Decay modes of ideally mixed narrow pentaquark states

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We investigate the decay modes of pentaquark baryons for both the antidecuplet and the octet states, which were recently claimed by Close and Dudek to be potentially narrow due to a hidden selection rule. We discuss how ideal mixing between the pentaquark octet and antidecuplet states can be related to the OZI rule, which also naturally leads to certain selection rules in the pentaquark decays. We then introduce a tensor representation for the pentaquark states, and present the possible decay modes for both the unmixed and ideally mixed pentaquark octet and antidecuplet states. The exclusive decay modes can be used to experimentally search for the other pentaquark states.

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INTRODUCTION

The recent discovery of the Θ⁺ baryon by LEPS Collaboration at SPring-8 [1], which has been subsequently confirmed by several groups [2, 3, 4, 5, 6, 7, 8], initiated a lot of theoretical works in the field of hadron physics. Experimentally, Θ⁺ is observed to have a mass of 1540 MeV and a decay width of < 25 MeV, but there are indications that the width should in fact be much smaller [9, 10, 11, 12]. The Θ⁺(1540) together with the recently discovered Ξ⁺(1862) [13] (See also Ref. [14].) are expected to be the members of the antidecuplet pentaquark states [15]. At this stage, a pressing issue is to identify the other members of the antidecuplet and the pentaquark octet states. What is also potentially interesting about these pentaquark states is their small decay widths. Although Θ⁺ can decay into a kaon and a nucleon, its width is smaller than that obtained from any reasonable estimate based on the assumption that it is a meson-nucleon quasi-bound state [14]. Hence, the picture of pentaquark states by Jaffe and Wilczek [17] based on a strong diquark states seems favorable, as the regrouping of the color, flavor-spin and spatial wavefunctions of the diquarks and an antiquark into the nucleon and kaon is expected to be quite costly [12, 14, 20, 21, 22]. In this picture, one can make two further interesting assumptions. The first is the fall-apart mechanism in their decays [10] and the other is the ideal mixing [17]. As we will discuss, both assumptions are the consequences of a generalized OZI rule.

In the formalism of Jaffe and Wilczek [17], two quarks form a color triplet boson state, two of which with a relative $p$-wave form a color $3$ state. These two diquarks then combine with an antiquark to form a color singlet state. If Θ⁺ belongs to antidecuplet, then it is an isospin singlet state. In this case, the diquarks, each in flavor $3$, will form a $6$, which then combines with an antiquark to give a $10$. However, since $6 \otimes 3 = 10 \oplus 8$, we also expect a pentaquark octet naturally, which differs from the soliton model picture [15]. As has been discussed by Close and Dudek in Ref. [19], the pentaquark octet is also expected to share a common feature with the antidecuplet and have a small decay width. The decays of these pentaquark octet and antidecuplet states are expected to proceed through a “fall-apart” mechanism. This mechanism states that there is no annihilation or creation of quark pairs and the decay proceeds through a recombination of diquarks from the strongly correlated diquarks into the final decay products. This follows from a generalized OZI rule in pentaquark decays, which also leads to ideal mixing between the octet and antidecuplet pentaquark states. If we apply this rule to the decay for both the pentaquark octet and antidecuplet decays, we obtain the “selection rules” advocated by Close and Dudek in Ref. [19]. As we will show below, one can generalize the rule by using the tensor representation for the pentaquark states.

In this paper, we start with how the generalized OZI rule leads to ideal mixing of the pentaquark octet and antidecuplet states, and to the fall-apart mechanism for the pentaquark decays. We then introduce a tensor formalism for the pentaquarks and derive the selection rules for their decays. Using the tensor method, we present all the possible decay modes of the pentaquark octet and antidecuplet states into the baryon octet and meson octet. We also discuss why certain decays are not allowed. In addition, the pentaquark octet decays into baryon decuplet and meson octet are shown to be prohibited by the OZI rule even if they are allowed energetically and by SU(3) symmetry. The SU(3) symmetry breaking effects of this decay is also discussed.

IDEAL MIXING

Quantum mechanical example

The selection rule follows from the fall-apart mechanism, which effectively is the OZI rule for the pentaquark decays. Here, we investigate the idea of the generalized
OZI rule in the present context and show that it also naturally leads to ideal mixing among the pentaquark states.

