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Pentaquark symmetries, selection rules and another potentially narrow state

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Abstract

We identify essential differences between the pentaquark and chiral soliton models of $\overline{\mathbf{10}}_5$ and $\mathbf{8}_5$ pentaquarks and conventional $\mathbf{8}_3$ states, which are experimentally measurable. We show how the decays of \mathcal{E}_5 states in particular can test models of the pentaquarks, recommend study of the relative branching ratios of, e.g., $\mathcal{E}_5^- \rightarrow \mathcal{E}^- \pi^0 : \mathcal{E}^0 \pi^-$, and predict that the decay amplitude $\mathcal{E}_5 \rightarrow \mathcal{E}^* \pi$ is zero at leading order in pentaquark models for any mixture of $\overline{\mathbf{10}}$ and the associated $\mathbf{8}_5$. We also include a pedagogic discussion of wavefunctions in the pentaquark picture and show that pentaquark models have this $\mathbf{8}_5$ with $F/D = 1/3$, in leading order forbidding $\mathcal{E}_5 \rightarrow \Lambda K$. The role of Fermi–Dirac symmetry in the $qqqq$ wavefunction and its implications for the width of pentaquarks are briefly discussed. The relative couplings $g^2(\Theta_Q N K_Q^*)/g^2(\Theta_Q N K_Q) = 3$ for $Q \equiv s, c, b$. A further potentially narrow state Λ in $\mathbf{8}_5$ with $J^P = 3/2^+$ is predicted around 1650 MeV.

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1. Introduction

The possible discovery of an exotic and metastable baryon with positive strangeness, the $\Theta^+(1540)$, has led to an explosion of interest in chiral soliton models (a version of which is cited as having predicted this state) and their relation to quark models. In this Letter we propose explicit experimental tests that are sensitive to the assumed dynamics and thereby can distinguish among models.

Such a state was predicted in a version of the chiral soliton model [1,2] to be in a $\overline{\mathbf{10}}$ of flavour $SU(3)$ and to have $J^P = 1/2^+$. Subsequent to the

original observation [3], several interpretations of the state have been suggested within quark models (for an early review see [4]). Their common feature is that its constitution be $udud\bar{s}$ with one unit of orbital angular momentum in the wavefunction; they differ in the ways that the interquark dynamics causes the $1/2^+$ state to be the lightest and to have an anomalously narrow width.

The lightness and narrow width pose severe challenges for constituent quark model descriptions of the Θ . If it is indeed a positive parity state, this requires internal orbital angular momentum and associated excitation energy. Strong attractive correlations between (ud) have to be invoked, motivated by QCD or flavour-spin interactions, that are able to make this configuration lighter than the negative parity S -wave. Furthermore, there remains open the possibility that

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large $SU(3)_F$ violating effects may occur such that the concept of states forming a $\overline{\mathbf{10}}$ may be lost. Identifying the dynamics that could create such a resonance remains a central problem.

If the Θ^+ is an isoscalar, then a common feature of current models is that it is a member of a $\overline{\mathbf{10}}$, which contains further exotic states, $\mathcal{E}^{+,-}$. There are three main differences among the implications of models that can distinguish among various dynamics.

- (i) The magnitude of the mass gap spanning the $\overline{\mathbf{10}}$ from Θ to \mathcal{E} is significantly smaller in pentaquark models than in the original formulation of the chiral soliton model [2,4–6] though the latter is somewhat flexible as has recently been noted [7].
- (ii) The first excited state of Θ is predicted [8] in pentaquark models to be a $J^P = 3/2^+$ isoscalar, in a $J^P = 3/2^+$ $\overline{\mathbf{10}}$, whereas there is no place for such a state in the present formulation of chiral soliton models.
- (iii) The hadron decays of non-exotic members of the $\overline{\mathbf{10}}$, in particular those of $\mathcal{E}^{0,-}$ are especially sensitive to the interquark dynamics in pentaquark models. A specific example has been discussed in [9] but we shall show here that there is a more extensive set of relations and selection rules that arise in pentaquark models and which can discriminate among various dynamic and mixing schemes. In particular the relative strengths of decays $\mathcal{E}^- \rightarrow \mathcal{E}^- \pi^0 : \mathcal{E}^0 \pi^-$ test $\overline{\mathbf{10}}\text{--}\mathbf{8}$ mixing [9]; $\mathcal{E}^- \rightarrow \Lambda K^- : \Sigma^0 K^-$ have selection rules that test the decay dynamics that have been hypothesised [4,10,11] to suppress the pentaquark widths; and $\mathcal{E} \rightarrow \mathcal{E}^* \pi$ is predicted to vanish for both $\overline{\mathbf{10}}$ and $\mathbf{8}_5$ initial pentaquark states in such dynamics. The electromagnetic mass splittings of the \mathcal{E} states also contain important information.
- (iv) In pentaquark models the spin–orbit forces suggest the existence of a nearby $\mathbf{8}_5$ $J^P = 3/2^+$ multiplet containing a Λ that should be narrow and unmixed barring isospin violating effects.

