**Introduction.**—Very recently, the BESIII Collaboration reported a novel structure $Z_{cs}(\ell 985)^-$ in the $K^+$ recoil-mass spectrum near the $D_s^+D^{*0}/D_s^{*+}D^0$ mass thresholds in the processes of $e^+e^-\to K^+(D_s^+D^{*0}+D_s^{*+}D^0)$ at $\sqrt{s} = 4.681$ GeV [1]. The pole mass and width were determined with a mass-dependent-width Breit-Wigner line shape,

$$M_{Z_{cs}}^{\text{Pole}} = (3982.5^{+1.8}_{-2.6} \pm 2.1) \text{ MeV},$$

$$\Gamma_{Z_{cs}}^{\text{Pole}} = (12.8^{+5.3}_{-4.4} \pm 3.0) \text{ MeV},$$

where the first and the second uncertainties are statistical and systematic, respectively. The significance of the resonance hypothesis is estimated to be $5.3\,\sigma$ over the pure contributions from the conventional charmed mesons.

The minimum quark constituents of $Z_{cs}(\ell 985)$ are $(c\bar{c} s\bar{s} n)$, where $n$ represents the $u/d$ quark. While the number of exotic states is rapidly growing (see Refs. [2–7] for recent reviews), $Z_{cs}(\ell 985)$ is still a very unusual state by current standards. Most XYZ states are isospin singlet, in which the numbers of constituent quark are not fixed. The unquenched quark dynamics [8, 9] would mix the two quark components with four quark components. However, the charged $Z_c/Z_b$ states [10–12], the $P_c$ states [13, 14] and the newly observed $Z_{cs}(\ell 985)$ are multiquark states without much doubt. In addition, $Z_{cs}(\ell 985)$ might be the rare hidden charm exotic candidate with strange number. Another candidate is the $P_c$ states reported recently by LHCb Collaboration [15].

The $Z_{cs}(\ell 900)$ and $Z_{cs}(\ell 4020)$ states are above the threshold of $D^+D_s^*/D^0D^*$ and $D^*D^*$ by several MeVs, respectively. $Z_{cs}(\ell 4020)$ states are likely the heavy quark spin symmetry (HQSS) partners of the $Z_{cs}(\ell 900)$ states. The theoretical interpretations of $Z_{cs}(\ell 900)$ and $Z_{cs}(\ell 4020)$ range from the threshold effect [16], to compact tetraquark states [17, 18], or hadronic molecules [19, 20]. The threshold effect picture of $Z_c$ states was challenged by Guo et al. [21] and JPAC Collaboration [22]. In the tetraquark scheme, it is hard to understand their proximity to the di-meson thresholds. In molecular scenario, the one-pion-exchange interaction for $I = 1 D^{(*)}D^{(*)}$ systems is mainly repulsive. Therefore, theorists resorted to the coupled-channel calculation to interpret $Z_c$ states as the molecular-type resonances, virtual states or bound states [23–35].

In the history, the successful prediction of $\Omega^-$ taught us the importance of SU(3)$_F$ symmetry [36] in hadron spectroscopy. In the exotic hadron sector, within a dominant short-range interaction from SU(3)$_F$ symmetry, we predicted the strange hidden charm pentaquark state as the $Z_cD^*$ bound state with a mass $4456.9$ MeV [37]. Recently, our prediction was supported by the observation of $P_{cs}(4459)^+$ by LHCb collaboration [15]. The $Z_c(\ell 900)$ and $Z_{cs}(\ell 985)$ states are in the proximity of the threshold of $DD^*/D^0D$ and $D_s^*D^*/D^*D$, respectively. It is natural to conjecture that they belong to the same SU(3)$_F$ multiplet. Therefore, it is crucial to investigate $Z_c(\ell 900)/Z_{cs}(\ell 4020)$ and the newly observed $Z_{cs}(\ell 985)$ state in a unified framework with SU(3)$_F$ symmetry and HQSS. In this Letter, we first make use of the HQSS to interpret the $Z_c(\ell 900)$ and $Z_{cs}(\ell 4020)$ as resonances in $J/\psi\pi, DD^*/D^0D$ and $D^*D^*$ coupled-channel calculation. Then, within the SU(3)$_F$ symmetry, we extend the calculation to the strange channels without unknown parameters. We aim at obtaining the mass and width of $Z_{cs}$ state and its HQSS partner state. Meanwhile, the $Z_{cs}$ was firstly observed in $D_s^*D^*/D^*D_s$ channel rather than in the hidden channels like $J/\psi K$. Another question addressed in this Letter is to determine the dominant decay channels of $Z_{cs}(\ell 985)$ and its possible HQSS.
partners.

