Identifying $D_{sJ}(2700)$ through its decay modes

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We study how to assign the recently observed $D_{sJ}(2700)$ meson to an appropriate level of the $c\bar{s}$ spectrum by the analysis of its decay modes in final states comprising a light pseudoscalar meson. We use an effective Lagrangian approach with heavy quark and chiral symmetries, obtaining that the measurement of the $D^* K$ decay width would allow to distinguish between two possible assignments.

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I. INTRODUCTION

Recently, Belle Collaboration collected an important new piece of information on charm spectroscopy from the first Dalitz plot analysis of the decay process $B^+ \rightarrow \bar{D}^0 D^0 K^+$. The Dalitz plot shows an accumulation of events at $M^2(D^0 D^0) = 16 - 18$ GeV$^2$ and $M^2(D^0 K^+) = 7 - 8$ GeV$^2$. This could be interpreted as the overlap of two contributions, a (horizontal) band in the $\bar{D}^0 D^0$ channel due to a state possibly identified with $\psi(4160)$, and a (vertical) band in the $D^0 K^+$ channel due to a $c\bar{s}$ state which does not correspond to any of the previously observed mesons with this quark content. A horizontal band at $M^2(D^0 D^0) \approx 14.2$ GeV$^2$ is also present, due to the production of $\psi(3770)$.

Information from the Dalitz plot was enriched by the study of the individual invariant mass distributions. The $M^2(D^0 K^+)$ distribution shows the presence of a resonance, $D_{sJ}(2700)$, with parameters:

$$M = 2708 \pm 9^{+11}_{-10} \text{ MeV}$$
$$\Gamma = 108 \pm 23^{+36}_{-31} \text{ MeV}.$$  \hspace{1cm} (1.1)

Moreover, from the distribution in the helicity angle $\theta$, the angle between the $D^0$ momentum and the opposite of the kaon momentum in the $D^0 D^0$ rest frame, it was possible to assign the spin-parity $J^P = 1^-$ to $D_{sJ}(2700)$. A measurement of the branching fractions

$$B(B^+ \rightarrow \bar{D}^0 D^0 K^+) = (22.2 \pm 2.2) \times 10^{-4}$$
$$B(B^+ \rightarrow D_{sJ}(2700) D^0) \times B(D_{sJ}(2700) \rightarrow D^0 K^+) = (11.3 \pm 2.2) \times 10^{-4}$$  \hspace{1cm} (1.2)

was also carried out.

Belle’s observation is just the last one of a series of discoveries of open and hidden charm hadrons which have greatly enriched the charm spectroscopy in the past few years. In case of open charm, which $D_{sJ}(2700)$ belongs to, before the B-factories era the known $c\bar{s}$ spectrum consisted of four states only: the pseudoscalar $D_s(1686)$ and vector $D^*_s(2112)$ mesons, corresponding to the $s$-wave states of the quark model, and the axial-vector $D_s(2536)$ and tensor $D_{s2}(2573)$ mesons, $p$-wave states. In 2003, two narrow resonances: $D_{sJ}(2317)$ and $D_{sJ}(2460)$ were discovered by BaBar and CLEO Collaborations, respectively, and later on confirmed by other experiments. They were assigned spin-parities $J^P = 0^+$ and $J^P = 1^+$. Their identification as proper $c\bar{s}$ states has been controversial since then; however, they have the quantum numbers of the two states needed to complete the $p$-wave multiplet, and their radiative decays occur accordingly, so that their interpretation as ordinary $c\bar{s}$ configurations is natural.

Last year, BaBar Collaboration announced the observation of another $c\bar{s}$ meson, $D_{sJ}(2860)$ decaying to $D^0 K^+$ and $D^+ K_S$, with mass $M(D_{sJ}(2860)) = 2856.6 \pm 1.5 \pm 5.0$ MeV and width $\Gamma(D_{sJ}(2860)) = 47 \pm 7 \pm 10$ MeV. The quantum numbers of this meson have still to be assigned, and two proposals have been put forward. One is that this state is the first radial excitation of $D_{sJ}(2317)$, hence with $J^P = 0^+$, the other one is that the meson has $J^P = 3^-$, an interpretation proposed in [11] and supported by a lattice QCD study [12]. In the same analysis of the $D K$ mass distribution, BaBar noticed a broad structure with $M = 2688 \pm 4 \pm 3$ MeV and $\Gamma = 112 \pm 7 \pm 36$ MeV, likely the same resonance $D_{sJ}(2700)$ found by Belle.