We make some brief comments on point (ii) and then develop our main thesis, which focuses on points (iii) and (iv). We assume that $SU(3)_F$ violating effects are not so large as to hide all underlying flavour symmetry in the wavefunctions.

2. A low-lying $J^P = 3/2^+$ $\overline{\mathbf{10}}$ multiplet [8]

An essential difference between the pentaquark and chiral soliton (Skyrme) models appears to be in their implications for the first excited state of the Θ . In $qqqq\bar{q}$ with positive parity $1/2^+$ there is necessarily angular momentum present, which implies a family of siblings but with $J^P = 3/2^+$. The spin–orbit forces among the quarks and antiquark lead to a mass gap between any member of the $J^P = 1/2^+$ and its $J^P = 3/2^+$ counterpart, which was calculated in Ref. [8] to be significantly less than m_π and possibly only $O(10\text{--}50)$ MeV in the models of [5,6]. Similar remarks hold for all the members of the $\overline{\mathbf{10}}$, such as $\mathcal{E}^{+,-}$, and their non-exotic analogues that can also occur in $\mathbf{8}_5$, such as $\mathcal{E}^{0,-}$. Such a $\overline{\mathbf{10}}$ family of $J^P = 3/2^+$ states does not occur in the present formulation of chiral soliton models, nor can it if the Wess–Zumino constraint selects allowed multiplets [12].

In the Skyrme model there are exotic states with $J^P = 3/2^+$ or higher but these are in $\mathbf{27}$ and $\mathbf{35}$ multiplets of $SU(3)_F$. Such states are also expected in pentaquark models (e.g., isotensor resonance with states ranging from $uuuu\bar{s}$ with charge +3 to $ddd\bar{s}$ with charge –1) [13]. The essential difference then is that in the chiral soliton Skyrme models any spin $3/2$ partner of the Θ will exist in a variety of charge states with $I = 1, 2$ whereas the unique feature of the pentaquark models [5,6] is that the first excited state is an *isoscalar* analogue of the Θ . (There may be versions of pentaquark models where this state is higher in mass but that it is isoscalar is universal in any quark model description.)

3. Pentaquark wavefunctions, mixing and decays

In pentaquark models where the $(qqqq)$ is in $\overline{\mathbf{6}}_F$, then $\overline{\mathbf{6}} \otimes \overline{\mathbf{3}} = \overline{\mathbf{10}} \oplus \mathbf{8}_5$ leading to an $\mathbf{8}_5$ that is degenerate with the $\overline{\mathbf{10}}_5$ before mixing; chiral soliton models can accommodate an $\mathbf{8}$ (as a radial excitation of the ground state nucleon octet) though degeneracy is accidental. A challenge will be to decode the mixings between $\overline{\mathbf{10}}_5$, this $\mathbf{8}_5$ and possible contamination with excited $\mathbf{8}_3$ in experiment. This is our main focus.

The essential dynamics that underpins correlations among the flavours and spins of quarks in QCD derives from a considerable literature that recognises that ud

Table 1

Pentaquark wavefunctions where ABC are defined in the text. Note that consistency requires the meson octet to be defined with each $q\bar{q}$ positive except for $\pi^+ = -u\bar{d}$; $\bar{K}^0 = -s\bar{d}$ and then $\pi^0 = (u\bar{u} - d\bar{d})/\sqrt{2}$. In this convention $\eta_8 = (2s\bar{s} - u\bar{u} - d\bar{d})/\sqrt{6}$. For JW [5] A_1, A_2 refer to diquarks and A_3 to the antiquark; for KL [6] A_1 is diquark and $[A_2 A_3]$ is the triquark in $\bar{\mathbf{6}}_F$. The $\bar{\mathbf{6}} \otimes \bar{\mathbf{3}}$ gives the $\bar{\mathbf{10}}$ and $\mathbf{8}$ as listed in text