\[ U/V\text{-spin partners of } Z_c(3900) \pm \text{ and } Z_c(4020) \pm. \]

The quantum numbers of \( Z_c(3900) \) are \( J^{G}(J^{PC}) = 1^+ (1^{-+}) \) \((C\text{ parity only for the neutral states here and below})\) [38]. For the S-wave \( \bar{D}^* D/\bar{D} D^* \) channel, we could construct two orthogonal basis vectors,

\[
\frac{1}{\sqrt{2}} \left( |\bar{D} D^*\rangle + \eta |D^* D\rangle \right),
\]

where \( J^{PC} = 1^{\pm \mp} \) for \( \eta = \mp 1 \). We omit the isospin information in Eq. (2). For the \( I = 1 \) channels, the \( G\)-parity (eigenvalue of \( \hat{G} = \hat{C} e^{i\pi} \)) is \( \eta \). Thus, the \( Z_c(3900) \) states correspond to the isovector channel with \( \eta = +1 \) in Eq. (2). The quantum numbers of \( Z_c(4020) \) are \( J^{G}(J^{PC}) = 1^+ (??^-) \) [38]. As the HQSS partner states of \( Z_c(3900) \), \( Z_c(4020) \) states will couple with the S-wave \( D^* D^* \) isovector channel, which implies its possible \( J^P \) could be \( 0^+, 1^+, 2^+ \). We will assume the \( J^P \) of \( Z_c(4020) \) is \( 1^+ \) and the reason will be given later.

We assume \( Z_{cs}(3985) \) state is the strange partner of \( Z_c(3900) \) in the \( SU(3)_F \) symmetry. To be specific, the \( Z_{cs} \) states are related to the \( Z_c \) states with the rotation in \( U/V\)-spin space as shown in Fig. 1.

\[
Z_{cs}^+ \leftrightarrow Z_{cs}^\Lambda, \quad Z_{cs}^- \leftrightarrow Z_{cs}^0.
\]

U-spin and V-spin are the \( SU(2) \) subgroups of the \( SU(3) \) group just like the isospin subgroup. The \( SU(2) \) doublets for these subgroup are

\[
u, d (I); \quad d, s (U); \quad u, s (V).
\]

The thresholds of \( D^{-} D^{*0} \) (3975 MeV) and \( D^{*-} D^0 \) (3977 MeV) are very close. In the heavy quark limit, \( \bar{D} D^* \) and \( \bar{D}^* D \) are degenerate. We construct the basis of the di-meson channel \( D_s^* D_s^* / \bar{D}_s D_s \) like Eq. (2),

\[
|G_U = \eta \rangle = \frac{1}{\sqrt{2}} \left( |D_s^* D_s^0 \rangle + \eta |D_s^* D^0 \rangle \right),
\]

\[
|G_U = \eta \rangle = \frac{1}{\sqrt{2}} \left( |D_s^* D_s^+ \rangle + \eta |D_s^* D_s^- \rangle \right),
\]

where \( \hat{G}_U \) and \( \hat{G}_V \) transformations are defined like \( \hat{G} \),

\[
\hat{G}_U = \hat{C} e^{iV_2}, \quad \hat{G}_V = \hat{C} e^{iV_2}.
\]

The di-meson channels with \( \eta = +1 \) correspond to \( Z_{cs}^- \) and \( Z_{cs}^0 \) with \( G_{U/V} = +1 \). Similarly, we construct the \( D_s^* D_s^* \) di-meson channels with \( J^P = 1^+ \) and \( G_{U/V} = +1 \), which are the HQSS partner channels of Eq. (5) with \( \eta = +1 \). These channels correspond to the HQSS partner states of \( Z_{cs} \) and \( U/V\)-spin partner states of \( Z_c(4020) \), which are named as \( Z_{cs}^\Lambda \) here and below.

