Understanding $D_{sJ}(2317)$

P. Colangelo and F. De Fazio

*Istituto Nazionale di Fisica Nucleare, Sezione di Bari, Italy*

**Abstract**

We analyze the hadronic and radiative decay modes of the recently observed $D_{sJ}(2317)$ meson, in the hypothesis that it can be identified with the scalar $s_{1/2}^P$ state of $c\bar{s}$ spectrum ($D_{s0}$). The method is based on heavy quark symmetries and Vector Meson Dominance ansatz. We find that the hadronic isospin violating mode $D_{s0} \rightarrow D_s\pi^0$ is enhanced with respect to the radiative mode $D_{s0} \rightarrow D^*_s\gamma$. The estimated width of the meson is $\Gamma(D_{s0}) \simeq 7$ KeV.
1 Introduction

The BaBar Collaboration has reported the observation of a narrow peak in the $D_s^+ \pi^0$ invariant mass distribution, corresponding to a state of mass 2.317 GeV [1]. The observed width is consistent with the resolution of the detector, thus $\Gamma \leq 10$ MeV. In the same analysis no significant signals are found in the $D_s \gamma$ and $D_s \gamma \gamma$ mass distributions. The meson has been denoted as $D_{sJ}(2317)$; the announcement has immediately prompted different interpretations [2, 3].

A possible quantum number assignment to $D_{sJ}(2317)$ is $J^P = 0^+$, as suggested by the angular distribution of the meson decay with respect to its direction in the $e^+ - e^-$ center of mass frame. This assignment can identify the meson with the $D_{s0}$ state in the spectrum of the $c\bar{s}$ system. Considering the masses of the other observed states belonging to the same system, $D_{s1}(2536)$ and $D_{sJ}(2573)$, the mass of the scalar $D_{s0}$ meson was expected in the range 2.45–2.5 GeV, therefore $\sim 150$ MeV higher than the observed 2.317 GeV. A $D_{s0}$ meson with such a large mass would be above the threshold $M_{DK} = 2.359$ GeV to strongly decay by $S$-wave Kaon emission to $DK$, with a consequent broad width. For a mass below the $DK$ threshold the meson has to decay by different modes, namely the isospin-breaking $D_s \pi^0$ mode observed by BaBar, or radiatively. The $J^P = 0^+$ assignment excludes the final state $D_s \gamma$, due to angular momentum and parity conservation; indeed such a final state has not been observed. On the other hand, for a scalar $c\bar{s}$ meson the decay $D_{s0} \to D^*_s \gamma$ is allowed. However, no evidence is reported yet of the $D_s \gamma \gamma$ final state resulting from the decay chain $D_{s0} \to D^*_s \gamma \to D_s \gamma \gamma$. In order to confirm the identification of $D_{sJ}(2317)$ with the scalar $D_{s0}$, one has at first to understand whether the decay modes of a scalar particle with mass of 2317 GeV can be predicted in agreement with the experimental findings presently available. In particular, the isospin violating decay to $D_s \pi^0$ should proceed at a rate larger than the radiative mode $D_{s0} \to D^*_s \gamma$, though not exceeding the experimental upper bound $\Gamma \leq 10$ MeV. This letter is devoted to such an issue.

2 Mode $D_{s0} \to D_s \pi^0$

In order to analyze the isospin violating transition $D_{s0} \to D_s \pi^0$ one can use a formalism that accounts for the heavy quark spin-flavour symmetries in hadrons containing a single heavy quark, and the chiral symmetry in the interaction with the octet of light pseudoscalar states.

In the heavy quark limit, the heavy quark spin $\vec{s}_Q$ and the light degrees of freedom
total angular momentum $\vec{s}_t$ are separately conserved. This allows to classify hadrons with a single heavy quark $Q$ in terms of $s_t$ by collecting them in doublets the members of which only differ for the relative orientation of $\vec{s}_Q$ and $\vec{s}_t$.

The doublets with $J^P = (0^-, 1^-)$ and $J^P = (0^+, 1^+)$ (corresponding to $s_t^P = \frac{1}{2}^-$ and $s_t^P = \frac{1}{2}^+$, respectively) can be described by the effective fields

$$H_a = \frac{1}{2} \left( 1 + \frac{\gamma^5}{2} \right) \left[ P_{a\mu} \gamma^\mu - P_a \gamma_5 \right]$$

$$S_a = \frac{1}{2} \left( 1 + \frac{\gamma^5}{2} \right) \left[ P_{1a}^{\mu} \gamma_\mu \gamma_5 - P_{0a} \right]$$

where $v$ is the four-velocity of the meson and $a$ is a light quark flavour index. In particular in the charm sector the components of the field $H_a$ are $P_{a}^{(*)} = D_0^{(*)}, D_1^{(*)}, D_2^{(*)}$ (for $a = 1, 2, 3$); analogously, the components of $S_a$ are $P_{0a} = D_0^{0}, D_1^{0}, D_0^{1}$ and $P_{1a} = D_1^{0}, D_1^{1}, D'^{1}_{s1}$.

