, 0)_{2^-}\}, \\ (\epsilon_1 \times \epsilon_3^\dagger) \cdot (\epsilon_2 \times \epsilon_4^\dagger) &\rightarrow \{(1)_{0^-}, (\text{diag}(2, 1, -1))_{1^-}, (\text{diag}(1, 0))_{2^-}\}, \\ S(\hat{r}, \epsilon_2, \epsilon_4^\dagger) &\rightarrow \left\{ (2)_{0^-}, (-1)_{1^-}, \left(\frac{1}{5}\right)_{2^-} \right\}, \\ S(\hat{r}, i\epsilon_2 \times \epsilon_3^\dagger, \epsilon_4^\dagger) &\rightarrow \left\{ \left(\frac{-1}{\sqrt{2}}\right)_{0^-}, \left(0, \frac{1}{\sqrt{2}}, 0\right)_{1^-}, \left(\frac{-1}{\sqrt{10}}, 0\right)_{2^-} \right\}, \\ S(\hat{r}, \epsilon_1 \times \epsilon_3^\dagger, \epsilon_2 \times \epsilon_4^\dagger) &\rightarrow \left\{ (2)_{0^-}, \left(\begin{array}{ccc} 0 & 0 & -\frac{2}{\sqrt{5}} \\ 0 & -1 & 0 \\ -\frac{2}{\sqrt{5}} & 0 & \frac{7}{5} \end{array} \right)_{1^-}, \left(\frac{1}{5} \quad 0 \\ 0 \quad -\frac{7}{5} \end{array} \right)_{2^-} \right\}. \end{aligned} \quad (17)$$

IV. SUMMARY

In recent decades, people have paid more and more attention to the study of multi-quark states. In particular, the observations of the charged charmoniumlike structures such as $Z_c(3900)^\pm$, $Z_c \pm(4020)$, and $Z_c^+(4430)$ provide important evidence of the existence of the tetraquark state in the charm sector. Very recently, the BESIII Collaboration reported a new charged charmoniumlike

state $Z_{cs}^-(3985)$ in the K^+ recoil-mass spectrum of the $e^+e^- \rightarrow (D^{*0}D_s^-/D^0D_s^{*-})K^+$ processes, which could be the first candidate of the charged hidden-charm tetraquark state with strangeness.

Stimulated by the observation of the strange charmoniumlike state $Z_{cs}(3985)$ and its behavior of being closed to threshold, in this work, we perform a systematic investigation on the $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ interactions

by using the OBE model. In this situation, only the intermediate range interaction from the σ and η exchanges contributes to the coupled $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ system, and the η exchange effective potential is dominant. These are remarkably different with the $D^{(*)}\bar{D}^*$ interactions since, where the long range force from the π exchange and the short range force from the vector meson ρ/ω exchanges are allowed, in particular, the π exchanges are the most important.

After gradually introducing the $S - D$ wave mixing and the coupled channel effect to study the phase shifts of the coupled $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ systems, finally, our results exclude the newly $Z_{cs}(3985)$ as a shape-type or

Feshbach-type strange hidden-charm tetraquark resonance with $I(J^P) = 1/2(1^+)$. In addition, we further study the P wave $D^{*0}D_s^-/D^0D_s^{*-}/D^{*0}D_s^{*-}$ interactions and find that the newly Z_{cs} cannot be interpreted as a $D^{(*)0}D_s^{*-}$ resonance with $I(J^P) = 1/2(0^-, 1^-, 2^-)$.

ACKNOWLEDGMENTS

R. C. is very grateful to Yasuhiro Yamaguchi for his guidance in the numerical calculation and Xiang Liu for his helpful discussions. R. C. is supported by the National Postdoctoral Program for Innovative Talent.

