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Possible interpretations of $D_{sJ}^{+}(2632)$ if it really exists

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Abstract: We analyze various possible interpretations of the narrow state $D_{sJ}^{+}(2632)$ observed by SELEX Collaboration recently, which lies above threshold and has abnormal decay pattern. These interpretations include: (1) several versions of tetraquarks; (2) conventional $c\bar{s}$ meson such as the first radial excitation of $D_s(2112)$ with abnormally large SU(3) symmetry breaking; (3) conventional $c\bar{s}$ meson with abnormally large η_1 coupling; (4) heavy hybrid meson. We discuss the physical implications of each interpretation. For example, if the existence of $D_{sJ}^{+}(2632)$ is confirmed as the first radial excitation of $D_s(2112)$ by other experiments, it will be helpful to look for (1) its SU(3) flavor partners $D_J^{0,+}(2530)$; (2) its B-meson analogues $B_J^{0,+}(5840), B_{sJ}^{+}(5940)$; (3) S-wave two pion decay modes. We also discuss possible reasons of negative results in search of $D_{sJ}^{+}(2632)$ by other experiments.

Keywords: Heavy Quarks Physics, QCD, Phenomenological Models.

Contents

1. Difficulty with $c\bar{s}$ assignment and SU(3) flavor symmetry

The experimental discovery of the narrow low-lying charm-strange mesons $D_{sJ}(2317)$, $D_{sJ}(2457)$ is an important event in heavy meson spectroscopy [1]–[4]. Many theoretical interpretations were proposed [5]–[9]. However, the large electromagnetic branching ratio of $D_{sJ}(2457)$ observed by BELLE Collaboration favors the assignment of these two states as conventional $c\bar{s}$ states [7, 8, 9]. These two states belong to the $(0^+, 1^+)$ doublet with $j_l = \frac{1}{2}^+$ in the heavy quark effective field theory.

Recently SELEX Collaboration observed an exotic charm-strange meson $D_{sJ}(2632)$. Its decay width is very narrow, $\Gamma < 17\,\mathrm{MeV}$ at 90% C.L. The decay channels are $D_s\eta$ and D^0K^+ with the unusual relative branching ratio: $\Gamma(D^0K^+)/\Gamma(D_s\eta) = 0.16 \pm 0.06$ [10]. This observation has inspired a lot of theoretical papers [11]–[20].

In order to explain why this relative branching ratio is abnormal, let's first assume (1) SU(3) flavor symmetry and (2) that the quark content of $D^+_{sJ}(2632)$ is $c\bar{s}$. The charge conjugate state $D^-_{sJ}(2632)$ belongs to $\mathrm{SU}(3)_F$ triplet when the heavy quark transforms as a $\mathrm{SU}(3)_F$ singlet. We denote the triplet as

$$(T^i) = \begin{pmatrix} \bar{c}u \\ \bar{c}d \\ \bar{c}s \end{pmatrix} = \begin{pmatrix} \bar{D}^0_J \\ \bar{D}^-_J \\ \bar{D}_{sJ}(2632) \end{pmatrix} \tag{1.1}$$

where $\bar{D}^0_J, \bar{D}^-_J$ are SU(3) flavor partners of $\bar{D}_{sJ}(2632)$.

Then the decay of heavy meson triplet T' can be described with the effective lagrangian in the $\mathrm{SU}(3)_F$ favor symmetry limit

$$L_8 = g_8 T'^{\dagger}_i M^i_j T^j + \mathrm{H.C.} \tag{1.2}$$

where M^i_j is the matrix of pseudoscalar meson octet

$$(M^i_j) = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & K^0 \\ K^- & \bar{K}^0 & -\frac{2\eta_8}{\sqrt{6}} \end{pmatrix}, \tag{1.3}$$

and T^j is the final state heavy pseudoscalar meson triplet. We suppressed the Lorentz indices for mesons.