The dynamics of the \( \bar{D}^{(*)}_s D^{(*)}_s \) di-meson systems are constrained by both \( SU(3)_F \) symmetry and HQSS. In the heavy quark limit, the \( c \) and \( \bar{c} \) are the spectators in the di-meson systems. For the S-wave channel, \( \bar{u}_s \) and \( t_1 \cdot t_2 \) are the only interaction operators in spin space, where \( \bar{u}_s \) is the unit operator in spin space and \( t_1 \) is the light spin operator of the heavy meson. In the \( SU(3)_F \) symmetry, the interaction operators in flavor space between light degrees of freedom are \( \mathbb{I}_F \) and \( \mathbb{C}_2 \), where \( \mathbb{I}_F \) is the unit operator and \( \mathbb{C}_2 = -\sum_{i} \lambda_i \lambda_i^2 \) is the Casimir operator. Therefore, the general interaction for \( \bar{D}^{(*)}_s D^{(*)}_s \) could be parameterized as,

\[
V_{q\bar{q}} = c_1 + c_2 t_1 \cdot t_2 + c_3 \mathbb{C}_2 + c_4 (t_1 \cdot t_2) \mathbb{C}_2.
\]

In the \( SU(3)_F \) symmetry, the \( \bar{D}^{(*)}_s D^{(*)}_s \) systems could be classified by \( 3_F \otimes 3_F \rightarrow 8_F \oplus 1_F \) as shown in Fig. 1. The matrix elements of the Casimir operator \( \mathbb{C}_2 \) read:

\[
\langle \mathbb{C}_2 \rangle_{s_F} = \frac{2}{3}; \quad \langle \mathbb{C}_2 \rangle_{1_F} = -\frac{16}{3}.
\]

Thus, in the \( SU(3)_F \) symmetry, the \( \bar{D}^{(*)}_s D^{(*)}_s \) interactions for the \( Z_c(3900) \) and \( Z_c(4020) \) are the same as those of \( \bar{D}^{(*)}_s D^{(*)}_s \) concerning with the \( Z_{cs} \) states.

In the spin space, we write the spin wave function of \( \bar{D}^{(*)}_s D^{(*)}_s \) as \( |l_i h_i (j_1) l_2 h_2 (j_2) J M \rangle \), where \( i, h_i \) and \( j_1 \) are the light spin, heavy spin and total spin of the heavy meson with label \( i \). \( J \) and \( M \) are the total spin and its third component of the di-meson system. One can use the \( 9j \) symbols to relate the above spin wave function to state \( |l_1 l_2 (L) h_1 h_2 (H) J M \rangle \), where \( L \) and \( H \) are total light spin and total heavy spin, respectively. The matrix elements of \( t_1 \cdot t_2 \) can be calculated,

\[
\langle l_1 \cdot l_2 \rangle^{0^+\pm}_{\{PP,WW\}} = \begin{pmatrix} 0 & \frac{\sqrt{7}}{4} \\ \frac{\sqrt{7}}{4} & -\frac{1}{2} \end{pmatrix},
\]

\[
\langle l_1 \cdot l_2 \rangle^{1^+\pm}_{\{PV\eta=\pm1,WW\}} = \begin{pmatrix} -\frac{1}{4} & -\frac{1}{2} \\ -\frac{1}{2} & \frac{1}{4} \end{pmatrix},
\]

\[
\langle l_1 \cdot l_2 \rangle^{1^+\pm}_{\{PV\eta=-1,WW\}} = \begin{pmatrix} 1 \frac{1}{4} \end{pmatrix}, \quad \langle l_1 \cdot l_2 \rangle^{2^+}_{\{WW\}} = \frac{1}{4},
\]

