Since the BESIII Collaboration reported a discovery of $Z_c(3900)$ in $e^+e^-\rightarrow \pi^+\pi^-J/\psi$ 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 as 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, 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 $e^+e^-\rightarrow (D^{*0}\bar{D}^-_s/D^{*+}\bar{D}^-_{s0})K^+$ and observed a structure $Z_{c*}(3885)$ in the $K^+$ recoil-mass spectrum when $\sqrt{s} = 4.681$ GeV [8], whose pole position is

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

$$\Gamma = 12.8^{+3.2}_{-4.4} \pm 3.0 \text{ MeV}.$$  

The observation of this $Z_{c*}(3885)$ obtains widely discussions [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 [$c\bar{c}s\bar{s}$]. 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 current existing $Z_c(3885)$.

As shown in Fig. 1, the $Z_{c*}(3885)$ is a little above the $D^{*0}\bar{D}^-_{s0}$ and $D^{*0}\bar{D}^-_{s0}$ thresholds (right panel), compared to the positions of the $Z_c(3885)$ and $Z_c(4025)$ (left panel), we can easily see its similarity with the $Z_c(3885)$ and $Z_c(4025)$. Since the $Z_c(3885)$ and $Z_c(4025)$ have often been assigned as hidden-charm hadronic molecules or resonances [13–16], one may propose whether the newly $Z_{c*}(3885)$ can be a strange hidden-charm tetraquark resonance. In Ref. [17], the mass of the $c\bar{c}q\bar{s}$ was predicted as 3982 MeV within the relativistic diquark-antidiquark picture. In this work, we study the $D^{*0}\bar{D}^{*0}_{s1}$ 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 known, 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 $Z_{cs}$ isn’t fixed yet, the different spin-parities of $I(J^{P}) = 1/2(1^+, 0^-, 1^-, 2^-)$ for the coupled $D^{*0}\bar{D}^{*0}_{s1}$ systems will be discussed.

This paper is organized as follows. After this introduction, we present the $D^{*0}\bar{D}^{*0}_{s1}$ 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.
II. INTERACTIONS

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

$$
\mathcal{L} = g_\sigma \left( H_a^{(0)} \sigma H_a^{(0)} \right) + ig \left( H_b^{(0)} \gamma_\mu A_\mu^{ab} \gamma_5 H_a^{(0)} \right) + g_\sigma \left( \bar{H}_a^{(0)} \sigma H_a^{(0)} \right) + ig \left( \bar{H}_b^{(0)} \gamma_\mu A_\mu^{ab} \gamma_5 H_b^{(0)} \right),
$$

where the subpotentials are

$$
\mathcal{V}_{\text{OBE}} = \begin{pmatrix}
(D^0 D_s^- | V_s | D^0 D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) \\
(D^0 D_s^- | V_s | D^0 D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) \\
(D^0 D_s^- | V_s | D^0 D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-)
\end{pmatrix},
$$

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.,

$$
\mathcal{V}_{E_{h_1 h_2 \rightarrow h_3 h_4}(q)} = -\frac{M(h_1 h_2 \rightarrow h_3 h_4)}{V^2(q^2, m^2_E)},
$$

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

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

$$
\mathcal{V}_{E_{h_1 h_2 \rightarrow h_3 h_4}(q)} = \begin{pmatrix}
(D^0 D_s^- | V_s | D^0 D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) \\
(D^0 D_s^- | V_s | D^0 D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) \\
(D^0 D_s^- | V_s | D^0 D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-) & (D^0 D_s^- | V_s | D^{*0} D_s^-)
\end{pmatrix}.
$$

where the subpotentials are

$$
\mathcal{V}_{\eta}^{D^0 D_s^- \rightarrow D_s^0 D_s^-} = \frac{g_\eta^2}{f_\pi^2} \left[ (\epsilon_1 \epsilon_2^i \epsilon_3^j) Z(\Lambda, m_{\eta_1}, r) \right],
$$

$$
\mathcal{V}_{\eta}^{D^0 D_s^- \rightarrow D_s^- D_s^-} = \frac{g_\eta^2}{f_\pi^2} \left[ (\epsilon_1 \epsilon_2^i \epsilon_3^j) Z(\Lambda, m_{\eta_1}, r) \right],
$$

$$
\mathcal{V}_{\eta}^{D^{*0} D_s^- \rightarrow D_s^0 D_s^-} = \frac{g_\eta^2}{f_\pi^2} \left[ (\epsilon_1 \epsilon_2^i \epsilon_3^j) Z(\Lambda, m_{\eta_1}, r) \right],
$$