In this scenario, $D_{sJ}(2700)$ needs to be properly placed in the spectrum of mesons with charm and strangeness. In our study we investigate its decay modes with the aim of finding signatures allowing to assign the meson to a particular $c\bar{s}$ level. In this respect we follow a different strategy from the one adopted in the framework of quark models [10], where the assignment is made by comparing the observed mass to the predicted value in a suitably chosen interquark potential, a procedure which can find difficulties in the cases where neglected effects, such as the threshold effects, are relevant.

In Section IV we present a theoretical framework based upon an effective Lagrangian describing strong decays of heavy mesons to final states comprising a light pseudoscalar meson, and displaying heavy quark and chiral
symmetries. We analyze the various decay modes and compare the predictions that follow from different assignments. The implications for the still unobserved states, including $c\bar{q}$ mesons, are discussed in Sections III-V, then our conclusions are presented.

II. HEAVY MESONS DECAYS TO LIGHT PSEUDOSCALAR MESONS

The study of properties and decays of hadrons containing a single heavy quark $Q = c, b$ is suitably carried out in the $m_Q \rightarrow \infty$ limit, which is formulated in the Heavy Quark Effective Theory [13]. In such a limit, the heavy quark acts as a static colour source for the rest of the hadron; its spin $s_Q$ is decoupled from the total angular momentum of the light degrees of freedom $\vec{s}$, and they are separately conserved. Hadrons can be classified according to the values of $s_\ell$ and of the total spin $J = s_Q + s_\ell$. In particular, heavy mesons can be organized in doublets, each one corresponding to a particular value of $s_\ell$ and parity; the members of each doublet differ for the orientation of $s_Q$ with respect to $s_\ell$ and, in the heavy quark limit, are degenerate. Mass degeneracy is broken at order $1/m_Q$.

For $Q\bar{q}$ states one can write $s_\ell = s_\gamma + \ell$, where $s_\gamma$ is the light antiquark spin and $\ell$ is the orbital angular momentum of the light degrees of freedom relative to the heavy quark. The lowest lying $Q\bar{q}$ mesons correspond to $\ell = 0$ (s-wave states of the quark model) with $s_\ell^P = 1^-$ and $J^P = (0^-,1^-)$. For $\ell = 1$ (p-wave states of the quark model), it could be either $s_\ell^P = 1^+$ or $s_\ell^P = 3^+$, the two corresponding doublets having $J^P = (0^+,1^+)$ and $J^P = (1^+,2^+)$. The mesons with $\ell = 2$ (d-wave states) are collected either in the $s_\ell^P = 3^-$ doublet, consisting of states with $J^P = (1^-,2^-)$, or in the $s_\ell^P = 5^-$ doublet with $J^P = (2^-,3^-)$ states. And so on.

The two states $D_s(1686)$ and $D_s^*(2112)$ can be identified with the members of the lowest lying $s_\ell^P = 1^-$ doublet. The resonances $D_{s1}(2536)$ and $D_{s2}(2573)$, together with $D_{s1J}(2317)$ and $D_{s2J}^*(2460)$, fill the four p-wave levels: in particular, $D_{s2}(2573)$ corresponds to $s_\ell^P = 1^+$, $J^P = 2^+$ state, while $D_{s1J}(2317)$ to $s_\ell^P = 1^+$, $J^P = 0^+$. The two axial-vector mesons $D_{s1}(2536)$ and $D_{s2J}(2460)$ could be assigned to a linear combination of $s_\ell^P = 3^-$ and $s_\ell^P = 1^-$ states, which is allowed at $O(1/m_Q)$; however, in case of non-strange charm mesons such a mixing has been measured and found to be small [14,15], so that we can identify $D_{s1J}(2536)$ and $D_{s2J}^*(2460)$ with the $J^P = 1^+$ $s_\ell^P = 4^-$ and $s_\ell^P = 1^+$ states, respectively.

In the interpretation in [14] $D_s(2860)$ corresponds to the $J^P = 3^-$ component of the $s_\ell^P = 5^-$ doublet, so that the still unassigned levels are the two states belonging to the $s_\ell^P = 3^-$ doublet, with $J^P = (1^-,2^-)$, and the partner of $D_{sJ}(2860)$ in the $s_\ell^P = 5^-$ doublet, which has spin two.