	$\bar{\mathbf{10}}$	$\mathbf{8}_5$
Θ^+	AAA	
p	$-(ACA + CAA + AAC)/\sqrt{3}$	$-(ACA + CAA - 2AAC)/\sqrt{6}$
n	$(ABA + BAA + AAB)/\sqrt{3}$	$(ABA + BAA - 2AAB)/\sqrt{6}$
Σ^+	$(CAC + ACC + CCA)/\sqrt{3}$	$(CAC + ACC - 2CCA)/\sqrt{6}$
Σ^0	$-(ABC + BAC + ACB + CAB + BCA + CBA)/\sqrt{6}$	$-(ABC + BAC + ACB + CAB - 2BCA - 2CBA)/\sqrt{12}$
Λ^0		$-(ABC - ACB + BAC - CAB)/2$
Σ^-	$(BAB + ABB + BBA)/\sqrt{3}$	$(BAB + ABB - 2BBA)/\sqrt{6}$
Ξ^+	$-CCC$	
Ξ^0	$(CBC + BCC + CCB)/\sqrt{3}$	$(CBC + BCC - 2CBB)/\sqrt{6}$
Ξ^-	$-(CBB + BCB + BBC)/\sqrt{3}$	$-(CBB + BCB - 2BBC)/\sqrt{6}$
Ξ^{--}	BBB	

in colour $\bar{\mathbf{3}}$ with net spin 0 feel a strong attraction [14]. This might even cause the S -wave combination to cluster as $[udu][d\bar{s}]$ which is the S -wave KN system, while the P -wave positive parity exhibits a metastability such as seen for the Θ . Two particular ways of realising this are due to Karliner and Lipkin [6] and Jaffe and Wilczek [5].

In such models the basic correlation among quarks is to form antisymmetric flavour pairs, in $\bar{\mathbf{3}}$ of $SU(3)_F$. In order to study the decays and mixings of these states it is important to have a well-defined convention for their wavefunctions [15]. We define the $\bar{\mathbf{3}}_F = (\mathbf{3}_F \otimes \mathbf{3}_F)$ basis states as

$$\begin{aligned}
 A &\equiv (ud) \equiv (ud - du)/\sqrt{2} \sim \bar{s}, \\
 B &\equiv (ds) \equiv (ds - sd)/\sqrt{2} \sim \bar{u}, \\
 C &\equiv (su) \equiv (su - us)/\sqrt{2} \sim \bar{d}
 \end{aligned} \tag{1}$$

for which $U_- A = -C$; $V_- B = -A$; $I_- C = -B$. The $\Theta^+ \equiv AAA \equiv (ud)(ud)\bar{s}$ and all other members of the $\bar{\mathbf{10}}$ follow by operating on this state sequentially by U_- and I_- until all states have been achieved. For reference they are listed in Table 1. We shall always understand the first two labels to refer to the diquarks and the rightmost to refer to the antiquark in JW [5], and for KL [6] the latter pair of labels is understood to be in the triquark. The flavour correlations in the two models are thus identical.

In addition to the three manifestly exotic combinations AAA , BBB , CCC the non-exotic states can also form an octet. In the specific dynamics advocated in [5], the quark pairs are strongly correlated into scalar pairs with colour $\bar{\mathbf{3}}$. These scalar “diquarks” are then forced to satisfy Bose symmetry, which leads naturally to the following correlations. Their colour degree of freedom is antisymmetric $\bar{\mathbf{3}} \otimes \bar{\mathbf{3}} \rightarrow \mathbf{3}$; their relative $L = 1$ provides an antisymmetric spatial state; their spin coupling is trivially symmetric; and Bose symmetry is completed by their flavour pairings being symmetric. This leads naturally to the positive parity $\bar{\mathbf{10}}$. For the $\mathbf{8}_5$ it leads to the mixed symmetric $\mathbf{8}^{Ms}$ states of Table 1; in this extreme dynamics there are no mixed antisymmetric $\mathbf{8}^{Ma}$ analogues (e.g., $p \equiv (AC - CA)A/\sqrt{2}$). This $\mathbf{8}^{Ms}$ decays to $\mathbf{8} \otimes \mathbf{8}$ with $F/D = 1/3$ as will become apparent later. Similar occurs for the KL correlation where the assumption that the triquark is in a $\bar{\mathbf{6}}_F$ implies that the pentaquark system form $\bar{\mathbf{10}} \oplus \mathbf{8}$ with the same symmetry type as in Table 1.