Let us start with a simple pedagogical example of two-level quantum mechanical system, with the Hamiltonian

$$H = \begin{pmatrix} E_8 & \Delta \\ \Delta & E_1 \end{pmatrix}.$$  \hspace{1cm} (1)

Here, we call the two states as

$$\psi_8 = \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad \psi_1 = \begin{pmatrix} 0 \\ 1 \end{pmatrix}.$$  \hspace{1cm} (2)

If there is a mixing in the mass matrix as in Eq. (1), the physical states are the eigenstates of the Hamiltonian. The two eigenvalues are

$$\lambda_{\pm} = \frac{1}{2} \left( E_8 + E_1 \pm \sqrt{(E_8 - E_1)^2 + 4\Delta^2} \right),$$  \hspace{1cm} (3)

and the corresponding eigenvectors are

$$\phi = \cos \theta \psi_8 + \sin \theta \psi_1,$$

$$\omega = -\sin \theta \psi_8 + \cos \theta \psi_1,$$  \hspace{1cm} (4)

where

$$\tan \theta = \frac{\lambda_+ - E_8}{\Delta},$$  \hspace{1cm} (5)

and equivalently,

$$E_8 - E_1 = \frac{1 - \tan^2 \theta}{\tan \theta} \Delta.$$  \hspace{1cm} (6)

Formally, one can also write $E_8 = \langle \psi_8 | H | \psi_8 \rangle$, $E_1 = \langle \psi_1 | H | \psi_1 \rangle$, $\Delta = \langle \psi_1 | H | \psi_8 \rangle$. An alternative way to obtain the mixing angle, Eq. (5) or Eq. (6), is substituting Eq. (1) into the following condition,

$$\langle \phi | H | \omega \rangle = 0.$$  \hspace{1cm} (7)

In the case of ideal mixing, the mixing angle thus obtained also isolates the strange and nonstrange quark-antiquark pair into separate wavefunctions. That is, in such an ideal case, one finds that $\phi$ is a pure $\bar{s}s$ state while $\omega$ is purely a nonstrange $\bar{q}q$ state. Stated in a diagrammatic language, ideal mixing follows when connected quark graphs dominate over disconnected quark graphs, which is also known as the OZI rule. This is so because the eigenvalues satisfy $\lambda_{+} \gg \langle \bar{s}s | H | \bar{q}q \rangle$, where $\lambda_{+} = \langle \bar{s}s | H | \bar{s}s \rangle$ and $\lambda_{-} = \langle \bar{q}q | H | \bar{q}q \rangle$.

**Vector and pseudoscalar channel**

First, one should note that a prerequisite for mixing is SU(3) symmetry breaking. That is, the Hamiltonian in Eq. (1) should have explicit symmetry breaking such that

$$\Delta = \langle \psi_1 | H | \psi_8 \rangle \neq 0.$$  \hspace{1cm} (8)

For the vector meson channel, the physical wavefunctions obtained from isolating the strange and nonstrange quark-antiquark parts are,

$$\phi = -\bar{s}s = \sqrt{\frac{2}{3}} \psi_8 - \sqrt{\frac{1}{3}} \psi_1,$$

$$\omega = \frac{1}{\sqrt{2}} (\bar{u}u + \bar{d}d) = \sqrt{\frac{2}{3}} \psi_1 + \sqrt{\frac{1}{3}} \psi_8,$$  \hspace{1cm} (9)

where

$$\psi_1 = \frac{1}{\sqrt{3}} (\bar{u}u + \bar{d}d + \bar{s}s),$$

$$\psi_8 = \frac{1}{\sqrt{6}} (\bar{u}u + \bar{d}d - 2\bar{s}s).$$  \hspace{1cm} (10)

Hence, the mixing angle coming from isolating the nonstrange and strange part in the wavefunction is $\tan \theta = -1/\sqrt{2}$, which is also close to the physical mixing angle. Hence, ideal mixing in the vector channel implies

$$\langle \bar{s}\gamma_{\mu}s | H | \bar{u}\gamma_{\mu}u + \bar{d}\gamma_{\mu}d \rangle = 0,$$  \hspace{1cm} (11)

where we have included $\gamma_{\mu}$ to represent that the states are vector mesons.