$$
\mathcal{V}_{\eta}^{D^{*0} D_s^- \rightarrow D_s^- D_s^-} = \frac{g_\eta^2}{f_\pi^2} \left[ (\epsilon_1 \epsilon_2^i \epsilon_3^j) Z(\Lambda, m_{\eta_1}, r) \right].
$$
In Eqs. (6)-(12), for conciseness, we introduce many variables, and their definitions are

$$
\begin{align*}
\Lambda_1^2 &= \Lambda^2 - q_1^2, \\
m_{n1}^2 &= m_n^2 - q_1^2, \\
q_1 &= \frac{(m_D^2 + m_{D'}^2) - (m_{D2}^2 + m_{D'}^2)}{2(m_D^2 + m_{D'}^2)}, \\
\Lambda_2^2 &= \Lambda^2 - q_2^2, \\
m_{n2}^2 &= m_n^2 - q_2^2, \\
\Lambda_3^2 &= \Lambda^2 - q_3^2, \\
m_{n3}^2 &= m_n^2 - q_3^2, \\
q_2 &= \frac{m_{D2}^2 - m_{D'}^2}{2(m_D^2 + m_{D'}^2)}, \\
q_3 &= \frac{m_D^2 - m_{D'}^2}{2(m_D^2 + m_{D'}^2)}.
\end{align*}
$$

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 $D_{1}^{*}O_{D_{1}}^{*}/D_{0}^{*}D_{0}^{*-}/D_{0}^{*}D_{0}^{*-}$ systems, i.e., $\langle 2S+1L_{j}|O|2S+1L_{j}\rangle$. For the $J^{P} = 1^{+}$ case, the spin-orbit wave functions $\langle 2S+1L_{j}|O|2S+1L_{j}\rangle$ are $D_{0}^{*}D_{1}^{*}|3S_{1}, 3D_{1}\rangle$, $D_{0}^{*}D_{0}^{*}|3S_{1}, 3D_{1}\rangle$, and $D_{0}^{*}D_{1}^{*}|3S_{1}, 3D_{1}\rangle$. Therefore, the matrix elements for all the operators are

$$
\begin{align*}
&\begin{pmatrix}
\epsilon_{1} \cdot \epsilon_{1} \cdot \\
\epsilon_{2} \cdot \epsilon_{2} \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{1} \times \epsilon_{2}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{2} \times \epsilon_{2}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{1} \times \epsilon_{2}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{2} \times \epsilon_{1}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{1} \times \epsilon_{1}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{2} \times \epsilon_{1}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{1} \times \epsilon_{2}) \cdot \\
\frac{1}{\sqrt{2}} (\epsilon_{2} \times \epsilon_{2}) \cdot \\
\end{pmatrix}
\end{align*}
\rightarrow (I),
\tag{13}
$$

$$
\begin{align*}
S(\hat{r}, \epsilon_{1}, \epsilon_{j}) - \sqrt{2} S(\hat{r}, \epsilon_{2} \times \epsilon_{1}, \epsilon_{j})
\end{align*}
\rightarrow \begin{pmatrix}
0 & -\sqrt{2} \\
-\sqrt{2} & 1
\end{pmatrix},
\tag{14}
$$

where $I$ 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 $\Lambda$ 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 $D_{0}^{*}D_{1}^{*-}$ system with several typical cutoff values. Here, we find that the interaction from the $\sigma$ exchange is weak attractive, and it becomes stronger with the increasing of cutoff $\Lambda$. It is obvious that this interaction isn’t stronger enough to form a bound state in the current cutoff choices.

As above mentioned, a repulsive barrier plays an important role in generating a resonance. However, as we can see from Fig. 2, there doesn’t exist any barrier in the attractive interaction from the $\sigma$ exchange, thus, the $S$ wave $D_{0}^{*}D_{1}^{*-}$ cannot bind as a scattering state. So, we further introduce the $S-D$ wave mixing since it may be helpful to generate a $D_{0}^{*}D_{1}^{*-}$ 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, comparing to the weak attractive $\sigma$ exchange interaction, the repulsive centrifugal force is too much stronger. Therefore, we can give an ambitious estimation that there cannot exist a $D_{0}^{*}D_{1}^{*-}$ scattering state with $I(J^{P}) = 1/2(1^{+})$ when we only consider the single $D_{0}^{*}D_{1}^{*-}$ 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 $D_{0}^{*}D_{1}^{*-}/D_{0}^{*}D_{0}^{*-}/D_{0}^{*}D_{1}^{*-}$ systems. As is well known, when the phase shift satisfies $\delta = (n + 1/2)\pi$ with $n = 0, 1, 2, \ldots$, the cross section $\sigma(E)$ takes the maximum value $\sigma(E_0)$. It corresponds to a resonance, its mass and width are

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

respectively.