Thus the selection rules that we obtain are common to all these pentaquark models and a consequence of the assumed decay dynamics. The proposal of Refs. [4,10,11] is that such pentaquarks can naturally have narrow widths due to the mismatch between the colour-flavour-spin state in an initial pentaquark and the meson–baryon colour singlet states into which they decay. For a simple attractive square well potential of

range 1 fm the width of a P -wave resonance 100 MeV above KN threshold is of order 200 MeV [5,10]. However, this has not yet taken into account any price for recoupling colour and flavour-spin to overlap the $(ud)(ud)\bar{s}$ onto colour singlets uud and $d\bar{s}$ say for the KN .

If decays are assumed to arise by “fall-apart” [4,10, 11,16] without need for gluon exchange to trigger the decay (even though gluon exchange may be important in determining the eigenstates), then in amplitude, starting with the Jaffe–Wilczek configuration, the colour recoupling costs $\frac{1}{\sqrt{3}}$. It is further implicitly assumed that the fall-apart decay to a specific channel occurs only when the flavour-spin correlation in the initial wavefunction matches that of the said channel. In such a case the flavour-spin correlation to any particular channel (e.g., K^+n) costs a further $\frac{1}{2\sqrt{2}}$, hence a total suppression in rate of $\frac{1}{24}$. This was originally noted in [10].

The minimal assumption then is that a diquark must cleave such that one quark enters the baryon and the other enters the meson. *While this is necessary, implicitly it is assumed also to be sufficient:* any components in the wavefunction that are not kinematically allowed to decay are assumed to be absolutely forbidden. Selection rules that we obtain here assume this and therefore are implicitly a test of this decay dynamics. There is also a penalty for the spatial overlaps. If once organised into colour singlets, the constituents then simply fall apart in a P -wave with no momentum transfer, only the $L_z = 0$ part of the wavefunction contributes. This implies a further suppression from the $L = 1 \otimes S = 1/2 \rightarrow J = 1/2; 3/2$ coupling. Thus a total suppression of $1/72$ for the $1/2^+$ and $1/36$ for $3/2^+$ may be expected [8].

The general conclusion is that if such dynamics govern the decays, then in such models a width of $O(1-10)$ MeV for $\Theta \rightarrow KN$ may be reasonable. The above dynamics also implies that $g^2(\Theta NK^*)/g^2(\Theta NK) = 3$ (this is also implicit in [11]). Although the NK^* decay mode is kinematically inaccessible this relation may eventually be tested in photoproduction experiments [17]. Analogously this implies that $g^2(\Theta_c ND^*)/g^2(\Theta_c ND) = 3$. The Θ_c is predicted in Ref. [5] to lie below strong decay threshold but spin-orbit effects [8] could elevate its mass such that it is even above D^*N threshold (see, e.g., [6]). Thus if

$m(\Theta_c) > 2.95$ GeV, an enhanced intrinsic coupling to D^*N could be searched for.

With the wavefunctions in Table 1 we can immediately account for the relative strengths of final states by carefully exploiting the symmetries of the wavefunctions. For example, the $\Theta \equiv (ud - du)(ud - du)\bar{s}/2 \rightarrow [(ud - du)u][d\bar{s}]/2 - [(ud - du)d][u\bar{s}]/2$ which maps onto $\Theta \rightarrow pK^0/\sqrt{2} - nK^+/\sqrt{2}$.

These amplitudes for decays into meson (M) and baryon (B) also depend on the flavour-spin symmetry of the baryon. If we make this explicit (ϕ, χ referring to the flavour and spin wavefunctions respectively and M_A, M_S denoting their mixed symmetry properties under interchange [18]) we have

$$(\bar{\mathbf{3}}_F, S = 0)(\bar{\mathbf{3}}_F, S = 0) \rightarrow M + B(\phi^{M_A} \chi^{M_A}).$$

The same colour-orbital configuration for tetraquarks ($qqqq$) in overall spin $S = 0$ can be realised with diquarks in $\mathbf{6}_F, S = 1$. The pattern of decays from this configuration mirror those above except that the baryon’s flavour-spin symmetry is swapped

$$(\mathbf{6}_F, S = 1)(\mathbf{6}_F, S = 1) \rightarrow M + B(\phi^{M_S} \chi^{M_S}).$$