The classification can be continued for higher values of $s_\ell^P$ to describe high spin mesons; moreover, it can be replicated for radial excitations, and it is expected to hold as far as $O(1/m_Q)$ effects are small.

Since the spin-parity of $D_{sJ}(2700)$ has been determined: $J^P = 1^-$, the state fits either in the doublet with $s_\ell^P = 1^-$ or in the one with $s_\ell^P = 3^-$. However, a $1^-$ state belonging to the $s_\ell^P = 1^-$ doublet is already known, $D_s^*(2112)$, so that in this case $D_{sJ}(2700)$ would be a radial excitation. Combinations of the two cases are not allowed in the heavy quark limit; a possible role of $1/m_Q$ effects will be discussed below. Therefore, two possibilities must be considered:

\begin{itemize}
  \item $D_{sJ}(2700)$ belongs to the doublet with $s_\ell^P = 1^-$ and is the first radial excitation; we denote this state as $D_s^{*\ell}$;
  \item $D_{sJ}(2700)$ is the low lying state with $s_\ell^P = 3^-$, denoted as $D_s^{*1}$.
\end{itemize}

These two possibilities can be conveniently analyzed adopting a formalism which represents the various doublets by effective fields being 4 $\times$ 4 matrices [10]. The two doublets of interest, with $s_\ell^P = 1^-$ denoted by $H_a$, and with $s_\ell^P = 3^-$ denoted by $X_a$ ($a$ is a light flavour index) are given by:

\begin{equation}
H_a = \frac{1 + \gamma_5}{2} \left[ P_{a\nu} \gamma^\mu - P_{a\gamma5} \right]
\end{equation}

\begin{equation}
X_a^{\mu} = \frac{1 + \gamma_5}{2} \times \left\{ P_{2a\nu} \gamma_5 \gamma^\nu - P_{1a\nu} \sqrt{\frac{3}{2}} \left[ g^{\mu\nu} - \frac{3}{2} (\gamma^\mu v - \gamma^\nu v) \right] \right\}
\end{equation}

with $v$ the meson four-velocity. Notice that $H_a$ describes the fundamental $s_\ell^P = 1^-$ doublet; the doublet corresponding to the first radial excitations is described by an identical structure $H_a$. The various operators in (2.1) annihilate mesons of four-velocity $v$ which is conserved in strong interaction processes: the heavy field operators contain a factor $\sqrt{m_P}$ and have dimension 3/2.

The interaction of these heavy mesons with the octet of light pseudoscalar mesons, introduced using the fields $\xi = e^{i\eta}$, with the matrix $M$ containing $\pi, K$ and $\eta$ fields:

\begin{equation}
M = \begin{pmatrix}
\sqrt{\frac{1}{8}} \pi^0 + \sqrt{\frac{1}{8}} \eta & \pi^+ & K^+
\pi^- & -\sqrt{\frac{1}{8}} \pi^0 + \sqrt{\frac{1}{8}} \eta & \pi^0 - \frac{1}{2} \eta \\
K^0 & K^0 & \eta
\end{pmatrix}
\end{equation}

and $f_\pi = 132$ MeV, can be described by an effective Lagrangian invariant under chiral transformations of the
light fields and heavy-quark spin-flavour transformations of the heavy fields. At the leading order in the heavy quark mass and light meson momentum expansion the decays $F \to HM$ ($F = H'$ and $X$, and $M$ a light pseudoscalar or vector meson) can be described by the Lagrangian interaction terms:

\[ \mathcal{L}_{H'} = \frac{g}{\Lambda_{\chi}} Tr[\bar{H}_a H_b \gamma_{\mu} \gamma_5 A^{\mu}_{a b}] \]

\[ \mathcal{L}_X = \frac{k'}{\Lambda_{\chi}} Tr[\bar{H}_a X^b \{i d_{\mu} A^\mu + i D A_{a b} \} \gamma_5] + h.c. \]

where $A_{a b} = \frac{1}{2} (\partial^\mu \bar{q}_a - \partial^\mu q_a^3)$ and $D$ is the covariant derivative $D_{a b} = -\delta_{a b} \partial_\mu + V_{a b}$ with $V_{a b} = \frac{1}{2} (\partial^\mu \bar{q}_a - \partial^\mu q_a^3)$. $\Lambda_{\chi}$ is the chiral symmetry-breaking scale and $\tilde{g}$, $k'$ are effective couplings. We use $\Lambda_{\chi} = 1$ GeV, while at present the value of the couplings is not known neither from experiment nor from theoretical considerations.