Finally, we cannot find any possible resonances for the $S$ wave single $D_{0}^{*}D_{1}^{*-}$ when cutoff $\Lambda$ is taken from 1.0 to 2.0 GeV. If the $S-D$ wave mixing is considered, unfortunately, we cannot find any resonant evidences for the single $D_{0}^{*}D_{1}^{*-}$ and $D_{0}^{*}D_{0}^{*-}$ systems either. Therefore, there doesn’t exist possible $D_{0}^{*}D_{1}^{*-}$ and $D_{0}^{*}D_{0}^{*-}$ resonance candidates from the single channel analysis. Here, we can conclude that the newly $Z_{2}$ is excluded as a $D_{0}^{*}D_{1}^{*-}$ or $D_{0}^{*}D_{0}^{*-}$ shape-type resonance.

In the following, we further consider the coupled channel effect to check whether the newly $Z_{2}$($3985$) can be a possible Feshbach-type resonance from the $D_{0}^{*}D_{1}^{*-}/D_{0}^{*}D_{1}^{*-}/D_{0}^{*}D_{1}^{*-}$ interactions. In Fig. 3, we present the OBE effective potentials for the $D_{0}^{*}D_{1}^{*-} \rightarrow D_{0}^{*}D_{1}^{*-}$, $D_{0}^{*}D_{1}^{*-}(D_{0}^{*}D_{1}^{*-}) \rightarrow D_{0}^{*}D_{1}^{*-}$, and $D_{0}^{*}D_{1}^{*-} \rightarrow D_{0}^{*}D_{1}^{*-}$ processes with $\Lambda = 1.5$ GeV. Here,
and the corresponding matrix elements for the spin-spin interaction are
the OBE effective potentials from $D^0D_s^− \rightarrow D^0D_s^−$ scattering process has the same numerical values with those in the $D^0D_s^− \rightarrow D^0D_s^−$ process. Compared to the $\sigma$ exchange, the $\eta$ 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 $\delta$ for the coupled $D^0D_s^−/D^0D_s^−/D^0D_s^−$ systems with $I(J^P) = 1/2(1^+)$. As shown in Fig. 4, with the increasing of the cutoff $\Lambda$, the OBE interactions become much stronger, which leads to the absolute values of the corresponding phase shifts for all channels become much larger. Also, small peaks appear at the $D^0D_s^−$ threshold in the phase shifts for all the $D^0D_s^-$ 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 $D^0D_s^-$ threshold 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 $D^0D_s^−/D^0D_s^−/D^0D_s^−$ Fechbash-type resonance.

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

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

\[
J^P = 0^− : D^0D_s^−|P_0\rangle, D^0D_s^−|P_3\rangle, D^0D_s^−|P_6\rangle, \\
J^P = 1^− : D^0D_s^−|P_1\rangle, D^0D_s^−|P_4\rangle, D^0D_s^−|P_7\rangle, \\
J^P = 2^− : D^0D_s^−|P_2\rangle, D^0D_s^−|P_5\rangle, D^0D_s^−|P_8\rangle,
\]

and the corresponding matrix elements for the spin-spin interaction and tensor force operators are summarized in Eq. (16)-(17).

In recent decades, people pay more and more attention on the study of multiquark states. In particular, the observations of the charged charmonium-like structures such as $Z_c(3900)^±, Z_c(4020)^±$ and $Z_c(4430)$ provide important evidences of the existences of tetraquark state in the charm sector. Very recently, the BESIII Collaboration reported a new charged charmonium-like state $Z_{cs}(3895)$ in in the $K^+$ recoil-mass spectrum of the $e^+e^- \rightarrow (D^0D_s^-/D^0D_s^-)K^+$ processes, which could be the first candidate of the charged hidden-charm tetraquark state with strangeness.

Stimulated by the observation of the strange charmonium-like state $Z_{cs}(3895)$ and its behavior of being closed to threshold, in this work, we perform a systematic investigation on the $D^0D_s^-/D^0D_s^-/D^0D_s^-$ interactions by using the OBE model. In this situation, only the intermediate range interaction from the $\sigma$ and $\eta$ exchanges contributes to the coupled $D^0D_s^-/D^0D_s^-/D^0D_s^-$ system, and the $\eta$ exchange effective potential is dominant. These are remarkable different with the $D^{(*)}_s$ interactions since where the long range force from the $\pi$ exchange and the short range force from the vector meson $\rho/\omega$ exchanges are allowed, in particular, the $\pi$ exchange 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 $D^0 D^+_s/D^0 D^-_s/D^0 D^-_s$ systems, finally, our results exclude the newly $Z_{c5}(3985)$ 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 $D^0 D^-_s/D^0 D^-_s/D^0 D^-_s$ interactions, and find that the newly $Z_{c5}$ cannot be interpreted as a $D^0 D^-_s$-type resonance with $I(J^P) = 1/2(1^-, 2^-)$.

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