In terms of these fields it is possible to build up an effective Lagrange density describing the low energy interactions of heavy mesons with the pseudo Goldstone $\pi$, $K$ and $\eta$ bosons [4, 5, 6, 7]:

$$\mathcal{L} = i \, \text{Tr} \{ H_b v^\mu D_{\mu ba} \overline{H}_a \} + \frac{f^2}{8} \text{Tr} \{ \partial^\mu \Sigma \partial_\mu \Sigma^\dagger \} + \text{Tr} \{ S_b ( i \, v^\mu D_{\mu ba} - \delta_{ba} \Delta ) \overline{S}_a \} + i g \text{Tr} \{ H_b \gamma_\mu \gamma_5 A_{\mu ba} \overline{H}_a \} + i g' \text{Tr} \{ S_b \gamma_\mu \gamma_5 A_{\mu ba} \overline{S}_a \} + [ i h \text{Tr} \{ S_b \gamma_\mu \gamma_5 A_{\mu ba} \overline{H}_a \} + \text{h.c.}] \ .$$

(3)

In (3) $\overline{H}_a$ and $\overline{S}_a$ are defined as $\overline{H}_a = \gamma^0 H_1^a \gamma_0$ and $\overline{S}_a = \gamma^0 S_1^a \gamma_0$; all the heavy field operators contain a factor $\sqrt{M_F}$ and have dimension $3/2$. The parameter $\Delta$ represents the mass splitting between positive and negative parity states.

The $\pi$, $K$ and $\eta$ pseudo Goldstone bosons are included in the effective lagrangian (3) through the field $\xi = e^{i \frac{3}{f}}$ that represents a unitary matrix describing the pseudoscalar octet, with

$$\mathcal{M} = \left( \begin{array}{ccc} \sqrt{\frac{2}{3}} \pi^0 + \sqrt{\frac{1}{6}} \eta & \pi^+ & K^+ \\ \pi^- & -\sqrt{\frac{2}{3}} \pi^0 + \sqrt{\frac{1}{6}} \eta & K^0 \\ K^- & K^0 & -\sqrt{\frac{1}{2}} \eta \end{array} \right)$$

(4)

and $f \simeq f_\pi$. In eq.(3) $\Sigma = \xi^2$, while the operators $D$ and $A$ are given by:

$$D_{\mu ba} = \delta_{ba} \partial_\mu + V_{\mu ba} = \delta_{ba} \partial_\mu + \frac{1}{2} \left( \xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger \right)_{ba}$$

(5)

$$A_{\mu ba} = \frac{1}{2} \left( \xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger \right)_{ba}$$

(6)

The strong interactions between the heavy $H_a$ and $S_a$ mesons with the light pseudoscalar mesons are thus governed, in the heavy quark limit, by three dimensionless
couplings: $g$, $h$ and $g'$. In particular, $h$ describes the coupling between a member of the $H_a$ doublet and one of the $S_a$ doublet to a light pseudoscalar meson, and is the one relevant for our discussion.

Isospin violation enters in the low energy Lagrangian of $\pi$, $K$ and $\eta$ mesons through the mass term

$$\mathcal{L}_{\text{mass}} = \frac{\bar{\mu} f^2}{4} \text{Tr} \{ \xi m_q \xi - \xi \xi^\dagger \}$$

with $m_q$ the light quark mass matrix:

$$m_q = \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_d & 0 \\ 0 & 0 & m_s \end{pmatrix}.$$  \hspace{5cm} (8)

