I briefly review the experimental observations concerning the charmed mesons $D_{sJ}^*(2317)$, $D_{sJ}(2460)$ and survey on some of the interpretations proposed in order to understand their nature. I present an analysis of their decay modes in the hypothesis that they can be identified with the scalar and axial vector $s^P = \frac{1}{2}^+$ states of the $c\bar{s}$ spectrum ($D_{s0}^*$, $D_{s1}^*$). The method is based on heavy quark symmetries and Vector Meson Dominance ansatz. Comparison with present data supports the interpretation.

1. Introduction

In April 2003, the BaBar Collaboration reported the observation of a narrow peak in the $D_s^+ n^0$ invariant mass distribution, corresponding to a state of mass 2317 MeV, denoted as $D_{sJ}^*(2317)$\(^1\). The state is produced from charm continuum and the observed width is consistent with the resolution of the detector, $\Gamma \leq 10$ MeV.

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 scalar $D_{s0}^*$ 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-violating $D_s n^0$ 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 $D^*_{s0} \rightarrow D^*_s\gamma$ is allowed. However, no evidence is reported yet of the $D_s\gamma\gamma$ final state resulting from the decay chain $D^*_{s0} \rightarrow D^*_s\gamma \rightarrow D_s\gamma\gamma$.

Later on, in May 2003, CLEO Collaboration confirmed the BaBar observation of $D^*_{sJ}(2317)$ with the same features outlined above; furthermore, CLEO Collaboration observed a second narrow peak, corresponding to a state with mass 2460 MeV decaying to $D^*_s\pi^0$ 2. Again, the width is compatible with the detector resolution. Evidence of this second state was present in the first analysis by BaBar Collaboration, which gave subsequent confirmation of the CLEO observation 3. BELLE Collaboration has confirmed both states 4, observing their production both from charm continuum, both in B decays; more recently also FOCUS Collaboration 5 has detected of a narrow peak at 2323 $\pm$ 2 MeV, slightly above the values obtained by the other three experiments for $D^*_{sJ}(2317)$.

The observation of the decay $D_{sJ}(2460) \rightarrow D^*_s\pi^0$ suggests that $D_{sJ}(2460)$ has $J^P = 1^+$. This assignment is supported also by the observation of the mode $D_{sJ}(2460) \rightarrow D_s\gamma$, forbidden to a $0^+$ state, and by the angular analysis performed by BELLE 6. Such an analysis was carried out for $D_{sJ}(2460)$ produced in B decays and favours the identification of $D_{sJ}(2460)$ with an axial-vector particle. Production of $D_{sJ}(2460)$ in B decays was observed also by BaBar 7. However, as in the case of $D^*_{sJ}(2317)$, the measured mass is below theoretical expectations for the $1^+ c\bar{s}$ state $D'_{sJ}$ and the narrow width contrasts with the expected broadness of the latter.

These peculiar features of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ have prompted a number of analyses, aimed either at refining previous results in order to support the $c\bar{s}$ interpretation of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$, or at explaining their nature in a different context. The various interpretations are reviewed in Ref. 8.

Among the non standard scenarios, it has been often considered the possibility of a sizeable four-quark component in $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$. Four-quark states could be baryonium-like or molecular-like, if they result from bound states of quarks or of hadrons, respectively, and examples of the second kind of states are the often discussed $f_0(980)$ and $a_0(980)$ when interpreted as $K\bar{K}$ molecules. In the molecular interpretation, the $D^*_{sJ}(2317)$ could be viewed as a $D\pi$ molecule 8, an interpretation supported by the fact that the mass 2.317 GeV is close to the $D\pi$ threshold, or as a $D_s\pi$ atom 9. Analogously, $D_{sJ}(2460)$ would be a $D^* K$ molecule. Mixing between ordinary $c\bar{s}$ state and a composite state has also been considered 10. No definite
answer comes from lattice QCD, since, according to Ref. 12, lattice predictions are inconsistent with the simple $q\bar{q}$ interpretation for $D_{sJ}^*(2317)$, while in Ref. 13 no exotic scenario is invoked to interpret this state. QCD sum rules are compatible with the $c\bar{s}$ interpretation.

To understand the structure of a particle one needs to analyse its decay modes under definite assumptions and compare the result with the experimental measurements. In the following we present an analysis based on such a strategy to discuss whether the identification of $D_{sJ}^*(2317)$ and $D_{sJ}(2460)$ is supported by data. To this end, we compute the decay modes of a scalar and an axial-vector particle with masses of 2317 MeV and 2460 MeV respectively, and check whether they can be predicted in agreement with the experimental findings presently available. In particular, the isospin violating decays to $D_{sJ}^0 \pi^0$ should proceed at a rate larger than the radiative modes, though not exceeding the experimental upper bounds on the total widths.