where \( P \) and \( V \) denote the pseudoscalar and vector heavy mesons, receptively. We use the superscript to denote the \( J^P \) of the di-meson channel. For the PV channels, we write the \( \eta \)
of Eqs. (2) and (5) explicitly. Eqs. (9)-(11) are the results in the heavy quark limit. We can see that among the \( VV \) channels, the \( J^P = 1^+ \) one has the same matrix elements with the \( PV \eta = +1 \) channel, which couples with the \( Z_c(3900) \). Thus, it is reasonable to assume the \( J^P = 1^+ \) for \( Z_c(4020) \). The matrix elements in spin space of channels corresponding to \( Z_c(3900)/Z_c(3985) \) and \( Z_c(4020)/Z_c' \) are equal.

The di-meson channels corresponding to \( Z_c(3900), Z_c(4020), Z_c(3985) \) and \( Z_c' \) have the same interaction, which is the result of the \( SU(3)_F \) symmetry and HQSS. We embed these symmetries in Eq. (7). One could adopt other equivalent approaches like superfield method in Ref. [39].

**Coupled-channel calculation.**—We list the channels considered in Table I. Apart from the \( \bar{D}^+(s)D(s) \) channels, we also include the \( J/\psi\pi \) channel for the \( Z_c(3900)/Z_c(4020) \) systems and \( J/\psi K \) channel for the \( Z_c(3985)/Z_c' \) systems. The coupled-channel \( T \)-matrix can be obtained by solving the Lippmann-Schwinger equations (LSEs):

\[
T_{ij} = V_{ij} + \sum_k V_{ik}G_kT_{kj},
\]

where \( G_i \) reads [40],

\[
G_i(E) = \int \frac{d^3l}{(2\pi)^3} \frac{w_{11} + w_{12}}{E^2 - (w_{11} + w_{12})^2 + i\epsilon},
\]

where \( w_{ia} = (p^2 + m_{ia}^2)^{1/2} \) and \( m_{ia} \) is the mass of the \( a \)-th particle in the channel \( i \). We take a hard cutoff \( \Lambda_i \) to regulate the integral. We vary the cutoff parameters \( \Lambda_2 = \Lambda_3 = 0.5 - 1.0 \) GeV but keep the same \( \Lambda_2 \) and \( \Lambda_3 \) to avoid the unintentional HQSS breaking effect. For the definiteness, we fix \( \Lambda_1 = 1.5 \) GeV. For the \( Z_c \) and \( Z_c' \) systems, we choose the same cutoff parameters to keep the \( SU(3)_F \) symmetry.

Following the pionless effective field theory [41, 42], we only introduce the contact interaction. For the off-diagonal potential \( V_{23} = V_{32} \), we take the leading order contact interaction as a constant \( v_{23} \). For the diagonal potential, we have \( V_{22} = V_{33} \) from Eq. (10). In order to obtain the resonances above thresholds, we introduce the next-to-leading order contact interaction for the elastic potential [31],

\[
V_{22} = V_{33} = C_d + \frac{g_i(E)}{2} \left( p^2 + p'^2 \right),
\]

where \( p \) and \( p' \) are the initial and final momenta in the center-of-mass system (c.m.s). The general terms at the next-to-leading order are \( (p + p')^2 \) and \( (p - p')^2 \), while the \( p \cdot p' \) term vanishes after partial wave expansion for the \( S \)-wave.

When the particles are on-shell, the magnitude \( p_i = |p_i| \) of channel \( i \) in c.m.s is

\[
p_i(E) = \sqrt{[E^2 - (m_{i1} + m_{i2})^2][E^2 - (m_{i1} - m_{i2})^2]} \frac{2E}{2E}.
\]

The elastic interaction for \( J/\psi\pi \) or \( J/\psi K \) is purely gluonic vander Waals force, which is known to be tiny [43–45]. We neglect the diagonal interaction in the first channel, \( V_{11} = 0 \). The processes \( D^{(*)}D^{(*)} \rightarrow J/\psi\pi \) and \( D_s^{(*)}D^{(*)} \rightarrow J/\psi K \) are related by the \( U/V \)-spin transformation. Thus the \( V_{13} \) for strange systems and non-strange systems are the same. In the heavy quark limit, the channel 2 and channel 3 have the same spatial wave function and flavor wave function, thus we focus on the spin wave function. The ratio \( V_{12}/V_{13} \) could be estimated by ratio of spin wave function overlaps.