The two-body decay width of a meson reads

$$\Gamma = g^2 \frac{k^{2L+1}}{m_0^{2L}}, \tag{1.4}$$

where g is the dimensionless effective coupling constant, L is the angular momentum for decay. m_0 is the parent mass and k is the decay momentum in the center of mass frame

$$k = \frac{1}{2m_0}\{[m_0^2 - (m_1+m_2)^2][m_0^2 - (m_1-m_2)^2]\}^{\frac{1}{2}}, \tag{1.5}$$

where m_1 and m_2 are the masses of final mesons. The ratio of decay widths of these two channels for $D^+_{sJ}(2632)$ is

$$\frac{\Gamma(D^0K^+)}{\Gamma(D^+_s\eta)} = \left(\frac{g_{D^0K^+}}{g_{D^+_s\eta}}\right)^2 \left(\frac{k_{D^0K^+}}{k_{D^+_s\eta}}\right)^{2L+1}. \tag{1.6}$$

Using $\lambda_{D^0K^+}/\lambda_{D_s\eta} = \sqrt{3/2}$, $k_{D^0K^+} = 499\,\mathrm{MeV}$, $k_{D_s\eta} = 325\,\mathrm{MeV}$, we get

$$\frac{\Gamma(D^0K^+)}{\Gamma(D_s\eta)} = 2.3 * (1.54)^{2L} \geq 2.3\,. \tag{1.7}$$

which is nearly 15 times larger than the experimental value [10]!

Clearly either some important physics is missing or one of the above two assumptions of SU(3) symmetry and $D_{sJ}(2632)$ as a $c\bar{s}$ state is wrong. In the following sections we analyze four different interpretations of this charming state and the resulting phenomenology.

2. Tetraquarks

One intriguing possibility is that $D_s(2632)$ is a tetraquark [11, 12, 13, 14, 16, 19]. Then, there should be other tetraquark partners of $D_{sJ}(2632)$. Moreover, its narrow width may require that quarks inside $D_{sJ}(2632)$ form tightly bound clusters like diquarks [14]. We give a brief review of different versions of tetraquark interpretations.

Tetraquarks with quark content $\bar{c}\bar{q}qq$ form four multiplets: two triplets, one anti-sextet and one 15-plet

$$3 \otimes 3 \otimes \bar{3} \otimes 1 = 3_1 \oplus 3_2 \oplus \bar{6} \oplus 15\,. \tag{2.1}$$

The wave functions of these states can be found in ref. [16]. The identification of $D_{sJ}(2632)$ as the $J^P = 0^+$ isoscalar member of the **15** tetraquarks with the quark content $\frac{1}{2\sqrt{2}}(ds\bar{d} + sd\bar{d} + su\bar{u} + us\bar{u} - 2ss\bar{s})\bar{c}$ leads to the relative branching ratio [16]

$$\frac{\Gamma(D_{sJ}^+(2632) \to D^0 K^+)}{\Gamma(D_{sJ}^+(2632) \to D_s^+ \eta)} = 0.25\,. \tag{2.2}$$

This decay pattern arises from the SU(3) Clebsch-Gordan coefficients very naturally.

Another possibility is that $D_{sJ}^+(2632)$ is dominated by $c\bar{s}s\bar{s}$ with $J^P = 0^+$ [12, 19]. The $s\bar{s}$ fluctuates into $1/\sqrt{3}\eta_1 - 2/\sqrt{6}\eta_8$. Hence, its decay mode is mainly $D_s^+\eta$. The final states D^0K^+, D^+K^0 are produced through the annihilation of $s\bar{s}$ into $u\bar{u} + d\bar{d}$ which requires η_1 component. Thus this process is OZI suppressed. In this way the anomalous decay pattern is achieved.

In the diquark correlation configuration [14], $D_{sJ}^+(2632)$ is suggested to be a $(cs)_{3^*} - (\bar{s}\bar{s})_3$ state where the subscript numbers are color representations. With the assumption that $(cs)_{3^*} - (\bar{s}\bar{s})_3$ has small mixing with $(c\bar{s})_1 - (s\bar{s})_1$, one can give a nice interpretation for the narrow width. The mixing between $s\bar{s}$ and $u\bar{u} + d\bar{d}$ can lead to the unusual branching ratio.

All the above three tetraquark interpretations predict the same production rates for D^+K^0, D^0K^+ final states. A serious challenge is that SELEX Collaboration didn't find any signal in the D^+K^0 channel [10].