2. Hadronic Modes
In order to analyze the isospin violating transitions $D_{s0} \rightarrow D_s \pi^0$ and $D_{s1}' \rightarrow D_s^* \pi^0$, 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 $\vec{s}_\ell$ are separately conserved. This allows to classify hadrons with a single heavy quark $Q$ in terms of $s_\ell$ by collecting them in doublets the members of which only differ for the relative orientation of $\vec{s}_Q$ and $\vec{s}_\ell$. The doublets with $J^P = (0^-, 1^-)$ and $J^P = (0^+, 1^+)$ (corresponding to $s_\ell^P = \frac{1}{2}^-$ and $s_\ell^P = \frac{1}{2}^+$, respectively) can be described by the effective fields

$$H_a = \frac{1 + \gamma^5}{2} [P^{a\mu}_0 \gamma^\mu - P_a \gamma^5]$$

$$S_a = \frac{1 + \gamma^5}{2} [P^{a\mu}_1 \gamma^\mu \gamma^5 - P_{0a}^*]$$

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 $P_a^{(*)} = D^{(*)0}, D^{(*)+}$ and $D_s^{(*)}$ (for $a = 1, 2, 3$); analogously, the components of $S_a$ are $P_{0a} = D_0^0, D_0^+, D_s^*$ and $P_{1a} = D_1^0, D_1^+, D_{s1}'$.

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\(^{15,16,17,18}\):

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

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

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

\[
L_{\text{mass}} = \frac{\mu f^2}{4} Tr\{\xi m_q \xi + \xi^\dagger m_q \xi^\dagger\}
\]

with \( m_q \) the light quark mass matrix:

\[
m_q = \begin{pmatrix}
m_u & 0 & 0 \\
0 & m_d & 0 \\
0 & 0 & m_s
\end{pmatrix} .
\]
In addition to the light meson mass terms, $L_{mass}$ contains an interaction term between $\pi^0$ ($I = 1$) and $\eta$ ($I = 0$) mesons: 

$$L_{mixing} = \frac{k}{2} \frac{m_s - m_u}{\sqrt{2} \sqrt{m_s + m_u}} \eta^0 \eta^0.$$

which vanishes in the limit $m_u = m_d$. Let us focus on the mode $D_s^* \rightarrow D_s \pi^0$. As in the case of $D_s^* \rightarrow D_s \pi^0$ studied in Ref. 19, the isospin mixing term can drive such a transition. The amplitude $A(D_s^* \rightarrow D_s \pi^0)$ is simply written in terms of $A(D_s^0 \rightarrow D_s \eta)$ obtained from (3), $A(\eta \rightarrow \pi^0)$ from (6) 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 $(m_d - m_u)/(m_s - m_u)$ accounting for isospin violation, so that the width $\Gamma(D_s^* \rightarrow D_s \pi^0)$ reads:

$$\Gamma(D_s^* \rightarrow D_s \pi^0) = \frac{1}{16\pi} \frac{h^2}{f^2} \frac{M_{D_s}}{M_{D_s^0}} \left( \frac{m_d - m_u}{m_s - m_u} \right)^2 \left( 1 + \frac{m_{\pi^0}^2}{|p_{\pi^0}|^2} \right) |\mathbf{q}_{\pi^0}| \ . \ (7)$$

As for $h$, the result of QCD sum rule analyses of various heavy-light quark current correlators is $|h| = 0.6 \pm 0.2 \, 17$. Using the central value, together with the factor $(m_d - m_u)/(m_s - m_u) \simeq \frac{1}{4}$ and $f = f_{\pi} = 132$ MeV we obtain$^{21}$:

$$\Gamma(D_s^* \rightarrow D_s \pi^0) = 7 \pm 1 \, K e V . \ (8)$$

The analogous calculation for $D_s^0 \rightarrow D_s \pi^0$ provides the result$^8$:

$$\Gamma(D_s^0 \rightarrow D_s \pi^0) = 7 \pm 1 \, K e V . \ (9)$$

3. Radiative Modes

Let us now turn to the calculation of radiative decay rates. We describe the procedure considering the mode $D_s^0 \rightarrow D_s^* \gamma$, the amplitude of which has the form:

$$A(D_s^0 \rightarrow D_s^* \gamma) = e \mathbf{d} \cdot \left( (\mathbf{e} \cdot \eta^*(p \cdot k) - (\eta^* \cdot p)(\mathbf{e} \cdot \mathbf{k}) \right) \ , \ (10)$$

where $p$ is the $D_s^0$ momentum, $\mathbf{e}$ the $D_s^*$ polarization vector, and $k$ and $\eta$ the photon momentum and polarization. The corresponding decay rate is:

$$\Gamma(D_s^0 \rightarrow D_s^* \gamma) = \alpha |\mathbf{d}|^2 |\mathbf{k}|^3 \ . \ (11)$$