The widths of the allowed decay modes of $D_{sJ}(2700)$ for the two possible interpretations $D_{s1}^{\ast}$ or $D_{s1}^{\ast\ast}$ can be computed using (2.3) and (2.4). Since a heavy vector state can decay to a light pseudoscalar and a heavy pseudoscalar or vector meson, we consider the modes: $D_{sJ}(2700) \to D^+ K_S$, $D^0 K^+$, $D_s \eta$ and $D_{sJ}(2700) \to D^+ K_S$, $D^{0*} K^+$, $D^{\ast\ast}_s \eta$, and compute the ratios of decay widths:

\[ R_1 = \frac{\Gamma(D_{sJ} \to D^+ K)}{\Gamma(D_{sJ} \to D K)} \]

\[ R_2 = \frac{\Gamma(D_{sJ} \to D_s \eta)}{\Gamma(D_{sJ} \to D K)} \]

\[ R_3 = \frac{\Gamma(D_{sJ} \to D^*_s \eta)}{\Gamma(D_{sJ} \to D K)} \]

(with $D^{(*)} K = D^{(*)0} K_S + D^{(*)+} K^+$) for both the assignments of $s_J^P$ to $D_{sJ}(2700)$. These ratios are useful to discriminate between the two assignments; moreover, the method is sensible since in the ratios, the dependence on the (unknown) effective couplings drops out, and the predictions are model independent.

Let us first identify $D_{sJ}(2700)$ with $D_{s1}^{\ast\ast}$. Using the effective Lagrangian (2.3), the $D_{s1}^{\ast\ast}$ decay widths can be written as follows:

\[ \Gamma(D_{s1}^{\ast\ast} \to D_s P_a) = C_P \frac{\tilde{g}^2}{6\pi f_Z^2} \frac{M_{D_{s1}^{\ast\ast}}}{M_{D^{\ast\ast}_s}} |q|^3 \]

\[ \Gamma(D_{s1}^{\ast\ast} \to D^*_s P_a) = C_P \frac{\tilde{g}^2}{3\pi f_Z^2} \frac{M_{D_{s1}^{\ast\ast}}}{M_{D^{\ast\ast}_s}} |q|^3 \]

where the light flavour index $a = u, d, s$ identifies $D_{s1}^{\ast\ast} = D^{(*)0}$, $D^{(*)+}$, $D_{s1}^{\ast\ast}$, respectively, and $P_a$ represents a light pseudoscalar meson with quark content $s\bar{a}$ ($P_a = K^+$, $K_S(L), \eta$). $C_P$ is a coefficient depending on the $P$ meson: $C_{K^+} = 1$, $C_{K_S} = \frac{1}{2}$, and $C_\eta = \frac{2}{3}$; the modulus of the three momentum $q$ reads: $|q| = \lambda^{1/2}(M_{D^{\ast\ast}_s}^2, M_{D^{(*)}_s}^2, M_{P_a}^2)/2M_{D^{\ast\ast}_s}$.

On the other hand, if $D_{sJ}(2700)$ is identified with $D_{s1}^{\ast}$ from the Lagrangian (2.4) we obtain:

\[ \Gamma(D_{s1}^{\ast} \to D_s P_a) = \]

\[ C_P \frac{16}{9\pi f_Z^2} \left( \frac{k'}{\Lambda_{\chi}} \right)^2 \frac{M_{D_{s1}^{\ast}}}{M_{D^{\ast}_s}} |P_a|^2 |q|^3 \]

\[ \Gamma(D_{s1}^{\ast} \to D^*_s P_a) = \]

\[ C_P \frac{2}{9\pi f_Z^2} \left( \frac{k'}{\Lambda_{\chi}} \right)^2 \frac{M_{D_{s1}^{\ast}}}{M_{D^{\ast}_s}} |P_a|^2 |q|^3 \]

These two sets of expressions produce different values for the ratios (2.5), as one can appreciate considering Table II where the numerical results are collected (with the errors obtained considering the uncertainty in the $D_{sJ}(2700)$ mass).

| $R_1 \times 10^2$ | $R_2 \times 10^2$ | $R_3 \times 10^2$ |
|-----------------|-----------------|-----------------|
| $D_{s1}^{\ast\ast}$ | 91 ± 4 | 20 ± 1 | 5 ± 2 |
| $D_{s1}^{\ast}$ | 4.3 ± 0.2 | 16.3 ± 0.9 | 0.18 ± 0.07 |

The ratios $R_1$ and $R_3$ are very different if $D_{sJ}(2700)$ is $D_{s1}^{\ast\ast}$ or $D_{s1}^{\ast}$, so that the measurements of these ratios allow to properly identify the $D_{sJ}(2700)$. Since $R_3$ is small, the $D^+ K$ decay mode is the main signal that must be investigated in order to distinguish between the two possible assignments for $D_{sJ}(2700)$.