A very different tetraquark version $cd\bar{d}\bar{s}$ is proposed in refs. [11, 13]. Naively, one would expect D^0K^+ decay channel is suppressed while D^+K^0 and $D_s\eta$ modes are both important. Interestingly, ref. [11] invoked the isospin symmetry breaking to explain the relative branching ratio. It was proposed that the mass eigenstate $D_{sJ}^+(2632)$ is the mixture between the two flavor eigenstates $a_{c\bar{s}}^+$ and $f_{c\bar{s}}^+$, where $a_{c\bar{s}}^+$ is the $I = 1, I_3 = 0$ state in $\mathrm{SU(3)}_F$ 6 representation and $f_{c\bar{s}}^+$ is the $I = 0$ state in $\bar{3}_1$ with our notation. With some special mixing scheme, the relative ratio is found to be $\frac{\Gamma(D^0K^+)}{\Gamma(D_s^+\eta)} = 0.16$. At the same time, the authors of ref. [11] predicted $4 < \frac{\Gamma(D^+K^0)}{\Gamma(D_s^+\eta)} < 7.6$, $1.7 < \frac{\Gamma(D_s\pi^0)}{\Gamma(D_s^+\eta)} < 6.5$.

3. Large $\mathrm{SU(3)}_F$ breaking and $D_{sJ}(2632)$ as a $c\bar{s}$ state

Some authors suggested that $D_{sJ}(2632)$ be the first radial excitation of $D_s(2112)$ with $J^P = 1^-$. The nodal structure of the radial wave function of $D_s(2632)$ ensures the narrow width while different decay momenta in two channels lead to anomalous decay pattern [14, 15].

We want to emphasize that two new narrow states D_J^0 and D_J^+ should also exist as flavor partners of $D_{sJ}(2632)$ within this scheme. The mass difference between $D_{sJ}(2632)$ and D_J^0 should roughly be the strange quark mass $m_s = 100\,\mathrm{MeV}$. So their masses are around 2532 MeV. If $D_{sJ}(2632)$ is interpreted as the first radial excitation of $D_s(2112)$, $D_J^0(2532)$ should be the first radial excitation of $D^*(2010)$. Similarly one expects the following B-meson analogues: $B_J^{0,+}(5840), B_{sJ}^+(5940)$. Moreover, the branching ratio of the S-wave two pion decay modes may be significant if $D_{sJ}(2632)$ is the radial excitation.

We want to point out that the decay momenta of the two channels $D_s^+\eta$ and D^0K^+ are the same in the exact $\mathrm{SU(3)}_F$ symmetry. Then there is no anomalous decay pattern.

In other words, the origin of this kind of explanation of the anomalous decay pattern can be traced back to the SU(3) flavor symmetry breaking effects. In the following we shall use the effective lagrangian formalism to analyze this effect in a model-independent way.

The symmetry breaking is caused by the quark mass matrix $m = diag(\hat{m}, \hat{m}, m_s)$. We assume the isospin symmetry. The symmetry breaking lagrangian reads

$$L_m = \alpha T_i'^{\dagger} m_j^i (S_3)^j + \beta T_i'^{\dagger} m_j^l (S_{\bar{6}})_{lk} \epsilon^{ijk} + \gamma T_i'^{\dagger} m_j^k (S_{15})_k^{ij} \tag{3.1}$$

where

$$\begin{aligned} (S_3)^i &= M_j^i T^j \,, \\ (S_{\bar{6}})_{ij} &= M_j^a T^b \epsilon_{iab} + M_i^a T^b \epsilon_{jab} \end{aligned}$$

and

$$(S_{15})_k^{ij} = M_k^i T^j + M_k^j T^i - \frac{1}{4}(\delta_k^i M_a^j T^a + \delta_k^j M_a^i T^a) \,.$$