In addition to the light meson mass terms, $\mathcal{L}_{\text{mass}}$ contains an interaction term between $\pi^0$ $(I = 1)$ and $\eta$ $(I = 0)$ mesons: $\mathcal{L}_{\text{mixing}} = \frac{\bar{\mu}}{2} \frac{m_d - m_u}{\sqrt{3}} \pi^0 \eta$ which vanishes in the limit $m_u = m_d$. As in the case of $D_s^* \rightarrow D_s \pi^0$ studied in [8], the isospin mixing term can drive the $D_{s0} \rightarrow D_s \pi^0$ transition.\footnote{Electromagnetic contributions to $D_{s0} \rightarrow D_s \pi^0$ are expected to be suppressed with respect to the strong interaction mechanism considered here.} The amplitude $A(D_{s0} \rightarrow D_s \pi^0)$ is simply written in terms of $A(D_{s0} \rightarrow D_s \eta)$ obtained from (3), $A(\eta \rightarrow \pi^0)$ from (7) and the $\eta$ propagator that puts the strange quark mass in the game. The resulting expression for the decay amplitude involves the coupling $h$ and the suppression factor $(m_d - m_u)/(m_s - m_d + m_u)$ accounting for isospin violation, so that the width $\Gamma(D_{s0} \rightarrow D_s \pi^0)$ reads:

$$\Gamma(D_{s0} \rightarrow D_s \pi^0) = \frac{1}{16\pi f^2} \frac{h^2}{M_{D_s}} \left( \frac{m_d - m_u}{m_s - m_d + m_u} \right)^2 (1 + \frac{m_{\pi^0}^2}{|\bar{p}_{\pi^0}|^2}) |\bar{p}_{\pi^0}|^3.$$  \hspace{5cm} (9)

As for $h$, the result of QCD sum rule analyses of various heavy-light quark current correlators is $|h| = 0.6 \pm 0.2$ [6]. Using the central value, together with the factor $(m_d - m_u)/(m_s - m_d + m_u) \simeq \frac{1}{437}$ [9] and $f = f_\pi = 132$ MeV we obtain:

$$\Gamma(D_{s0} \rightarrow D_s \pi^0) \simeq 6 \text{ KeV}.$$  \hspace{5cm} (10)

Eq.(9) can receive $SU(3)_F$ corrections: a hint on their size comes from the use of $f = f_\pi = 171$ MeV instead of $f_\pi$ in (9), which gives $\Gamma(D_{s0} \rightarrow D_s \pi^0) \simeq 4 \text{ KeV}$. On the other hand, we neglect corrections related to the finite charm quark mass.

The analogous calculation for $D_s^* \rightarrow D_s \pi^0$ involves the coupling $g$ in (3). Since $h$ and $g$ have similar sizes $(0.3 \leq g \leq 0.6)$, it turns out that the transitions $D_{s0} \rightarrow D_s \pi^0$ is enhanced with respect to $D_s^* \rightarrow D_s \pi^0$ essentially due to kinematics, being $|\bar{p}_{\pi^0}|^2 |\bar{p}_{\pi^0}|^3 \simeq 3 \times 10^2$.\footnote{The analogous calculation for $D_s^* \rightarrow D_s \pi^0$ involves the coupling $g$ in (3). Since $h$ and $g$ have similar sizes $(0.3 \leq g \leq 0.6)$, it turns out that the transitions $D_{s0} \rightarrow D_s \pi^0$ is enhanced with respect to $D_s^* \rightarrow D_s \pi^0$ essentially due to kinematics, being $|\bar{p}_{\pi^0}|^2 |\bar{p}_{\pi^0}|^3 \simeq 3 \times 10^2$.}
3 Radiative $D_{s0} \to D_s^*\gamma$ decay

Let us now turn to $D_{s0} \to D_s^*\gamma$, the amplitude of which has the form:

$$A(D_{s0} \to D_s^*\gamma) = e \, d \left[ (\epsilon^* \cdot \eta^*)(p \cdot k) - (\eta^* \cdot p)(\epsilon^* \cdot k) \right], \quad (11)$$

where $p$ is the $D_{s0}$ momentum, $\epsilon$ the $D_s^*$ polarization vector, and $k$ and $\eta$ the photon momentum and polarization. The corresponding decay rate is:

$$\Gamma(D_{s0} \to D_s^*\gamma) = \alpha |d|^2 |\vec{k}|^3. \quad (12)$$

The parameter $d$ gets contributions from the photon couplings to the light quark part $e_s \bar{s}\gamma_\mu s$ and to the heavy quark part $e_c \bar{c}\gamma_\mu c$ of the electromagnetic current, $e_s$ and $e_c$ being strange and charm quark charges in units of $e$. Its general structure is:

$$d = d^{(h)} + d^{(\ell)} = \frac{e_c}{\Lambda_c} + \frac{e_s}{\Lambda_s}, \quad (13)$$

where $\Lambda_a$ ($a = c, s$) have dimension of a mass. Such a structure is already known from the constituent quark model. In the case of $M1$ heavy meson transitions, an analogous structure predicts a relative suppression of the radiative rate of the charged $D^*$ mesons with respect to the neutral one [10, 11, 12, 13], suppression that has been experimentally confirmed [14]. From (12,13) one could expect a small width for the transition $D_{s0} \to D_s^*\gamma$, to be compared to the hadronic width $D_{s0} \to D_s\pi^0$ which is suppressed as well.