\[
\frac{V_{12}}{V_{13}} = \frac{(J/\psi\pi |PV\eta = +1, 1^+, \text{spin})}{(J/\psi\pi |VV, 1^+, \text{spin})} = 1.
\]

With Eq. (16), we can parameterize the \( V_{12} \) and \( V_{13} \) with one single coupling constant \( v_{12} \).

With the HQSS and \( SU(3)_F \) symmetries, the \( V_{ij} \) reads

\[
V_{ij} = \begin{bmatrix}
0 & v_{12} & v_{12} \\
v_{12} & C_d + \frac{g_i(E)}{2} \left( p^2 + p'^2 \right) & v_{23} \\
v_{12} & v_{23} & C_d + \frac{g_i(E)}{2} \left( p^2 + p'^2 \right)
\end{bmatrix},
\]

where \( C_d \) is the loop function in Eq. (13).

We have four unknown coupling constants, \( v_{12}, v_{23}, C_d \) and \( C_d' \). We shall solve the LSEs and fit the masses and widths of \( Z_c(3900) \) and \( Z_c(4020) \) to determine the four coupling constants. The resonances are located in the unphysical Riemann sheet which is accessed by analytical continuation [40, 46]. We replace \( G_i(E) \) with

\[
C_d^i(E) = G_d^i(E) + \frac{p_i(E)}{4\pi E},
\]

where \( G_d^i \) is the loop function in Eq. (13).

Since the widths of these resonances are narrow, \( \Gamma \ll M \), we could estimate the partial decay widths with the Breit-Wigner parameterization [40]. The \( T_{ij}(E) \) matrix reads

\[
T_{ij}(E) = \frac{1}{2M_R E - M_R + i\frac{\Gamma_R}{2}},
\]

where \( M_R \) and \( \Gamma_R \) are the mass and width of the resonance, respectively. \( g_i \) is the coupling vertex of the resonance and particles in channel \( i \). The partial decay width \( \Gamma_i \) reads,

\[
\Gamma_i = \int \frac{1}{2M_R} |M_{R \rightarrow i}|^2 2\pi\delta(M_R - E_{i1} - E_{i2}) \frac{d^3p_{i1}}{(2\pi)^3 2E_{i1} 2E_{i2}}.
\]

where \( M_{R \rightarrow i} = g_i \). We make the substitution in the narrow-width approximation,

\[
2\pi\delta(M_R - E_{i1} - E_{i2}) \rightarrow 2\text{Im} \frac{1}{M_R - E_{i1} - E_{i2} - i\frac{\Gamma_R}{2}}.
\]
We change the integral variable to \( E \) and the partial wave decay width becomes:

\[
\Gamma_i = -\frac{1}{16\pi^2} \int_{m_{1}+m_{2}}^{\infty} \frac{dE \rho_{11}(E)}{E^2} 4M_i \text{Im} T_{ii}(E). \tag{22}
\]

In practical calculation, we integrate in \( E \) around 2 widths up and down the pole mass (if allowed by the lower limit) to obtain \( \Gamma_i \), since we find \( \sum_i \Gamma_i \approx \Gamma_H \) in this integration range.

**Numerical results and Discussions.**—We choose the recent results of the charged \( Z_c(3900) \) and \( Z_c(4020) \) in Refs. [47, 48] as input. We could either choose the averaged results in Ref. [38], which would give the similar final results. We determine the coupling constants in either \( \Lambda_{2/3} = 1.0 \text{ GeV} \) or 0.5 GeV and then calculate the masses and widths of \( Z_{cs} \) and \( Z_{cs}' \). We present \( T_{11} \)-matrix in the unphysical sheet with \( \Lambda_{2/3} = 1.0 \text{ GeV} \) in Fig. 2. We can see two poles corresponding to \( Z_c(3900) \) and \( Z_c(4020) \), which are barely above the thresholds of \( D\bar{D}^* \) and \( D^*\bar{D}^* \) by several MeVs, respectively. The positions of poles are \( M - i\Gamma/2 \), where \( M \) and \( \Gamma \) are the mass and width of resonances.