Thus if one imposed overall antisymmetry on the tetraquark wavefunction one encounters for the flavour-spin part of the wavefunction

$$|(\bar{\mathbf{3}}_F, S = 0)(\bar{\mathbf{3}}_F, S = 0)\rangle \pm |(\mathbf{6}_F, S = 1)(\mathbf{6}_F, S = 1)\rangle.$$

Noting that there is an $L = 1$ within the ($qqqq$) system, the above wavefunctions imply that the (+) phase decays to $M + B(\mathbf{56})$ in a P -wave and the (−) phase decays to $M + B(\mathbf{70}(L = 1))$ in an S -wave. The latter would naively be kinematically forbidden and as such lead to a suppressed width if the Θ were in this representation (which is in the **105**-dimensional mixed symmetry representation of flavour-spin [16]). However, one needs also to confront the kinematically allowed decays to $M + B(\mathbf{56})$ from the (+) phase state (in the symmetric **126** representation and discussed in [11]). In practice decays shared by the $(\bar{\mathbf{3}}, S = 0) \times (\bar{\mathbf{3}}, S = 0)$ and $(\mathbf{6}, S = 1)(\mathbf{6}, S = 1)$ states lead to mixing. If this is stronger than the mass gap between these two states, then one would obtain the above two configurations, leading to the possibility of the Θ as a narrow state in **105** partnered by a (yet unobserved) broad partner in **126** (see also [19]). By contrast, if the mixing is small on the scale of the mass

gap, the wavefunction of the light eigenstate in this spin-zero tetraquark sector will be dominated by the $(\bar{\mathbf{3}}, S=0)(\bar{\mathbf{3}}, S=0)$ configuration, which is the Jaffe–Wilczek model [5]. As noted in Ref. [10], mixing with the spin-one tetraquark sector as manifested in the KL correlation can lead to a lower eigenstate. The discussion in the rest of our present Letter does not depend on this dynamical question.

Within the assumption that decays are driven by the fall-apart dynamics, the flavour patterns follow for all of these configurations.

It is especially instructive to apply our study of the fall-apart to the \mathcal{E}_5 states. In what follows we assume that only the ground state baryon + 0^- meson channels are kinematically accessible. If other channels such as 1^- mesons could be accessed these would cause the intrinsic suppression to be less dramatic.

4. Decays of \mathcal{E}_5 states

Starting with the wavefunction (see Table 1)

$$|\mathcal{E}^-(\bar{\mathbf{10}})\rangle = -\frac{1}{2\sqrt{3}} \left([(ds - sd)(su - us) + (su - us)(ds - sd)]\bar{u} + (ds - sd)^2\bar{d} \right)$$

we can rewrite this in flavour space in the form $(qqq)(q\bar{q})$.

$$|\mathcal{E}^-(\bar{\mathbf{10}})\rangle = -\frac{1}{2\sqrt{3}} \times \left([(ds - sd)s](u\bar{u} - d\bar{d}) + [(su - us)d + (sd - ds)u](s\bar{u}) - [(su - us)s](d\bar{u}) + [(ds - sd)d](s\bar{d}) \right)$$

which maps onto the following ground state hadrons

$$\mathcal{E}^-(\bar{\mathbf{10}}) \rightarrow -\frac{1}{\sqrt{6}} (\sqrt{2}\mathcal{E}^-\pi^0 + \mathcal{E}^0\pi^- - \sqrt{2}K^-\Sigma^0 - \bar{K}^0\Sigma^-).$$

These agree in relative magnitudes and phases with the standard de Swart results [20,21]; they agree in relative magnitudes with Oh et al. [22] but their phases differ from ours. Refs. [9,22] do not discuss the $\mathbf{8}$ decays as these depend in general on an undetermined

F/D . However with the pentaquark wavefunctions, as specified as in Table 1, the octet from $\bar{\mathbf{6}}_F \otimes \bar{\mathbf{3}}_F$ that is orthogonal to the $\bar{\mathbf{10}}$ is

$$|\mathcal{E}_5^-(\mathbf{8})\rangle = -\frac{1}{2\sqrt{3}} \left([(ds - sd)(su - us) + (su - us)(ds - sd)]\bar{u} - \sqrt{2}(ds - sd)^2\bar{d} \right)$$

and for the assumed decay dynamics employed in [4, 10,11], the particle decomposition is