For $D_{sJ}^-(2632)$, we have

$$\begin{aligned} L_{D_{sJ}} &= \left\{(g_8 + \alpha m_s) - \frac{1}{2}\gamma(m_s - \hat{m})\right\}(D_{sJ}^+ K^- \bar{D}^0 + D_{sJ}^+ \bar{K}^0 D^-) - \\ &\quad - \frac{2}{\sqrt{6}}\left\{(g_8 + \alpha m_s) + \frac{3}{2}\gamma(m_s - \hat{m})\right\} D_{sJ}^+ \eta_8 D_s^- \,. \end{aligned} \tag{3.2}$$

$$\frac{g_{\bar{D}^0 K^-}}{g_{D_s^- \eta_8}} = \frac{1}{\sqrt{6}}\left[1 - \frac{4(g_8 + \alpha m_s)}{(g_8 + \alpha m_s) + \frac{3}{2}\gamma(m_s - \hat{m})}\right] . \tag{3.3}$$

Naively one would expect $|\frac{\gamma(m_s - \hat{m})}{g_8 + \alpha m_s}| \ll 1$. If we assume the physical η meson is approximately η_8, we can extract the value $\frac{\gamma(m_s - \hat{m})}{g_8 + \alpha m_s}$ from the relative branching ratio $\frac{\Gamma(D^0 K^+)}{\Gamma(D_s^+ \eta)}$.

Numerically there are two possibilities:

$$\frac{\gamma(m_s - \hat{m})}{g_8 + \alpha m_s} = 4.8 \,. \tag{3.4}$$

Or

$$\frac{\gamma(m_s - \hat{m})}{g_8 + \alpha m_s} = 1.1 \,. \tag{3.5}$$

Although the above numbers seem quite unnatural, the anomalously large SU$(3)_F$ breaking can explain the special decay pattern of $D_{sJ}(2632)$ in principle.

For $\bar{D}_J^0$ and D_J^-,

$$\begin{aligned} L_{D_J} &= \left\{(g_8 + \alpha\hat{m}) - (\beta + \frac{1}{4}\gamma)(m_s - \hat{m})\right\} \times \\ &\quad \times \left\{\left(\frac{1}{\sqrt{2}} D_J^0 \pi^0 \bar{D}^0 + D_J^0 \pi^+ D^-\right) + (D_J^+ \pi^- \bar{D}^0 - \frac{1}{\sqrt{2}} D_J^+ \pi^0 D^-)\right\} + \\ &\quad + \left\{(g_8 + \alpha\hat{m}) + (\beta + \frac{3}{4}\gamma)(m_s - \hat{m})\right\}(D_J^0 K^+ D_s^- + D_J^+ K^0 D_s^-) + \\ &\quad + \frac{1}{\sqrt{6}}\left\{(g_8 + \alpha\hat{m}) + 3\left(\beta - \frac{3}{4}\gamma\right)(m_s - \hat{m})\right\}(D_J^0 \eta_8 \bar{D}^0 + D_J^+ \eta_8 D^-) \,. \end{aligned} \tag{3.6}$$

In general, $D_J^{0,+}$ may not have the same decay pattern as $D_{sJ}(2632)$. For example, in the extreme case $|\alpha| \ll |\gamma|, |\beta| \ll |\gamma|$, we have the following relative branching ratios

$$\Gamma(D_J^+ \to D^0\Pi^+) : \Gamma(D_J^+ \to D^+\Pi^0) : \Gamma(D_J^+ \to D_s^+\bar{K}^0) : \Gamma(D_J^+ \to D^+\eta) = \\ = 1 : 0.5 : 36.6 : 74.8 \tag{3.7}$$

for the ratio in eq. (3.4) and

$$\Gamma(D_J^+ \to D^0\Pi^+) : \Gamma(D_J^+ \to D^+\Pi^0) : \Gamma(D_J^+ \to D_s^+\bar{K}^0) : \Gamma(D_J^+ \to D^+\eta) = \\ = 1 : 0.5 : 0.4 : 0.1 \tag{3.8}$$

for the ratio in eq. (3.5).