It is worth noticing that the results in Table II are different from the ones obtained in ref. [19] using the $3P_0$ model, a quark model with harmonic oscillator meson wave functions, where it turns out that the $DK$ mode is suppressed if $D_{sJ}(2700)$ is identified with $D_{s1}^{\ast\ast}$. Moreover, in the $D_{s1}^{\ast}$ identification, more reduction of the $D^+ K$ signal than obtained in ref. [19] is expected.

From the ratios in Table II assuming that the width of $D_{sJ}(2700)$ is saturated by decay modes with a heavy meson and a light pseudoscalar meson in the final state, we can determine the coupling constants appearing in (2.3) and (2.4), in correspondence of the two possible assignments for $D_{sJ}(2700)$. This assumption is reasonable, since decay modes with more than one light pseudoscalar meson, or a light vector meson in the final states are severely phase-space suppressed.

Identifying $D_{sJ}(2700)$ with $D_{s1}^{\ast\ast}$ we obtain:

\[ \tilde{g} = 0.26 \pm 0.05 \]

while if $D_{sJ}(2700)$ is $D_{s1}^{\ast}$ we obtain

\[ k' = 0.14 \pm 0.03 \]

These two values are similar to the results obtained for analogous coupling constants appearing in the effective heavy quark chiral Lagrangians. Using the results (2.8), (2.9) the branching fractions of the various decay
modes can be computed; the results are collected in Table II. The errors in eqs. (2.8) and (2.9), as well as those in Table II, are obtained from the uncertainties in \( M_{D_{sJ}} \) and \( \Gamma(D_{sJ}) \).

A comment is in order about the accuracy of the results. The effective Lagrangians (2.3)–(2.4) coincide with the first terms of an expansion in the light pseudoscalar meson momenta. Since in the decays we have considered, such momenta are not very small, one should in principle add other terms, which should be weighted by new, unknown coupling constants. Analogously, corrections to the heavy quark limit, which is also used in (2.3)–(2.4), could be considered, by adding \( \mathcal{O}(1/m_Q) \) terms, which would contain new unknown constants as well [15]. We cannot assess the role and the size of such corrections on general grounds, however we expect that they would largely cancel out in the ratios of widths. On the other hand, a mixing between the two \( J^P = 1^- \), \( s_\ell^P = \frac{1}{2}^- \), \( \frac{3}{2}^- \) states, which is possible at \( \mathcal{O}(1/m_Q) \), would involve a mixing angle that could be fixed by measuring the ratios in Table II and comparing the experimental results with the ratios computed in the heavy quark limit.

## III. PARTNERS WITHOUT STRANGENESS

It is interesting to study the charmed mesons with the same quantum numbers as \( D_{sJ}(2700) \), but with a different light quark flavour. These states are charged charmed mesons and a neutral one, denoted as \( D_{sJ}^+ \) and \( D_{sJ}^0 \), respectively. They have not been observed yet, so that their masses are unknown. We fix such masses to \( 2600 \pm 50 \) MeV by the reasonable assumption that \( D_{sJ}(2700) \) is heavier by an amount of the size of the strange quark mass.

Allowed decay modes for \( D_{sJ}^+(2600) \) are: \( D_{sJ}^+ \to D^0\pi^+, \ D_{sJ}^+\eta \), and \( D_{sJ}^+ \to D^{*0}\pi^+, \ D^{*+}\pi^0, \ D^{*+}\eta \), while for \( D_{sJ}^0 \) they are: \( D_{sJ}^0 \to D^{*0}\pi^+, \ D^{*0}\pi^0, \ D_{sJ}^0\eta \) and \( D_{sJ}^0 \to D^{*0}\pi^+, \ D^{*+}\pi^-, \ D^{*+}\eta \); the corresponding widths are obtained using eq. (2.6)–(2.7), and depend on the possible identification of \( D_{sJ}^{+(0)} \). The states having \( s_\ell = \frac{1}{2}^- \) are denoted as \( D_{sJ}^{+(0)} \) and are radial excitations, while the states having \( s_\ell = \frac{3}{2}^- \) are denoted as \( D_{sJ}^{+(0)} \).