$$\mathcal{E}_5^-(\mathbf{8}) \rightarrow -\frac{1}{\sqrt{24}} (\sqrt{6}\mathcal{E}^-\eta_1 - \sqrt{3}\mathcal{E}^-\eta_8 - \mathcal{E}^-\pi^0 + \sqrt{2}\mathcal{E}^0\pi^- + 2\sqrt{2}\Sigma^-\bar{K}^0 - 2\Sigma^0K^- + 0\Lambda K^-)$$

which corresponds to $\mathbf{8} \rightarrow \mathbf{8} \otimes \mathbf{8}$ with $F/D = 1/3$ (or $g_1 = \sqrt{5}g_2$ in the de Swart convention [20,21]). With this one can therefore deduce the branching ratios for N, Σ, Λ states in $\mathbf{8}_5$ immediately from existing tables [20,21] and we do not discuss them further here.

For the \mathcal{E}_5 we see immediately distinctions between the two states.

- (i) Isospin ($I = 3/2$ versus $I = 1/2$) is responsible for the distinctive ratios

$$\frac{\Gamma(\mathcal{E}_5^- \rightarrow \mathcal{E}^-\pi^0)}{\Gamma(\mathcal{E}_5^- \rightarrow \mathcal{E}^0\pi^-)} = \begin{cases} 1/2, & \mathbf{8}, \\ 2, & \bar{\mathbf{10}} \end{cases}$$

and analogous for the ΣK modes.

- (ii) There is a selection rule that ΛK^- modes vanish. For the $\bar{\mathbf{10}}$ this is a trivial consequence of isospin; for the $\mathbf{8}_5$ it is a result of the pentaquark wavefunction, in particular that the $qqqq$ flavour wavefunction of the pair of diquarks is symmetric in flavour, (i.e., $\bar{\mathbf{6}}_F = \bar{\mathbf{3}}_F \otimes \bar{\mathbf{3}}_F$) leading to $F/D = 1/3$.

A pedagogic explanation of the selection rule is as follows. The \mathcal{E}_5 state wavefunctions contain two pieces of generic structure $(dssu)\bar{u}$ and $(dssd)\bar{d}$. The $I = 3/2$ and $I = 1/2$ states differ in the relative proportions of these two. However, only the first component $(dssu)\bar{u}$ contains the \bar{u} required for the K^- and this is common to both the $\mathcal{E}(I = 3/2)$ and $\mathcal{E}(I = 1/2)$. Thus as the $\mathcal{E}(I = 3/2) \rightarrow K\Lambda$ is trivially forbidden by isospin, the $\mathcal{E}(I = 1/2) \rightarrow K\Lambda$ must be also unless there is cross-talk between the two components in the wavefunction. This would happen

if annihilation $(dssu)\bar{u} \rightarrow (dss) \rightarrow (dssd)\bar{d}$ occurs. Thus observation of ΛK^- could arise if there are admixtures of $\mathbf{8}_3$ in the wavefunction.

Rescattering from kinematically forbidden channels, such as $\mathcal{E}\eta$ can feed both $K\Sigma$ and $K\Lambda$, though this is not expected to be a large effect if experience with light hadrons is relevant (such as the small width of the $f_1(1285)$ not being affected by rescattering from the kinematically closed KK^* channel, and the predicted $\pi_2 \rightarrow b_1\pi \sim 0$ [23] not being affected by rescattering from the allowed channels $\pi f_2; \pi\rho$). Whether this carries over to pentaquarks may be tested qualitatively in models by comparing the relative suppression of Θ , \mathcal{E}^{--} and \mathcal{E}^- states; if there is no rescattering and the $\mathcal{E}\eta$ channels are closed in the initial pentaquark wavefunction, its width will be further suppressed from $1/24 \rightarrow \sim 3/115$ and $K\Lambda \sim 0$. In this case the width of \mathcal{E}^- (after phase space effects have been removed) will be less than that of \mathcal{E}^{--} . A dominance of $K\Lambda > K\Sigma$ can arise if there are pentaquark configurations having $F = D$. In this latter case the $\Sigma^0 K^-$ would be forbidden but ΛK^- allowed. The $\Lambda K : \Sigma K$ ratio in general can be used to constrain the F/D ratio and begin to discriminate between various dynamical schemes.