In order to determine the amplitude of $D_{s0} \to D_s^*\gamma$ we follow a method based again on the use of heavy quark symmetries, together with the vector meson dominance (VMD) ansatz [11, 13]. We first consider the coupling of the photon to the heavy quark part of the e.m. current. The matrix element $\langle D_s^*(v', \epsilon)|\bar{c}\gamma_\mu c|D_{s0}(v)\rangle$ ($v, v'$ meson four-velocities) can be computed in the heavy quark limit, matching the QCD $\bar{c}\gamma_\mu c$ current onto the corresponding HQET expression [15]:

$$J^{HQET}_\mu = \bar{h}_v[v_\mu + \frac{i}{2m_Q}(\overrightarrow{\partial}_\mu - \overleftarrow{\partial}_\mu) + \frac{i}{2m_Q}\sigma_{\mu\nu}(\overrightarrow{\partial}^\nu + \overleftarrow{\partial}^\nu) + \ldots]h_v \quad (14)$$

where $h_v$ is the effective field of the heavy quark. For transitions involving $D_{s0}$ and $D_s^*$, and for $v = v'$ ($v \cdot v' = 1$), the matrix element of $J^{HQET}_\mu$ vanishes. The consequence is that $d^{(h)}$ is proportional to the inverse heavy quark mass $m_Q$ and presents a suppression factor since in the physical case $v \cdot v' = (m_{D_{s0}}^2 + m_{D_s}^2)/2m_{D_{s0}}m_{D_s} = 1.004$. Therefore, we neglect $d^{(h)}$ in (13).
To evaluate the coupling of the photon to the light quark part of the electromagnetic current we invoke the VMD ansatz and consider the contribution of the intermediate \( \phi(1020) \):\[
\langle D_s^*(v', \epsilon)|\bar{s}\gamma_\mu s|D_{s0}(v)\rangle = \sum_{\lambda}<D_s^*(v', \epsilon)|\bar{s}\gamma_\mu s|\phi(k, \epsilon_1(\lambda))|D_{s0}(v)\rangle \frac{i}{k^2 - M_\phi^2} \langle 0|\bar{s}\gamma_\mu s|\phi(k, \epsilon_1(\lambda))\rangle
\]
with \( k^2 = 0 \) and \( \langle 0|\bar{s}\gamma_\mu s|\phi(k, \epsilon_1)\rangle = M_\phi f_\phi \epsilon_{1\mu} \). The experimental value of \( f_\phi \) is \( f_\phi = 234 \text{ MeV} \). The matrix element \( \langle D_s^*(v', \epsilon)|\bar{s}\gamma_\mu s|\phi(k, \epsilon_1(\lambda))|D_{s0}(v)\rangle \) describes the strong interaction of a light vector meson (\( \phi \)) with two heavy mesons (\( D_s^* \) and \( D_{s0} \)). It can also be obtained through a low energy lagrangian in which the heavy fields \( H_a \) and \( S_a \) are coupled, this time, to the octet of light vector mesons.\(^2\) The Lagrange density is set up using the hidden gauge symmetry method [5], with the light vector mesons collected in a \( 3 \times 3 \) matrix \( \hat{\rho} \) analogous to \( M \) in (4). The lagrangian\(^3\) reads as [16]:
\[
\mathcal{L}' = i \hat{\mu} \text{Tr} \{ \bar{S}_a H_b \sigma^{\lambda\nu} V_{\lambda\nu}(\rho)_{ab} \} + h.c.,
\]
with \( V_{\lambda\nu}(\rho) = \partial_\lambda \rho_\nu - \partial_\nu \rho_\lambda + [\rho_\lambda, \rho_\nu] \) and \( \rho_\lambda = i \frac{g\sqrt{2}}{\sqrt{2}} \hat{\rho}_\lambda \), \( g \) being fixed to \( g = 5.8 \text{ GeV}^{-1} \) by the KSRF rule [17]. The coupling \( \hat{\mu} \) in (16) is constrained to \( \hat{\mu} = -0.13 \pm 0.05 \text{ GeV}^{-1} \) by the analysis of the \( D \to K^* \) semileptonic transitions induced by the axial weak current [16].