4. Large singlet coupling and $D_{sJ}(2632)$ as a $c\bar{s}$ state

As in the last section we assume $D_{sJ}(2632)$ is a $c\bar{s}$ state with $J^P = 1^-$. Instead of invoking very large SU(3) breaking effects to account for its decay pattern, we introduce the interaction between SU(3) singlet η_1 and $D_{sJ}(2632)$ and consider the mixing between η_8 and η_1. Now the effective lagrangian becomes

$$L_{\text{eff}} = g_8 T_i^{\prime\dagger} M_j^i T^j + g_1 \eta_1 T_i^{\prime\dagger} T^i \,. \tag{4.1}$$

The physical states η and η' are admixtures of η_8 and η_1,

$$\begin{aligned} \eta &= \eta_8 \cos\theta - \eta_1 \sin\theta \\ \eta' &= \eta_8 \sin\theta + \eta_1 \cos\theta \end{aligned} \tag{4.2}$$

with the mixing angle $\theta = -20^\circ$ [21]. Now we have

$$\begin{aligned} L_{D_{sJ}} = g_8 \Big\{ & D_{sJ}^+ K^- \bar{D}^0 + D_{sJ}^+ \bar{K}^0 D^- - \left(\frac{2}{\sqrt{6}}\cos\theta + \frac{g_1}{g_8}\sin\theta\right) D_{sJ}^+ \eta D_s^- + \\ & + \left(-\frac{2}{\sqrt{6}}\sin\theta + \frac{g_1}{g_8}\cos\theta\right) D_{sJ}^+ \eta' D_s^- \Big\} . \end{aligned} \tag{4.3}$$

The lagrangian involving $\bar{D}_J^0$ and D_J^- reads

$$\begin{aligned} L_{D_J} = g_8 \Big\{ & \frac{1}{\sqrt{2}} D_J^0 \pi^0 \bar{D}^0 + D_J^0 \pi^+ D^- + D_J^0 K^+ D_s^- - \\ & - \frac{1}{\sqrt{2}} D_J^+ \pi^0 D^- + D_J^+ \pi^- \bar{D}^0 + D_J^+ K^0 D_s^- + \\ & + \left(\frac{1}{\sqrt{6}}\cos\theta - \frac{g_1}{g_8}\sin\theta\right)(D_J^0 \eta \bar{D}^0 + D_J^+ \eta D^-) + \\ & + \left(\frac{1}{\sqrt{6}}\sin\theta + \frac{g_1}{g_8}\cos\theta\right)(D_J^0 \eta' \bar{D}^0 + D_J^+ \eta' D^-) \Big\} . \end{aligned} \tag{4.4}$$

From the relative decay ratio we can extract either

$$\frac{g_1}{g_8} = 16.2 \tag{4.5}$$

or

$$\frac{g_1}{g_8} = -11.7\,. \tag{4.6}$$

With these values, we have following the relative branching ratios

$$\Gamma(D_J^+ \to D^0\Pi^+) : \Gamma(D_J^+ \to D^+\Pi^0) : \Gamma(D_J^+ \to D_s^+\bar{K}^0) : \Gamma(D_J^+ \to D^+\eta) = \\ = 1 : 0.5 : 0.07 : 6.6 \tag{4.7}$$

for the ratio in eq. (4.5) and

$$\Gamma(D_J^+ \to D^0\Pi^+) : \Gamma(D_J^+ \to D^+\Pi^0) : \Gamma(D_J^+ \to D_s^+\bar{K}^0) : \Gamma(D_J^+ \to D^+\eta) = \\ 1 : 0.5 : 0.07 : 2.4 \tag{4.8}$$

for the ratio in eq. (4.6).

It is interesting to note that η_1 mixes strongly with $G\tilde{G}$ because of axial anomaly. The anomalously large coupling between η_1 and $D_{sJ}^+(2632)$ indicates that $D_{sJ}^+(2632)$ may be a heavy hybrid meson containing explicit glue [22]. That's the subject of the next section.