For the pseudoscalar channel, the situation is quite different and the mixing angle is very small, suggesting that the disconnected quark lines are not suppressed,

$$\langle \bar{s}\gamma_{5}s | H | \bar{u}\gamma_{5}u + \bar{d}\gamma_{5}d \rangle \neq 0.$$  \hspace{1cm} (12)

The reason for such a large difference between Eqs. (11) and (12) are relatively well understood in both the perturbative and nonperturbative regime. Perturbatively, three gluon intermediate state are needed to connect the vector currents in Eq. (11), while only two gluon states are needed in Eq. (12). In the nonperturbative regime, at least two classical instanton configuration is needed to connect the vector currents, while only a single instanton is needed in the pseudoscalar channel [25].

**Pentaquarks**

The realization of ideal mixing between pentaquark octet and antidecuplet depends on whether a condition similar to Eq. (7) is fulfilled by the pentaquark currents. We first note that the wavefunctions of the nucleon with positive charge in pentaquark antidecuplet and octet are obtained as

$$N_{\bar{W}}^{+} = \frac{1}{\sqrt{3}} \left( |ud][ud]\bar{d} + \sqrt{2} |ud][us]_{+}\bar{s} \right),$$

$$N_{s}^{+} = \frac{1}{\sqrt{3}} \left( -\sqrt{2} |ud][ud]\bar{d} + |ud][us]_{+}\bar{s} \right).$$  \hspace{1cm} (13)
Here, \([ud] = \frac{1}{\sqrt{2}}(ud - du)\) and \([A][B]_+ = \frac{1}{\sqrt{2}}([A][B] + [B][A])\). We also have suppressed color and Dirac indices. As before, the ideal mixing is achieved when the following rule is satisfied:

\[
\langle [ud][ud]\bar{d}|H|[ud][us]\rangle = \text{small.} \tag{14}
\]

We call it “generalized OZI rule”, as none of the previous theoretical arguments, which justify the OZI rule in the meson sector, applies in this case. It is through experiments that we can confirm ideal mixing. But since ideal mixing is a consequence of the “generalized OZI rule” it is important to probe other consequences of this rule.

A useful result, that can be directly verified experimentally, is the characteristic features of decay modes of the pentaquark baryons due to the “generalized OZI rule”. Consider the decays of the pentaquark octet and antidecuplet into a kaon and a hyperon. The respective mixing is a consequence of the “generalized OZI rule” it is important to probe other consequences of this rule.

In this section, we introduce the tensor notation for the pentaquark states and derive selection rules for their decays assuming the OZI rule. The pentaquark antidecuplet and octet states can be represented by \(T^{ijk}\) and \(P^i\), respectively, where \(i, j, k\) are the flavor indices \(u, d, s\). We can obtain the SU(3) relation in their decays into either baryon decuplet \(D_{ijk}\) or normal baryon octet \(B^i\) made of three quarks and a meson octet \(M^i\) by constructing SU(3) invariant Lagrangians. For example, the SU(3) invariant relations for the decay of pentaquark antidecuplet into a normal baryon octet and meson octet can be obtained as \(23^1\)

\[
\mathcal{L}_{10} = -g_{10}^{\pi\eta}T^{ijk}B^iJ^k_m + (\text{H.c.}), \tag{17}
\]

where the particle identification of \(T^{ijk}\) can be found in Ref. \(23\) and the forms for \(B_i^j\) and \(M^i_j\) can be found, e.g., in Ref. \(24\). By expanding Eq. \(17\), we explicitly have

\[
\mathcal{L}_{10} = -\sqrt{6}g_{10}^{\pi\eta}\bar{K}_cN - \bar{N}_{10}^{\pi\eta}\bar{T}\cdot\pi N + \sqrt{3}\eta\bar{N}_{10}N - \sqrt{3}\bar{N}_{10}\bar{K}\Lambda + \bar{N}_{10}\bar{T}\cdot\Sigma K + i\bar{N}_{10}(\bar{\Sigma}\times\Sigma)\cdot\pi - \sqrt{3}\bar{N}_{10}\cdot\bar{T}\cdot\tau\Xi - \bar{K}_{c}\bar{N}_{10}\cdot\tau\Sigma K_{c} + (\text{H.c.}),
\]