We give the numerical results in Table II. One can see that we could reproduce the mass and width of the newly observed \( Z_{cs}(3985) \) state with \( Z_c(3900) \) and \( Z_c(4020) \) as input. Our results are in good agreement with the experimental results, which strongly supports that isospin doublet \( Z_{cs}(3985) \) states are the \( U/V \)-spin partner of the charged \( Z_c(3900) \) as resonances. Meanwhile, we predict a new resonance above the threshold of \( D\bar{D} \) by 8 MeV with a width about 30 MeV, which is the HQSS partner of the \( Z_{cs}(3985) \) and \( U/V \)-spin partner of \( Z_c(4020) \).

In this calculation, we use the decay modes \( J/\psi \pi(K), D(s)D^*/D^*(s)D, D^*(s)D^* \) to saturate the total widths, which would bring some uncertainties. These uncertainties would be compensated in ratios of partial decay widths. The coupling constants \( v_{12} \) is very small and thus the resonances are dominated by the \( D(s)D^*/D^*(s)D \) and \( D^*(s)D^* \) components. With the central value of \( v_{12} \), the partial wave decay widths of \( \Gamma_1 \) are very small, thus we take the upper limit of \( v_{12} \) to give the lower limit of \( \Gamma_i/\Gamma_1 \). From Table II, we can see that the decay process to \( J/\psi \pi(K) \) is suppressed by at least one order compared with the open charmed final state decays. The dominant \( D(s)D^*/D^*(s)D \) decay modes lead to the observation in these channels in experiment [1]. In Eq. (5), the \( Z_{cs}(3985) \) states with \( G_{V/U} = 1 \) have the same components of \( D\bar{D}^* \) and \( D^*\bar{D} \), which is constrained by the \( SU(3)_F \) symmetry.

We have assigned seven states \( Z_{cs}^{0/\pm}, Z_{cs}^{0/+}, Z_{cs}^{0/-} \) into the \( SU(3) \) octet in Fig. 1. The left isospin singlet in \( S_F \) representation might mix with the isospin singlet in \( 1_F \) like the \( \phi \) and \( \omega \) mesons. The \( C \) matrix elements read

\[
\langle C_2 \rangle_{s\bar{s}} = -\frac{4}{3}, \quad \langle C_2 \rangle_{(u\bar{u}+d\bar{d})}/\sqrt{2} = -\frac{10}{3}.
\]

Both matrix elements have different signs from those of the octet in ideal \( SU(3) \) symmetry. Therefore, the mixture would make the eighth state disappear. In the compact tetraquark scheme, the existence of the tetraquark states do not depend on the flavor. Searching for the eighth state would help to distinguish the compact tetraquark states from the di-meson states.

We further extend the calculation to the hidden bottom sector with heavy quark flavor symmetry (HQFS). We use the coupling constants determined with \( \Lambda_{2/3} = 0.5 \text{ GeV} \) and obtain two poles \((M, \Gamma)\) in the non-strange channel,

\[
(10612.0, 32.2) \text{ MeV}, \quad (10656.9, 32.3) \text{ MeV}, \quad (23)
\]

which correspond to the \( Z_c(10610) \) and \( Z_c(10650) \). We also predict two strange hidden bottom states \( Z_{bs} \) near the \( B_sB^*/B^*_s\bar{B} \) and \( B^*_sB^* \) threshold respectively,

\[
(10699.9, 32.3) \text{ MeV}, \quad (10747.9, 32.2) \text{ MeV}. \tag{24}
\]

The resonance \( Z_{bs} \) could be searched for in the \( B_sB^*, B^*_s\bar{B}, \bar{B}K \) final states and \( Z_{bs} \) in the \( B^*_sB^*, B_sB^*, B^*_sB, \bar{B}K \) final states.