-
- [1] M. Ablikim *et al.* (BESIII Collaboration), Observation of a Charged Charmoniumlike Structure in $e^+e^- \rightarrow J/\psi\pi^+\pi^-$ at $\sqrt{s} = 4.26$ GeV, *Phys. Rev. Lett.* **110**, 252001 (2013).
 - [2] Z. Q. Liu *et al.* (Belle Collaboration), Study of $e^+e^- \rightarrow \pi^+\pi^-J/\psi$ and Observation of a Charged Charmoniumlike State at Belle, *Phys. Rev. Lett.* **110**, 252002 (2013).
 - [3] T. Xiao, S. Dobbs, A. Tomaradze, and K. K. Seth, Observation of the charged hadron $Z_c^\pm(3900)$ and evidence for the neutral $Z_c^0(3900)$ in $e^+e^- \rightarrow \pi\pi J/\psi$ at $\sqrt{s} = 4170$ MeV, *Phys. Lett. B* **727**, 366 (2013).
 - [4] M. Ablikim *et al.* (BESIII Collaboration), Observation of a Charged Charmoniumlike Structure $Z_c(4020)$ and Search for the $Z_c(3900)$ in $e^+e^- \rightarrow \pi^+\pi^-h_c$, *Phys. Rev. Lett.* **111**, 242001 (2013).
 - [5] M. Ablikim *et al.* (BESIII Collaboration), Observation of a Charged ($D\bar{D}^*$) $^\pm$ Mass Peak in $e^+e^- \rightarrow \pi D\bar{D}^*$ at $\sqrt{s} = 4.26$ GeV, *Phys. Rev. Lett.* **112**, 022001 (2014).
 - [6] M. Ablikim *et al.* (BESIII Collaboration), Observation of a Charged Charmoniumlike Structure in $e^+e^- \rightarrow (D^*\bar{D}^*)^\pm\pi^\mp$ at $\sqrt{s} = 4.26$ GeV, *Phys. Rev. Lett.* **112**, 132001 (2014).
 - [7] M. Ablikim *et al.* (BESIII Collaboration), Measurement of $e^+e^- \rightarrow \pi^+\pi^-\psi(3686)$ from 4.008 to 4.600 GeV and observation of a charged structure in the $\pi^\pm\psi(3686)$ mass spectrum, *Phys. Rev. D* **96**, 032004 (2017).
 - [8] M. Ablikim *et al.* (BESIII Collaboration), Observation of a near-threshold structure in the K^+ recoil-mass spectra in $e^+e^- \rightarrow K^+(D_s^-D^{*0} + D_s^{*-}D^0)$, [arXiv:2011.07855](https://arxiv.org/abs/2011.07855).
 - [9] L. Meng, B. Wang, and S. L. Zhu, $Z_{cs}(3985)$ as the U -spin partner of $Z_c(3900)$ and implication of other states in the $SU(3)_F$ symmetry and heavy quark symmetry, *Phys. Rev. D* **102**, 111502 (2020).
 - [10] J. Z. Wang, Q. S. Zhou, X. Liu, and T. Matsuki, Toward charged $Z_{cs}(3985)$ structure under a reflection mechanism, *Eur. Phys. J. C* **81**, 51 (2021).
 - [11] B. D. Wan and C. F. Qiao, About the exotic structure of Z_{cs} , [arXiv:2011.08747](https://arxiv.org/abs/2011.08747).
 - [12] Z. Yang, X. Cao, F. K. Guo, J. Nieves, and M. P. Valderrama, Strange molecular partners of the $Z_c(3900)$ and $Z_c(4020)$, [arXiv:2011.08725](https://arxiv.org/abs/2011.08725).
 - [13] H. X. Chen, W. Chen, X. Liu, and S. L. Zhu, The hidden-charm pentaquark and tetraquark states, *Phys. Rep.* **639**, 1 (2016).
 - [14] X. Liu, An overview of XYZ new particles, *Chin. Sci. Bull.* **59**, 3815 (2014).
 - [15] F. K. Guo, C. Hanhart, U. G. Meissner, Q. Wang, Q. Zhao, and B. S. Zou, Hadronic molecules, *Rev. Mod. Phys.* **90**, 015004 (2018).
 - [16] N. Brambilla, S. Eidelman, C. Hanhart, A. Nefediev, C. P. Shen, C. E. Thomas, A. Vairo, and C. Z. Yuan, The XYZ states: Experimental and theoretical status and perspectives, *Phys. Rep.* **873**, 1 (2020).
 - [17] D. Ebert, R. N. Faustov, and V. O. Galkin, Excited heavy tetraquarks with hidden charm, *Eur. Phys. J. C* **58**, 399 (2008).
 - [18] T. M. Yan, H. Y. Cheng, C. Y. Cheung, G. L. Lin, Y. C. Lin, and H. L. Yu, Heavy quark symmetry and chiral dynamics, *Phys. Rev. D* **46**, 1148 (1992); Erratum, *Phys. Rev. D* **55**, 5851 (1997).
 - [19] M. B. Wise, Chiral perturbation theory for hadrons containing a heavy quark, *Phys. Rev. D* **45**, R2188 (1992).
 - [20] G. Burdman and J. F. Donoghue, Union of chiral and heavy quark symmetries, *Phys. Lett. B* **280**, 287 (1992).
 - [21] R. Casalbuoni, A. Deandrea, N. Di Bartolomeo, R. Gatto, F. Feruglio, and G. Nardulli, Phenomenology of heavy meson chiral Lagrangians, *Phys. Rep.* **281**, 145 (1997).
 - [22] A. F. Falk and M. E. Luke, Strong decays of excited heavy mesons in chiral perturbation theory, *Phys. Lett. B* **292**, 119 (1992).
 - [23] Y. R. Liu and M. Oka, $\Lambda_c N$ bound states revisited, *Phys. Rev. D* **85**, 014015 (2012).
 - [24] M. Tanabashi *et al.* (Particle Data Group), Review of particle physics, *Phys. Rev. D* **98**, 030001 (2018).
 - [25] R. Machleidt, K. Holinde, and C. Elster, The Bonn meson exchange model for the nucleon nucleon interaction, *Phys. Rep.* **149**, 1 (1987).