where the overall factor \(g_{10}^{\pi\eta}\) is understood and

\[
N = \left( \begin{array}{c} p \\ n \end{array} \right), \quad \Xi = \left( \begin{array}{c} \Xi^0 \\ \Xi^- \end{array} \right), \quad K = \left( \begin{array}{c} K^+ \\ K^0 \end{array} \right),
\]

\[
K_{c} = \left( \begin{array}{c} \bar{K}^0 \\ -K^- \end{array} \right), \quad \tau \cdot \pi = \left( \begin{array}{c} \pi^0 \\ \sqrt{2}\pi^- \\ -\pi^0 \end{array} \right),
\]

etc. The isospin transition operator \(T\) is defined as \(25\)

\[
T^{(+1)} = \frac{1}{\sqrt{3}}\left( \begin{array}{c} \sqrt{6} \\ 0 \\ 0 \\ 0 \\ \sqrt{2} \\ 0 \\ 0 \\ 0 \\ 0 \end{array} \right), \quad T^0 = \frac{1}{\sqrt{3}}\left( \begin{array}{c} 0 \\ 0 \\ 0 \\ 0 \\ 2 \\ 0 \\ 0 \\ 0 \\ 2 \end{array} \right),
\]

\[
T^{(-1)} = \frac{1}{\sqrt{3}}\left( \begin{array}{c} 0 \\ 0 \\ 0 \\ 0 \\ 0 \\ \sqrt{2} \\ 0 \\ 0 \\ 0 \end{array} \right),
\]

so that \(T \cdot \pi = -T^{(+1)}\pi^+ + T^{(-1)}\pi^- + T^0\pi^0\) are on and so on.

The obtained couplings are listed in Table \(1\). Moreover, as shown in Ref. \(25\), it is straightforward to show that one can not form a coupling between the antidecuplet and baryon decuplet and meson octet, which suggests that the \(N(1710)\) can not be a member of the pure antidecuplet as it has large decay width into the \(\Delta\pi\) channel \(17, 14, 23\).

In order to construct SU(3) invariant interactions for the pentaquark octet, which respects the OZI rule, we have to be careful with the contraction of the SU(3) tensors. To do that, we note that the pentaquark octet and
TABLE I: Couplings of the pentaquark antidecuplet with the baryon octet and pseudoscalar meson octet. Multiplying the universal coupling constant \( g_{\pi \pi} \) is understood.

| \( \Theta \) | \( N_{10}^+ \) | \( N_0^{10} \) | \( \Sigma_{10}^+ \) |
|---|---|---|---|
| \( K^+ n \) | \( \sqrt{6} \) | \( \pi^+ n \) | \( \pi^+ \Lambda - \sqrt{3} \) |
| \( K^0 p \) | \( -\sqrt{6} \) | \( \pi^0 p \) | \( \pi^0 \Sigma^0 + \sqrt{3} \) |
| \( \eta \eta \) | \( \sqrt{6} \) | \( \eta \pi \) | \( \eta \Sigma^0 + \sqrt{3} \) |
| \( K^+ \Lambda \) | \( -\sqrt{3} \) | \( K^+ \Sigma^0 + \sqrt{3} \) | \( \eta \Sigma^0 + \sqrt{3} \) |
| \( K^0 \Sigma^0 \) | \( 1 \) | \( K^0 \Lambda - \sqrt{3} \) | \( K^0 \sqrt{2} \) |
| \( K^0 \Sigma^0 \) | \( -1 \) | \( K^0 \eta - \sqrt{3} \) | \( K^0 \sqrt{2} \) |

| \( \bar{\Sigma}_{10}^+ \) | \( \bar{\Sigma}_{10}^+ \) | \( \bar{\Sigma}_{10}^+ \) |
|---|---|---|
| \( \pi^\pm \Sigma^- \) | \( \pi^\pm \Sigma^- \) | \( \pi^\pm \Sigma^- \) |
| \( \pi^\pm \Lambda \) | \( \pi^\pm \Lambda \) | \( \pi^\pm \Lambda \) |
| \( \eta \Sigma^0 \) | \( \eta \Sigma^0 \) | \( \eta \Sigma^0 \) |
| \( K^+ \Xi^- \) | \( K^+ \Xi^- \) | \( K^+ \Xi^- \) |
| \( K^0 \Xi^- \) | \( K^0 \Xi^- \) | \( K^0 \Xi^- \) |
| \( K^0 \eta \) | \( K^0 \eta \) | \( K^0 \eta \) |
| \( K^- p \) | \( K^- p \) | \( K^- p \) |