**Summary and Outlook.**—In summary, we perform the \( J/\psi \pi(K), D(s)D^*/D^*(s)D, D^*(s)D^* \) coupled-channel calculation in the contact interaction with the \( SU(3)_F \) symmetry and HQSS. We fit the masses and widths of \( Z_{cs}(3900) \) and \( Z_c(4020) \) to determine the coupling constants. We reproduce the mass and width of \( Z_{cs}(3985) \) very well as the \( U/V \)-spin partner of \( Z_c(3900) \). We also obtain the ratios of the partial decay widths of \( Z_{cs}(3985) \) and obtain its dominant \( D\bar{D}^*/D^*(s)D \) decay modes. We introduce the \( G_{V/U} \) parity to label \( Z_{cs} \) states. In the \( SU(3)_F \) limit, the partial decay
widths to $\bar{D}_s D^*$ and $\bar{D}_s^* D$ are equal. We also predict the $Z'_cs$ with the mass around 4130 MeV, which are the HQSS partner states of the $Z_{cs}(3985)$ and $U/V$-spin partner states of the $Z_{s}(4020)$. With the HQFS, we reproduce the masses and widths of $Z_{cs}(10610)$ and $Z_{s}(10650)$ and predict their $G_{U/V}$ partner states $Z_{bs}$ and $Z'_{bs}$ with a mass around 10700 MeV and 10750 MeV.