The parameter $d$ gets contributions from the photon couplings to the light quark part $e_s s \gamma_\mu s$ and to the heavy quark part $e_c \bar{c} \gamma_\mu c$ of the electromagnetic

$^a$Electromagnetic contributions to $D_s^0 \rightarrow D_s \pi^0$ are expected to be suppressed with respect to the strong interaction mechanism considered here.
current, $e_s$ and $e_c$ being strange and charm quark charges in units of $e$. Its general structure is:

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

(12)

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, suppression that has been experimentally confirmed. From (11,12) 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. We first consider the coupling of the photon to the heavy quark part of the e.m. current. The matrix element $\langle D_{s0}^*(v',\epsilon)|\bar{c}\gamma_{\mu}c|D_{s0}^*(v)\rangle(v,v'\text{ meson four-velocities})$ can be computed in the heavy quark limit, matching the QCD $\bar{c}\gamma_{\mu}c$ current onto the corresponding HQET expression:

$$J_{\mu}^{HQET} = \tilde{h}_\nu[v_\mu + \frac{i}{2m_Q}(\tilde{\sigma}_{\mu} - \frac{v_\mu}{v^\nu}(\tilde{\sigma}_\nu + \frac{v_\nu}{v^\mu}(\tilde{\sigma}_\mu)) + \cdots]h_v$$

(13)

where $h_v$ is the effective field of the heavy quark. For transitions involving $D_{s0}^*$ and $D_{s0}^*$, and for $v = v'$ ($v \cdot v' = 1$), the matrix element of $J_{\mu}^{HQET}$ 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 $v \cdot v' = (m_{D_s}^2 + m_{D_s}^2)/(2m_{D_{s0}}^2 m_{D_s}^2) = 1.004$. Therefore, we neglect $d^{(h)}$ in (12).

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 $\phi(1020)$:

$$\langle D_{s0}^*(v',\epsilon)|\bar{s}\gamma_{\mu}s|D_{s0}^*(v)\rangle =$$

$$\sum_\lambda \langle D_{s0}^*(v',\epsilon)\phi(k,\epsilon_1(\lambda))|D_{s0}^*(v)\rangle \frac{i}{k^2 - M_\phi^2} \langle 0|\bar{s}\gamma_{\mu}s|\phi(k,\epsilon_1(\lambda))\rangle$$

(14)

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 $f_\phi = 234$ MeV. The matrix element $\langle D_{s0}^*(v',\epsilon)\phi(k,\epsilon_1)|D_{s0}^*(v)\rangle$ describes the strong interaction of a light vector meson ($\phi$) with two heavy mesons ($D_{s0}^*$ 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.
The Lagrange density is set up using the hidden gauge symmetry method, with the light vector mesons collected in a $3 \times 3$ matrix $\hat{\rho}_\mu$ analogous to $M$ in (4). The lagrangian reads as:

$$L' = i \hat{\mu} Tr \{ \bar{S}_a H_b \sigma^{\lambda\nu} V_{\lambda\nu} (\rho)_{ba} \} + h.c.,$$

with $V_{\lambda\nu} (\rho) = \partial_\lambda \rho_\nu - \partial_\nu \rho_\lambda + [\rho_\lambda, \rho_\nu]$ and $\rho_\lambda = i \frac{g_V}{\sqrt{2}} \hat{\rho}_\lambda$, $g_V$ being fixed to $g_V = 5.8$ by the KSRF rule. The coupling $\hat{\mu}$ in (15) is constrained to $\hat{\mu} = -0.1 \text{GeV}^{-1}$ by the analysis of the $D \to K^*$ semileptonic transitions induced by the axial weak current. The resulting expression for $\frac{1}{\Lambda_s}$ is:

$$\frac{1}{\Lambda_s} = -4 \hat{\mu} \frac{g_V}{\sqrt{2}} \frac{M_{D^*_s}}{M_{D^*_s} M_{\phi}}.$$  