When we construct an effective interaction of the pentaquark octet baryons with the baryon octet and pseudoscalar meson octet, one must to have the constraint \( f/d = 1/3 \), and the above selection rules hold regardless of the mixing with the antidecuplet members due to isospin, the possible interactions of the pentaquark octet baryons can be calculated as

\[
\mathcal{L}_8 = g_8 \epsilon^{i m \bar{c} n} \mathcal{S}_{[i,j,k]} \mathcal{B}_1^l M_m^k + (\text{H.c.})
\]

(24)

Substituting Eq. (22) into Eq. (23), one has

\[
\mathcal{L}_8 = 2 g_\pi \mathcal{P}_1^l M_l^m + g_\pi \mathcal{P}_1^l M_l^m + (\text{H.c.}),
\]

(25)

which is equivalent to Eq. (23) with \( f/d = 1/3 \). With this information, all the possible interactions of the pentaquark octet baryons can be calculated as

\[
\mathcal{L}_8 = g_{N_8 \pi} N_8 \tau \cdot \pi N + g_{N_8 \eta N_8} N_8 \n \cdot \n + i g_{N_8 \Sigma N_8} N_8 \tau \cdot \Sigma K + i g_{N_8 \Xi N_8} N_8 \tau \cdot \Xi K
\]

(26)

where

\[
\begin{align*}
g_{N_8 \pi} & = \frac{1}{\sqrt{2}}(d + f) g_8, \\
g_{N_8 \eta N_8} & = -\frac{1}{\sqrt{6}}(d - 3f) g_8, \\
g_{N_8 \Sigma N_8} & = -\sqrt{2f} g_8, \\
g_{N_8 \Xi N_8} & = \sqrt{2} \eta_8 \eta_8, \\
g_{N_8 \Lambda N_8} & = \frac{1}{\sqrt{2}}(d - f) g_8, \\
g_{N_8 \Lambda \n \cdot \n} & = \frac{1}{\sqrt{3}}(d + f) g_8, \\
g_{N_8 \Lambda \Sigma \cdot \Sigma} & = \frac{1}{\sqrt{2}} (d + f) g_8, \\
g_{N_8 \Lambda \Xi \cdot \Xi} & = \frac{1}{\sqrt{2}} (d + f) g_8, \\
g_{N_8 \n \cdot \n} & = \frac{1}{\sqrt{2}} (d + 3f) g_8, \\
g_{N_8 \n \cdot \n} & = -\frac{1}{\sqrt{2}} (d - f) g_8, \\
g_{N_8 \Sigma \cdot \Sigma} & = \frac{1}{\sqrt{2}} (d + f) g_8, \\
g_{N_8 \Xi \cdot \Xi} & = -\frac{1}{\sqrt{2}} (d - 3f) g_8.
\end{align*}
\]

(27)

With \( f/d = 1/3 \) and \( d = 3/2 \), the couplings are given explicitly in each channel in Table I. It can be easily seen that the constraint \( f/d = 1/3 \) gives rise to some selection rules so that the \( N_8 \n \cdot \n \), \( \Lambda_8 \Xi \) and \( \Xi \Lambda \) couplings vanish, as has been noted by Close and Dudek [19] from analyzing the decay of \( \Xi \). Since \( \Lambda_8 \) and \( \Xi_8 \) do not mix with the antidecuplet members due to isospin, the above selection rules hold regardless of the mixing with the antidecuplet.