Using the effective coupling constants \( \tilde{g} \) and \( k' \) in (2.8), (2.9), we obtain:

\[
\Gamma(D_{sJ}^{+(0)}) = (128 \pm 61) \text{ MeV} \quad (3.1)
\]
\[
\Gamma(D_{sJ}^{+(0)}) = (85 \pm 46) \text{ MeV} \quad (3.2)
\]

so that the \( c\bar{q} \) partners have widths which are different in the case of the two assignments, although with a sizeable uncertainty. Since the mesons are not very broad, it should be possible to observe them. The predicted branching fractions, collected in Table III, confirm that the two assignments produce different results. In the identification with the state \( D^{*'} \), the mode \( D^{*'} \to D^{*+}\pi^- \) has the largest branching fraction, while in the second hypothesis, i.e. \( D_{sJ} = D_{sJ}^+ \), the mode with the largest branching ratio is \( D_{sJ}^+ \to D^{*+}\pi^- \).

## IV. SPIN PARTNERS

Since in the heavy quark limit the heavy mesons are collected in doublets with a definite value of \( s_\ell^P \), the state \( D_{sJ}(2700) \) has a partner from which it differs only for the value of the total spin.

The partner of \( D_{sJ}^{*'} \) (\( s_\ell^P = \frac{1}{2}^- \)) has \( J^P = 0^- \); it is denoted \( D_{sJ}' \), the first radial excitation of \( D_{sJ} \). On the other hand, the partner of \( D_{sJ}^{*1} \) (\( s_\ell^P = \frac{3}{2}^- \)) has \( J^P = 2^- \); we refer to this state as to \( D_{sJ}^{*2} \). In both cases, the decay modes \( D_{sJ}' \to D^{*0} K^+, \ D_{sJ}^{*2} \to D^{*0} K^0 \), \( D_{sJ}^{*2} \to D^{*+} \eta \), are permitted. Using the effective Lagrangians (2.3)–(2.4) we find:

\[
\Gamma(D_{sJ}' \to D^{*}\pi) = \mathcal{C}_P \bar{q}^2 \frac{M_{D_{sJ}^*}}{2k'_{\pi}^2 \bar{q}} |\bar{q}|^3
\]
\[
\Gamma(D_{sJ}^{*2} \to D^{*}\pi) = \mathcal{C}_P \frac{4}{9\pi f_{\pi}^2} \left( \frac{k'}{\Lambda} \right)^2 \left( \frac{M_{D_{sJ}^*}}{M_{D_{sJ}^{*2}}} \right) \left[ M_{D_{sJ}^*} + |\bar{q}|^2 \right] |\bar{q}|^3 .
\]

(4.1)

In the heavy quark limit, these partners are degenerate, hence, in the numerical analysis we assign them the same mass as \( D_{sJ}(2700) \). Using the obtained values for \( \tilde{g} \) and \( k' \), we get:

\[
\Gamma(D_{sJ}') = (70 \pm 30) \text{ MeV}
\]
\[
\Gamma(D_{sJ}^{*2}) = (12 \pm 5) \text{ MeV}
\]

(4.2)

and the branching fractions in Table IV. In the two assignments the spin partners differ for their decay width.

## V. \( D_{sJ}(2700) \) DECAY CONSTANT

We conclude our study observing that, since \( D_{sJ}(2700) \) has been discovered through the production in \( B^+ \to D^{*0} D_{sJ}(2700) \), it is possible to estimate its decay constant \( f_{D_{sJ}} \) defined as:

\[
< 0 |\bar{s} \gamma_\mu (1 - \gamma_5) c |D_{sJ}(p, \epsilon) > = f_{D_{sJ}} M_{D_{sJ}} \epsilon_\mu .
\]

(5.1)

The effective Hamiltonian governing \( b \to c\bar{c}s \) transitions:

\[
H_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{cb} V_{cs}^* \left[ C_1(\mu) O_1 + C_2(\mu) O_2 \right],
\]