5. Large N_c limit and $D_{sJ}^+(2632)$ as a heavy hybrid meson

Large N_c formalism is a very useful and unique tool which is applicable in the whole energy regime from zero to infinity. In this section, we explore the possibility of $D_s(2632)$ being a hybrid state containing explicit glue using large N_c expansion. Recall that the normalized interpolating currents for a pure glueball and a conventional meson A in the large N_c limit read

$$O_A = \frac{1}{\sqrt{N_c}}\bar{q}^i\Gamma_A q^i \tag{5.1}$$

$$O_G = g_s^2 G^2\,. \tag{5.2}$$

The factor $1/\sqrt{N_c}$ is introduced to ensure the creation amplitude by O_A from the vacuum is $\sim \mathcal{O}(N_c^0)$ in the large N_c limit. I.e., the following matrix elements are order unity when $N_c \to \infty$:

$$\langle|0|O_A(0)|A\rangle = F_A \sim \mathcal{O}(N_c^0) \tag{5.3}$$

$$\langle|0|O_G(0)|\mathrm{Glueball}\rangle = F_G \sim \mathcal{O}(N_c^0)\,. \tag{5.4}$$

We have depicted $q\bar{q}$ mesons, $\bar{q}Gq$ hybrid states and glueballs in figure 1.

In order to illustrate the formalism, we consider the N_c order of the mixing amplitude between a glueball and a $q\bar{q}$ meson, which can be extracted through the correlation function:

$$\langle 0|\hat{T}O_A(x)O_G(y)|0\rangle\,. \tag{5.5}$$

The Feynman diagram is presented in figure 2. Also shown is the double-line representation in the large N_c approach. There are two independent closed color loops which contribute N_c^2 while two vertices contribute g_s^2. Its N_c order reads

$$\langle 0|\hat{T}O_A(x)O_G(y)|0\rangle \sim \frac{1}{\sqrt{N_c}}g_s^4 N_c^2 = N_c^{-\frac{1}{2}}\,. \tag{5.6}$$

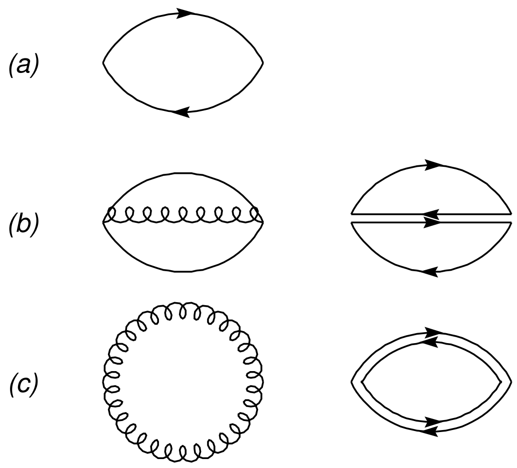

Figure 1: Feynman diagrams and the corresponding double-line representation for the decay modes of $J^P = 1^-$ heavy hybrid meson.

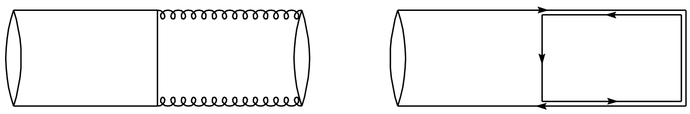

Figure 2: Feynman diagrams and the corresponding double-line representation for the decay modes of $J^P = 1^-$ heavy hybrid meson.

In other words, glueballs decouple from the conventional mesons when $N_c \to \infty$ [23].

We consider two kinds of heavy hybrid mesons: (1)

$$O_{H_1} = \frac{1}{\sqrt{N_c}} \bar{s}\gamma^\nu g_s G_{\mu\nu} c \tag{5.7}$$

with $J^P = 1^-$ and $c\bar{s}$ in the color octet state or (2)

$$O_{H_2} = \frac{1}{\sqrt{N_c}} \bar{s} i\gamma_5 c g_s^2 G_{\mu\nu} \tilde{G}^{\mu\nu} \tag{5.8}$$

with $J^P = 0^+$ and $c\bar{s}$ in the color singlet state.

$$\langle 0|O_{H_{1\mu}}(0)|\text{Vector Hybrid State} : H_1\rangle = F_{H_1}\epsilon_\mu \sim \mathcal{O}(N_c^0) \tag{5.9}$$

$$\langle 0|O_{H_2}(0)|\text{Scalar Hybrid State} : H_2\rangle = F_{H_2} \sim \mathcal{O}(N_c^0)\,. \tag{5.10}$$