Selection rule in ideal mixing

When there is ideal mixing between the pentaquark octet and antidecuplet states for the nucleon and Sigma...
resonances, the decay will separate into strange and nonstrange parts. Within ideal mixing, the nucleons with no $\bar{s}s$ component will decay into nonstrange hadrons and that with $\bar{s}s$ component will decay into strange hadrons only.\(^2\)

To separate the $\bar{s}s$ and light $\bar{q}q$ components in $\eta_8$, we introduce the interaction of the pentaquark octet with the normal baryon octet and meson singlet $\eta_1$ given by

$$L_1 = -\sqrt{3}g_1 T_i^i B_j^j \eta_1 + \text{(H.c.)}$$

$$= -\sqrt{3}g_1 \eta_1 \left( N_8 \bar{N}_N + \Lambda_8 \Lambda + \Xi_8 \Xi + \Sigma_8 \Sigma \right) + \text{(H.c.)}. \quad (28)$$

In ideal mixing, the couplings $g_8 = g_{10} = g_1 (= g)$ will be all equal \(^19\) and we have the interactions as

$$L_{\text{mixing}} = -\sqrt{3}g \sqrt{N_{\bar{q}q} \tau \cdot \pi N}$$

$$+ \sqrt{3}g \left( \eta_{\bar{q}q} N_{\bar{q}q} - \eta_{ss} N_{ss} \right) N$$

$$- \frac{3}{\sqrt{2}} g N_{ss} K \Lambda + \frac{3}{2} g N_{ss} \tau \cdot \Sigma K$$

$$+ i \left( \frac{3}{3} g \left( \Sigma_{\bar{q}q} \times \Sigma \right) \tau + \sqrt{\frac{3}{3} g \eta_{\bar{q}q} \Sigma_{\bar{q}q} \cdot \Sigma} - \sqrt{3} g \eta_{ss} \Sigma_{ss} \cdot \Sigma - \sqrt{3} g K \Sigma_{ss} \cdot \tau \Xi \right).$$

\(^2\) The nucleons with no $\bar{s}s$ component may couple to strange hadrons through quark-antiquark creation. However, this is possible only when the normal baryon has five quark component, which is suppressed in the OZI limit.
rule prohibits the coupling of the pentaquark octet with the baryon decuplet and meson octet unless other OZI-evading processes are allowed.

With our rule for tensor notation, this can be seen easily. That is, we can have a coupling between the pentaquark octet and baryon decuplet and meson octet in the form of

\[ \mathcal{L}_{int} = h_8 \epsilon^{ijk} \epsilon^{ipq} S_{i[p,q]} D_{jlm} M^m_k + (H.c.) \]

\[ = 3 h_8 \epsilon^{ijk} \eta^{ipq} D_{jlm} M^m_k + (H.c.). \quad (31) \]

However, as one can see, the antiquark index \( q \) in \( S_{i[p,q]} \) is not contracted with the index \( m \) in the meson field \( M \). Hence, unlike Eq. (24), the above equation does not respect the OZI rule. In this case, one can not make any SU(3) invariant terms where the index \( q \) in the pentaquark field is contracted with the index \( m \) in the meson field. This implies that \( h_8 = 0 \) in the OZI limit, and, therefore, rules out \( N(1440) \) and \( N(1710) \) as ideally mixed states of pentaquark octet and antidecuplet because the branching ratios of their decays into \( \Delta \pi \) are large. These decays are not possible even when the SU(3) symmetry breaking term is allowed for the interaction of pentaquark octet with baryon decuplet and meson octet. A leading order term that respects the generalized OZI rule can be written as

\[ \mathcal{L}_{int} = h'_8 \epsilon^{ijk} \epsilon^{ipq} S_{i[p,q]} D_{jlm} M^m_k Y^l_{i} + (H.c.), \quad (32) \]

where \( Y \) is the hypercharge operator. Since this symmetry breaking term still satisfies isospin symmetry, the interaction lagrangian is not affected in the SU(2) sector, i.e., \( N_8 \rightarrow \Delta \pi \). Therefore, the decay \( N_8 \rightarrow \Delta \pi \) is not allowed even in the SU(3) symmetry breaking interaction. \( \^2 \) One can explicitly work out the contractions to find that the decay is indeed forbidden. This shows that the \( N_8 \rightarrow \Delta \pi \) decay is an isospin-symmetry-breaking process, i.e., it is allowed only when we include the SU(2) symmetry breaking terms.

### CLOSING

We have discussed how the generalized OZI rule leads to ideal mixing and selection rules in the pentaquark decays. We have then introduced a simple tensor method that gives rise to the selection rules in the pentaquark baryon decays. Through experimental verification of the possible decay channels and their relative strength, one will be able to identify the other possibly narrow pentaquark states and test the picture of ideal mixing of pentaquark octet and antidecuplet states.

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