Decays to $\mathcal{E}^*\pi$ and Σ^*K for Σ^* , \mathcal{E}^* in the $\mathbf{10}$ are forbidden (even if allowed by phase space). For the $\bar{\mathbf{10}}$ this is a result of $\bar{\mathbf{10}} \neq \mathbf{8} \otimes \mathbf{10}$ as noted in Ref. [9] who also discuss $SU(3)_F$ breaking as a potential source of violation of this zero. However, this selection rule may be stronger in the pentaquark models of Refs. [5,6] due to the diquarks having antisymmetric flavour ($\mathbf{3}$) and spin zero, both of which prevent simple overlap of flavour-spin with the $\mathbf{10}$, $S = 3/2$ baryon decuplet resonances. Thus although $SU(3)_F$ allows $\mathbf{8} \rightarrow \mathbf{10} \otimes \mathbf{8}$ to occur, for the $\mathbf{8}_5$ states of Table 1 it is again forbidden as a consequence of the antisymmetric flavour content of the wavefunction, at least within the models of suppressed decay widths considered here. While we discussed this for the Jaffe–Wilczek wavefunction, Karliner and Lipkin have one of their quark pairs strongly correlated into a vector spin state within a triquark (e.g., $ud\bar{s}$) so the flavour antisymmetries and explicit scalar diquark in the residual wavefunctions suggest that this dynamics also would be challenged to accommodate a violation of this selection rule.

The $I = 3/2$ states will all be narrow. They are degenerate up to electromagnetic mass shifts. Across

the $I = 3/2$ multiplet the mass split is $\mathcal{E}^{--} - \mathcal{E}^+ = (d-u) + \langle e^2/R \rangle$ [18] where the Coulomb contribution in known hadrons is $\sim 2-9$ MeV, hence a spread of 3–10 MeV is expected. For the non-exotic states $m(\mathcal{E}_{\mathbf{8}_5}) > m(\mathcal{E}_{\bar{\mathbf{10}}})$, with

$$\begin{aligned} m(\mathcal{E}_{\mathbf{8}_5}^0) - m(\mathcal{E}_{\bar{\mathbf{10}}}^0) &= \frac{1}{2} [m(\mathcal{E}_{\mathbf{8}_5}^-) - m(\mathcal{E}_{\bar{\mathbf{10}}}^-)] \\ &= \frac{1}{3} [m(d) - m(u)] \sim 1.5-2.5 \text{ MeV} \end{aligned}$$

and hence degenerate to within better than 5 MeV. If the coupling to $\mathcal{E}^*\pi$ vanishes for the $\mathbf{8}_5$ as well as the $\bar{\mathbf{10}}$, then mixing by the common $\mathcal{E}\pi$ decay channels will be destructive. If the widths are truly narrow the mass eigenstates become $\mathcal{E}_u \equiv (ds)(su)\bar{u}$ and $\mathcal{E}_d \equiv (ds)(ds)\bar{d}$ separated by ~ 10 MeV. The heavier state $b.r.(\mathcal{E}_d \rightarrow \mathcal{E}^-\pi^0) = 2 \times b.r.(\mathcal{E}_d \rightarrow \Sigma^- \bar{K}^0)$ (apart from phase effects) but it does not decay to either $\mathcal{E}^0\pi^-$ nor $\Sigma^0 K^-$. In contrast the lighter state $\mathcal{E}_u \rightarrow \mathcal{E}^0\pi^- : \mathcal{E}^-\pi^0 = 2$, as for the pure $\mathbf{8}$ (but with opposite relative phase), while it does not decay to $\Sigma^- \bar{K}^0$.

Violation of these relations would imply either mixing with excited $\mathbf{8}_3$ states, be due to pentaquark components in the wavefunction beyond those above, or because the width suppression is realised by some dynamics other than implicit in Refs. [4,10,11]. In the former case one would expect the $\mathbf{8}_3$ components to decay without suppression and dominate the systematics of the widths. In this case there will be narrow $\mathcal{E}^{-,0}$ with $I = 3/2$ partnering the exotic $\mathcal{E}^{+,-}$ and broad $I = 1/2$ states that are akin to normal excited \mathcal{E} states. By contrast, were the $\mathcal{E}\pi$ charge ratios to show mixing between the two \mathcal{E}_5 states with two narrow states such as \mathcal{E}_d and \mathcal{E}_u , then observation of any ΛK or $\mathcal{E}^*\pi$ would require components in the pentaquark wavefunction with different symmetries to those above.