The resulting expression for \( \frac{1}{\Lambda_s} \) is:
\[
\frac{1}{\Lambda_s} = -4\hat{\mu} \frac{g}{\sqrt{2}} \sqrt{\frac{M_{D_s^*} f_\phi}{M_{D_{s0}} M_\phi}}.
\]

The parameters are obtained from independent channels; we use their central values.

The numerical result for the radiative width:
\[
\Gamma(D_{s0} \to D_s^*\gamma) \simeq 1 \text{ KeV}
\]
shows that the hadronic \( D_{s0} \to D_s^*\pi^0 \) transition is more probable than the radiative mode \( D_{s0} \to D_s^*\gamma \). In particular, if we assume that the two modes essentially saturate the \( D_{s0} \) width, we have
\[
\Gamma(D_{s0}) \simeq 7 \text{ KeV}
\]
and
\[
\mathcal{B}(D_{s0} \to D_s\pi^0) \simeq 0.85
\]
\[
\mathcal{B}(D_{s0} \to D_s^*\gamma) \simeq 0.15
\]
\(^2\)The standard \( \omega_8 - \omega_0 \) mixing is assumed, resulting in a pure \( \bar{s}s \) structure for \( \phi \).
\(^3\)The role of other possible structures in the effective lagrangian contributing to radiative decays is discussed in [13].
at odds with the case of the $D_s^*$ meson, where the radiative mode dominates the decay rate.

The same conclusion concerning the hierarchy of $D_{s0} \to D_s\pi^0$ versus $D_{s0} \to D_s^*\gamma$ is reached in [3] using the quark model. Since our calculation is based on a different method, the $s^P_\ell = \frac{1}{2}^-$ and $s^P_\ell = \frac{1}{2}^+$ doublets being treated as uncorrelated multiplets, we find the agreement noticeable.

4 Conclusions and perspectives

We have found that the observed narrow width and the enhancement of the $D_s\pi^0$ decay mode are compatible with the identification of $D_{sJ}(2317)$ with the scalar state belonging to the $s^P_\ell = \frac{1}{2}^+$ doublet of the $c\bar{s}$ spectrum. However, this conclusion does not avoid other questions raised by the BaBar observation, one being the low mass of the state. We believe that such a particular issue requires additional investigations. A second point is that the radiative mode, although suppressed, is not negligible, and should be observed at a level typically represented by the ratios in (20).

The quantum number assignment has two main and rather straightforward consequences. The first one is the existence of the axial vector partner $D_{s1}'$ belonging to the same spin doublet $s^P_\ell = \frac{1}{2}^+$. Even in the case where the hyperfine splittings between positive and negative parity states are similar: $M_{D_{s1}'} - M_{D_{s0}} \simeq M_{D_s^*} - M_{D_s}$, this meson is below the $D^*K$ threshold. Therefore, its hadronic decay to $D_s^*\pi^0$, at the rate

$$\Gamma(D_{s1}' \to D_s^*\pi^0) = \frac{\hbar^2}{48\pi f^2} \frac{M_{D_s^*}}{M_{D_{s1}'} M_{D_{s1}'}^3} \frac{M_{D_s} - m_s}{m_s} \left[2 + \frac{(M_{D_s}^2 + M_{D_{s1}'}^2 - M_{\pi^0}^2)^2}{4M_{D_s}^2 M_{D_{s1}'}^2}\right]$$

$$\times \left(1 + \frac{m_{D_{s0}}^2}{|\vec{p}_{\pi^0}|^2}\right)|\vec{p}_{\pi^0}|^3 \simeq \Gamma(D_{s0} \to D_s\pi^0)$$

would produce a narrow peak in the $D_s^*\pi^0$ mass distribution. The confirmation of such a state, the existence of which is suggested by the analysis of the $D_s\gamma\pi^0$ mass distribution [1], will support the interpretation.

The second consequence concerns the doublet of scalar and axial vector mesons in the $b\bar{s}$ spectrum. Since the mass splitting between $B$ and $D$ states is similar to the corresponding mass splitting between $B_s$ and $D_s$ states, such mesons should be below the $BK$ and $B^*K$ thresholds, thus producing narrow peaks in $B_s\pi^0$ and $B_s^*\pi^0$ mass
distributions, with rates resulting from expressions analogous to (9)-(21).

**Note added.** When this work was completed, the CLEO Collaboration announced the observation of a narrow resonance with mass 2.46 GeV in the $D^*_s \pi^0$ final state and the confirmation of $D_{sJ}(2317)$ [18]. Moreover, a theoretical analysis based on the quark model was posted on the Los Alamos arXive, with the same conclusions presented here [19].

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