For the $J^P = 1^-$ hybrid meson, its decay channels are shown in figure 3 using the Feynman diagram and double-line representation. The dominant decay mode in the large N_c limit is for the gluon to split into a color-octet $q\bar{q}$ pair in diagram (a) in figure 3. Then the $q\bar{q}$ recombine with $c\bar{s}$ pair to form $D^0K^+, D^+K^0, D_s\eta$. The N_c order of this decay amplitude is $\sim \left(\frac{1}{\sqrt{N_c}}\right)^3 g_s^2 N_c^2 = N_c^{-\frac{1}{2}}$. When it decays into a glueball and a $q\bar{q}$ meson, the decay amplitude is $\mathcal{O}(N_c^{-1})$, i.e., diagram (c). When it decays into two $q\bar{q}$ mesons through color-singlet two gluon intermediate states in diagram (b), the decay amplitude is $\mathcal{O}(N_c^{-\frac{3}{2}})$.

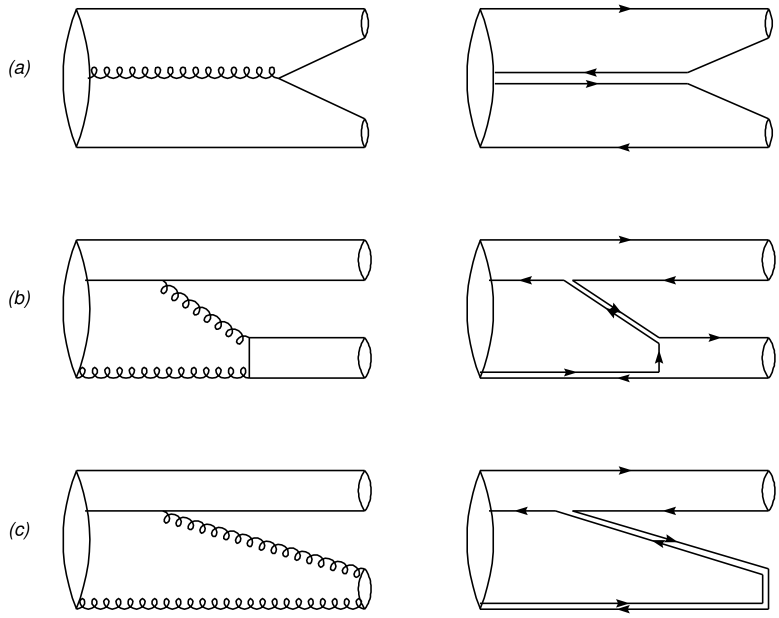

Figure 3: Feynman diagrams and the corresponding double-line representation for the decay modes of $J^P = 1^-$ heavy hybrid meson.

In other words, the assignment of $D_{sJ}(2632)$ as a $J^P = 1^-$ hybrid meson will not lead to the abnormal decay pattern observed by SELEX Collaboration.

Now we turn to the second case. With this configuration, it is possible to understand the narrowness of $D_s(2632)$. When it decays into a glueball and D_s in diagram (b) in figure 4, the amplitude is $\mathcal{O}(N_c^0)$ in the large N_c limit. Hence, this would be the most possible decay channel if kinematics allows. However, it cannot occur because the glueball is too heavy.

The decay mode $D_s\eta_1$ in diagram (c) in figure 4 is $\mathcal{O}(N_c^{-\frac{1}{2}})$ and subleading. However, it is also forbidden by kinematics. Therefore $D_s(2632)$ is narrow. We note that η meson is a mixture of η_8 and η_1. The mixing amplitude is $\mathcal{O}(N_c^0)$ arising from SU(3) symmetry breaking effect. Hence the kinematically allowed decay mode $D_s\eta$ becomes the dominant mode although it is $\mathcal{O}(N_c^{-\frac{1}{2}})$.