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1. M. Ablikim et al. (BESIII), (2020), arXiv:2011.07855 [hep-ex].
2. N. Brambilla, S. Eidelman, C. Hanhart, A. Nefediev, C.-P. Shen, C. E. Thomas, A. Vairo, and C.-Z. Yuan, Phys. Rept. 873, 1 (2020).
3. Y.-R. Liu, H.-X. Chen, W. Chen, X. Liu, and S.-L. Zhu, Prog. Part. Nucl. Phys. 107, 237 (2019).
4. F.-K. Guo, C. Hanhart, U.-G. Meißen, Q. Wang, Q. Zhao, and B.-S. Zou, Rev. Mod. Phys. 90, 015004 (2018).
5. S. L. Olsen, T. Skwarnicki, and D. Zieminska, Rev. Mod. Phys. 90, 015003 (2018).
6. H.-X. Chen, W. Chen, X. Liu, and S.-L. Zhu, Phys. Rept. 639, 1 (2016).
7. A. Esposito, A. Pilloni, and A. Polosa, Phys. Rept. 668, 1 (2017).
8. T. Barnes and E. Swanson, Phys. Rev. C 77, 055206 (2008).
9. E. J. Eichten, K. Lane, and C. Quigg, Phys. Rev. D 69, 094019 (2004).
10. A. Bondar et al. (Belle), Phys. Rev. Lett. 108, 122001 (2012).
11. M. Ablikim et al. (BESIII), Phys. Rev. Lett. 110, 252001 (2013).
12. Z. Liu et al. (Belle), Phys. Rev. Lett. 110, 252002 (2013); [Erratum: Phys.Rev.Lett. 111, 019901 (2013)].
13. R. Aaij et al. (LHCb), Phys. Rev. Lett. 122, 222001 (2019).
14. R. Aaij et al. (LHCb), Phys. Rev. Lett. 115, 072001 (2015).
15. M. Wang, in Implications of LHCb measurements and future prospects (2020) on behalf of the LHCb.
16. E. S. Swanson, Phys. Rev. D 91, 034009 (2015).
17. J. Dias, F. Navarra, M. Nielsen, and C. Zanetti, Phys. Rev. D 88, 016004 (2013).
18. S. Agaev, K. Azizi, and H. Sundu, Phys. Rev. D 96, 034026 (2017).
19. F.-K. Guo, C. Hidalgo-Duque, J. Nieves, and M. P. Valderrama, Phys. Rev. D 88, 054007 (2013).
20. J. He, X. Liu, Z.-F. Sun, and S.-L. Zhu, Eur. Phys. J. C 73, 2635 (2013).
21. F.-K. Guo, C. Hanhart, Q. Wang, and Q. Zhao, Phys. Rev. D 91, 051504 (2015).
22. A. Pilloni, C. Fernandez-Ramirez, A. Jackura, V. Mathieu, M. Mikhasenko, J. Nys, and A. Szczepaniak (JPAC), Phys. Lett. B 772, 200 (2017).
23. P. G. Ortega, J. Segovia, D. R. Entem, and F. Fernández, Eur. Phys. J. C 79, 78 (2019).
24. T. Chen, Y. Chen, M. Gong, C. Liu, L. Liu, Y.-B. Liu, Z. Liu, J.-P. Ma, M. Werner, and J.-B. Zhang (CLQCD), Chin. Phys. C 43, 103103 (2019).
25. Y. Ikeda, S. Aoki, T. Doki, S. Gongyo, T. Hatsuda, T. Inoue, T. Iritani, N. Ishii, K. Murano, and K. Sasaki (HAL QCD), Phys. Rev. Lett. 117, 242001 (2016).
26. S.-H. Lee, C. DeTar, D. Mohler, and H. Na (Fermilab Lattice, MILC), (2014), arXiv:1411.1389 [hep-lat].
27. S. Prelovsek, C. Lang, L. Leskovec, and D. Mohler, Phys. Rev. D 91, 014504 (2015).
28. F. Aceti, M. Bayar, E. Oset, A. Martinez Torres, K. Khemchandani, J. M. Dias, F. Navarra, and M. Nielsen, Phys. Rev. D 90, 016003 (2014).
29. J. He, Phys. Rev. D 92, 034004 (2015).
30. J. He and D.-Y. Chen, Eur. Phys. J. C 78, 94 (2018).
31. M. Albaladejo, F.-K. Guo, C. Hidalgo-Duque, and J. Nieves, Phys. Lett B 755, 337 (2016).
32. C. Hanhart, Y. S. Kalashnikova, P. Matsuakhe, R. Mizuk, A. Nefediev, and Q. Wang, Phys. Rev. Lett. 115, 202001 (2015).
33. F.-K. Guo, C. Hanhart, Y. S. Kalashnikova, P. Matsuakhe, R. Mizuk, A. Nefediev, Q. Wang, and J. L. Wynen, Phys. Rev. D 93, 074031 (2016).
34. Q. Wang, V. Baru, A. Filin, C. Hanhart, A. Nefediev, and J.-L. Wynen, Phys. Rev. D 98, 074023 (2018).
35. B. Wang, L. Meng, and S.-L. Zhu, (2020), arXiv:2009.01980 [hep-ph].
[36] V. Barnes et al., Phys. Rev. Lett. 12, 204 (1964).
[37] B. Wang, L. Meng, and S.-L. Zhu, Phys. Rev. D 101, 034018 (2020).
[38] P. Zyla et al. (Particle Data Group), PTEP 2020, 083C01 (2020).
[39] J. Nieves and M. Valderrama, Phys. Rev. D 86, 056004 (2012).
[40] J. Oller and E. Oset, Nucl. Phys. A 620, 438 (1997), [Erratum: Nucl.Phys.A 652, 407–409 (1999)].
[41] E. Epelbaum, H.-W. Hammer, and U.-G. Meissner, Rev. Mod. Phys. 81, 1773 (2009).
[42] R. Machleidt and D. Entem, Phys. Rept. 503, 1 (2011).
[43] K. Yokokawa, S. Sasaki, T. Hatsuda, and A. Hayashigaki, Phys. Rev. D 74, 034504 (2006).
[44] L. Liu, H.-W. Lin, and K. Orginos, PoS LATTICE2008, 112 (2008).
[45] X.-H. Liu, F.-K. Guo, and E. Epelbaum, Eur. Phys. J. C 73, 2284 (2013).
[46] J. Nieves and E. Ruiz Arriola, Phys. Rev. D 64, 116008 (2001).
[47] M. Ablikim et al. (BESIII), Phys. Rev. Lett. 112, 132001 (2014).
[48] M. Ablikim et al. (BESIII), Phys. Rev. D 92, 092006 (2015).