(5.2)
In principle, the dependence of the Wilson coefficients $C_i$ on the scale $\mu$ should cancel against the $\mu$-dependence of the matrix elements of the $O_i$. As it is well known, the calculation of matrix elements such as $\langle D_{sJ} | D | H_{e\ell f} | B \rangle$ is a difficult task. The simplest evaluation can be obtained by naive factorization [21], in which the $B^+ \to D^0 D_{sJ}(2700)$ decay amplitude is written as:

$$A(B^+ \to D^0 D_{sJ}(2700)) = a_1 G_F \sqrt{\frac{\pi}{3}} \frac{V_{cs} V_{cx}^*}{V_{ub} V_{ub}^*} \times$$

$$<\bar{D}^0 (v') | \bar{b} \gamma^\mu (1 - \gamma_5) c | B(v) > \times$$

$$<D_{sJ}(p, \epsilon) | \bar{c} \gamma_\mu (1 - \gamma_5) \eta | 0 >$$

with $a_1 = C_1 + C_2/3$, $G_F$ the Fermi constant and $V_{ij}$ elements of the Cabibbo Kobayashi Maskawa mixing matrix; we use $a_1 = 1.1$ together with the values quoted in [22] for $V_{ij}$.

In the heavy quark limit, the first matrix element in (5.4) can be written in terms of the Isgur-Wise function $\xi(v \cdot v')$ [13]:

$$<\bar{D}^0 (v') | \bar{b} \gamma^\mu (1 - \gamma_5) c | B(v) > = \xi(v \cdot v') \sqrt{M_B M_D} (v + v')^\mu$$

with $v$ and $v'$ the four-velocities of the heavy mesons. The linear parametrization:

$$\xi(v \cdot v') = 1 - \rho^2 (v \cdot v' - 1)$$

involves the slope $\rho^2$, for which various determinations are available: $\rho^2 = 0.83^{+0.05}_{-0.01}$ from lattice QCD [23],

$$\rho^2 = 1.179 \pm 0.048 \pm 0.028 \text{ from a recent BaBar measurement [24]},$$

$$\rho^2 = 1.26 \pm 0.16 \pm 0.11 \text{ from a Belle measurement [25]}.\] The Heavy Flavour Averaging Group quotes $\rho^2 = 1.23 \pm 0.05$ for $B \to D^*$ decays, and $\rho^2 = 1.17 \pm 0.18$ from $B \to D$ transitions [26]: we use this last value. Computing $\mathcal{B}(B^+ \to D^0 D_{sJ}(2700))$ by naive factorization, considering the branching fraction $\mathcal{B}(D_{sJ}(2700) \to D^0 K^+)$ in Table III for the two possible assignments to $D_{sJ}(2700)$, and comparing the result to (1.2), we obtain:

$$\mathcal{B}_{D_{sJ}^+} = 243 \pm 41 \text{ MeV}$$

$$\mathcal{B}_{D_{sJ}^*} = 181 \pm 30 \text{ MeV}$$

VI. CONCLUSIONS

We have discussed possible ways to distinguish between two assignments for the state $D_{sJ}(2700)$. Our result is that the decay mode to $D^* K$ has very different branching ratios in the two possible assignments, so that a measurement of such a branching fraction would be useful to shed light on the identification of $D_{sJ}(2700)$.

We have also obtained several predictions, namely the effective couplings governing the strong decays of $D_{sJ}$ to $D(D^*)$ and a light pseudoscalar meson, which are different according to the adopted interpretation. We have also obtained predictions for non strange partners of $D_{sJ}$, for which we also found that the decay to $D^* \pi$ has the largest branching ratio, and predictions for the spin partners. Further research on $D_{sJ}(2700)$ according to these
suggestions would be useful to complete our understanding of the open charm meson spectrum.

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TABLE IV: Branching ratios of the spin partner of $D_s^*$ for the two quantum number assignments.

| $D_s^*$ ($J^P = 0^-$) | $B(D_s^*(D_{s2}^*) \to D^{0}K^+)$ | $B(D_s^*(D_{s2}^*) \to D^{*-}K_S)$ | $B(D_s^*(D_{s2}^*) \to D^*_S)$ |
|------------------------|-----------------|-----------------|-----------------|
| (50.0 ± 0.5)%         | (23.7 ± 0.2)%   | (2.6 ± 0.9)%    |
| (49.8 ± 0.6)%         | (23.6 ± 0.2)%   | (3.1 ± 1.0)%    |

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