DK channel can happen only when the color-singlet two gluons annihilate into a pair of light quarks, namely after the hybrid becomes a four-quark state. This hybrid to four-quark transition is $1/N_c$ suppressed in amplitude. The four-quark state can decay to both $D_s\eta$ and DK, and the later channel is further suppressed compared to the former due to the color mismatch factor which is $1/N_c$ in amplitude. The N_c order can be read from diagram (a) in figure 4

$$\sim \left(\frac{1}{\sqrt{N_c}}\right)^3 g_s^4 N_c^2 = N_c^{-\frac{3}{2}} . \tag{5.11}$$

Naively there are three closed color loops, hence one would get the factor N_c^3. However, the charm and strange quark are in the color singlet state. Hence their color indices are

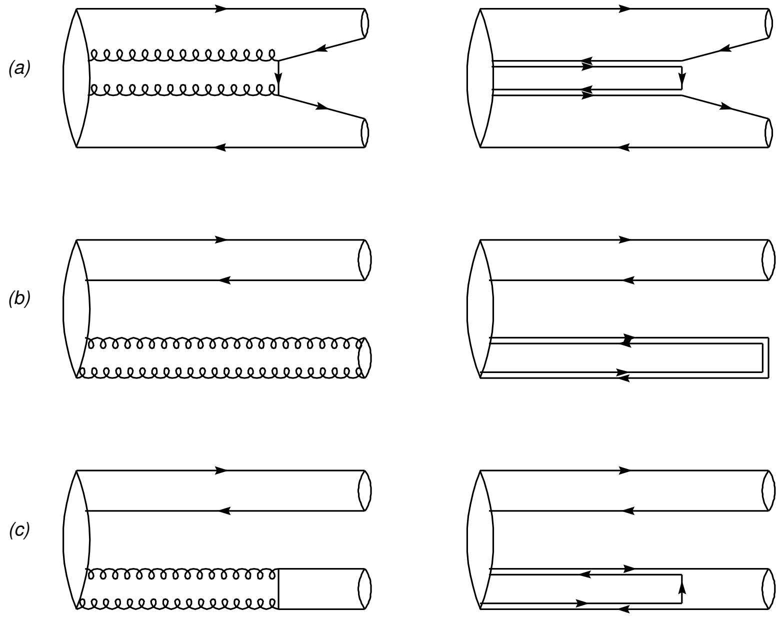

Figure 4: Feynman diagrams and the corresponding double-line representation for the decay modes of $J^P = 0^+$ heavy hybrid meson.

the same. In other words, there are only two independent color loops. Finally we obtain

$$\frac{\Gamma(D^0K^+)}{\Gamma(D_s\eta)} \sim \frac{\mathcal{O}(1/N_c^3)}{|\mathcal{O}(N_c^{-\frac{1}{2}}) + \mathcal{O}(N_c^{-\frac{3}{2}})|^2} = \mathcal{O}(N_c^{-2}) \sim \frac{1}{9} \times 1.5 \approx 0.16 \tag{5.12}$$

where the numerical value arises from using $N_c = 3$. The assignment of $D_{sJ}(2632)$ as a $c\bar{s}g_s^2 G\tilde{G}$ state explains the decay pattern quite naturally. A dynamical calculation of the mass of such a heavy hybrid meson will be very desirable, which is beyond the present note.

6. Discussions

In this note, we considered various possible interpretations of $D_{sJ}^+(2632)$ and discussed the resulting phenomenology. Future experimental confirmation of this state and discovery of its partners will help unveil the mysterious underlying dynamics of this narrow charm-strange meson.

Very recently, BABAR, CLEO and FOCUS reported negative results in the search of $D_{sJ}(2632)$ [24]. All these three experiments used different type of beams, and study different type of productions. BABAR and CLEO are e^+e^- collider experiments while FOCUS is the photoproduction experiment. In contrast, SELEX is a hadroproduction experiment. The major difference is that (1) SELEX utilized Σ^- beam which consists strange quark explicitly and (2) SELEX looks at very large X_f, which is not available for colliders [25].

At BABAR and CLEO, the center of mass is close to the Upsilon resonance while SELEX is a fixed target hadroproduction experiment with a beam momentum of 500 GeV. The production ratios of different charm states could be quite different between e^+e^- annihilation and gluon-gluon fusion which is the dominant production mechanism in hadroproduction [25]. If $D_{sJ}(2632)$ really exists, the reason that SELEX saw it while BABAR, CLEO and FOCUS didn't could be some special mechanism in hadroproduction at large X_f.

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