The numerical result for the radiative width shows that the hadronic $D^*_s \to D_\gamma$ transition is more probable than the radiative mode $D^*_s \to D_\gamma \pi^0$, at odds with the case of the $D^*_s$ meson, where the radiative mode dominates the decay rate. In particular, if we assume that the two modes essentially saturate the $D^*_s$ width, we have $\Gamma(D^*_s) = 8 \pm 1$ KeV. As for the two radiative modes allowed for $D_{sJ}(2460)$, one finds:

$$\Gamma(D'_{s1} \to D_{s\gamma}) = 3.3 \pm 0.6 \text{ KeV} \quad \Gamma(D'_{s1} \to D^*_s \gamma) = 1.5 \text{ KeV}$$

which in turn give a total width $\Gamma(D'_{s1}) = 12 \pm 1$ KeV.

4. Comparison with other approaches

The results of the previous two sections show that, within the described approach, the observed hierarchy of hadronic versus radiative modes is realized, supporting the identification of $D^*_s(2317)$ and $D_{sJ}(2460)$ with $(D_{s0}^*, D'_{s1})$. Other analyses have followed the same strategy of computing decay rates of the two narrow states in order to understand their structure. In Table 1 we compare our results with the outcome of other approaches based on the $c\bar{s}$ picture as well. Analyses in which the states are assumed to have an exotic structure provide larger values for the widths ($O(10^2)$ KeV).

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b The standard $\omega_8 - \omega_0$ mixing is assumed, resulting in a pure $\bar{s}s$ structure for $\phi$.

c The role of other possible structures in the effective lagrangian contributing to radiative decays is discussed in Ref. 25.
Table 1. Estimated width (KeV) of $D^*_{s0}$ and $D^*_{s1}$, using the $c\bar{s}$ picture. The results in column [35] are obtained using experimental inputs from Belle (Focus).

| Decay mode | [31] | [32] | [21, 8] | [33] | [34] |
|------------|------|------|---------|------|------|
| $D^*_{s0} \rightarrow D_s \pi^0$ | 21.5 | $\simeq 10$ | 7 ± 1 | 16 | 129 ± 43 |
| $D^*_{s0} \rightarrow D_s^* \gamma$ | 1.74 | 1.9 | 0.85 ± 0.05 | 0.2 | $\leq 1.4$ |
| $D^*_{s1} \rightarrow D_s^* \pi^0$ | 21.5 | $\simeq 10$ | 7 ± 1 | 32 | 187 ± 73 |
| $D^*_{s1} \rightarrow D_s^* \gamma$ | 5.08 | 6.2 | 3.3 ± 0.6 | | $\leq 5$ |
| $D^*_{s1} \rightarrow D_s^* \gamma$ | 4.66 | 5.5 | 1.5 | | |

In particular, we observe that conclusions analogous to those presented above have been reached in Ref. 31, which is based on the observation that heavy-light systems should appear as parity-doubled, i.e. in pairs differing for parity and transforming according to a linear representation of chiral symmetry. In particular, the doublet composed by the states having $J^P = (0^+, 1^+)$ can be considered as the chiral partner of that with $J^P = (0^-, 1^-)$\(^d\). 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.

5. Conclusions and perspectives

We presented the calculation of hadronic and radiative decay rates of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ in a framework based on heavy quark symmetries and on the Vector Meson Dominance ansatz. This analysis shows that the observed narrow widths and the enhancement of the $D^*_{sJ}(2317)$ decay modes are compatible with the identification of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ with the states belonging to the $J^P = (0^+, 1^+)\frac{1}{2}$ doublet of the $c\bar{s}$ spectrum. Nevertheless, unanswered questions remain, such as the near equality of the masses of $D^*_{sJ}(2317)$ and $D_{sJ}(2460)$ with their non-strange partners. The missing evidence of the radiative mode $D^*_{sJ}(2317) \rightarrow D^*_{sJ} \gamma$ is another puzzling aspect deserving further experimental investigations.

The quantum number assignment has a rather straightforward consequence concerning 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

\(^d\)This idea was first suggested in Ref. 35 in order to obtain a consistent implementation of chiral symmetry and reconsidered also in Ref. 36.
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

In conclusion, we expect that new experimental results will continue to enrich this scenario. Actually, another excited $c\bar{s}$ meson has been recently observed by Selex Collaboration$^{38}$, motivating further studies in this exciting period for hadron spectroscopy.

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