5. Decays of p_5 and n_5 states

These follow immediately from $SU(3)$ tables with $F/D = 1/3$. In general there will be mixing between these as suggested by Jaffe and Wilczek. For the

extreme $p_5(s\bar{s})$ and $p_5(d\bar{d})$ we have

$$p_5(s\bar{s}) \rightarrow \frac{1}{2} \left(\frac{1}{\sqrt{2}} \Sigma^0 K^+ + \sqrt{\frac{3}{2}} \Lambda K^+ - \Sigma^+ K^0 + p\eta_s \right),$$

where $\eta_s \equiv \eta_1/\sqrt{3} + 2\eta_8/\sqrt{6}$; and while phase space only admits trivial $p_5(d\bar{d})$ decays to $N\pi$.

It is immediately apparent that the decays of $P_{11}(1440; 1710)$ do not fit well with this scheme. First, there is a dominance of non-strange hadrons in the heavier $P_{11}(1710)$ with prominent $\Delta\pi$ in the decays of both $P_{11}(1440; 1710)$. This mode is not possible for the p_5 states in $\bar{\mathbf{10}}$ nor in $\mathbf{8}_5$ unless overwritten by rescattering or mixing with $\mathbf{8}_3$.

It is clear that $P_{11}(1440)$ is partnered by $P_{33}(1660)$ as in a traditional $\mathbf{56}$ multiplet of $SU(6)$ qqq states. There is no obvious sign of pentaquarks here. A possibility is that the states are linear combinations of p_3 and p_5 ; the p_5 could even dominate the wavefunction but its $O(1 \text{ MeV})$ width is swamped by the $O(100 \text{ MeV})$ width of the unsuppressed p_3 component. The p_5 decays listed above would then show up as rare decays at the $O(1\%)$ level.

6. Λ_5 state with $J^P = 3/2^+$

There is one further potentially narrow state in pentaquark models, which has little opportunity for mixing with qqq states. This is the Λ_5 state that is the $J^P = 3/2^+$ spin-orbit partner of Λ_5 in $\mathbf{8}_5$.

First note that $\bar{\mathbf{10}}$ contains Σ_5 but has no Λ_5 . The $\mathbf{8}_5$ contains a Λ_5 , and there will be no mixing with $\bar{\mathbf{10}}$ so long as isospin is good. If there were no mixing with $\Lambda(qqq)$ excited states, this Λ_5 would be narrow, with width identical to that of the Θ apart from phase space factors.

The Λ_5 wavefunction shows that it has only one strange mass quark and hence is similar to the Θ in this regard. Ref. [5] estimate $\sim 1600 \text{ MeV}$ for such a state (the excess $\sim 60 \text{ MeV}$ relative to the Θ arising because the mass of a $(us)\bar{d}$ set is larger than $(ud)\bar{s}$ due to the relatively smaller downward mass shift in the (us) diquark). Scaling the spin-orbit splitting from [8] and allowing for the relative masses of the \bar{s}/\bar{d} and $m(us)/m(ud)$ gives 40–70 MeV for the Λ_5 mass gap

of $3/2^+ - 1/2^+$ and hence 1600–1700 as a conservative estimate for the mass range for the partner $\Lambda_5(3/2^+)$.

Perusal of the data [21] shows that, for the $1/2^+$, mixing with qqq states is likely (given the existence of a candidate $\mathbf{56}, 0^+$ multiplet containing $P_{11}(1440)$, $\Lambda(1600)$, $\Sigma(1660)$, $\Xi(?)$). However there is no $3/2^+$ multiplet with a $\Lambda(1600-1700)$ seen, nor is one expected in standard qqq models. The first such is the set containing $P_{13}(1720)$, $\Lambda(1890)$, \dots . Thus there is a significant gap between $\Lambda(1890)$ and our predicted $\Lambda_5(3/2^+)$.

The branching ratios for either the spin $1/2$ or $3/2$ states can be determined from the breakdown

$$\Lambda_5 \rightarrow -\frac{1}{2\sqrt{2}} (pK^- - n\bar{K}^0 + \Sigma^- \pi^+ - \Sigma^+ \pi^- - \Sigma^0 \pi^0 - \sqrt{3} \Lambda \eta \eta \bar{n}).$$

Decays to $\Sigma^* \pi$ should be suppressed, even if they are kinematically accessible. The production rate of the spin $1/2$ state in $\gamma p \rightarrow K^+ \Lambda_5$ should be similar to that of $\gamma n \rightarrow K^- \Theta$ (perhaps a factor of four smaller if K exchange drives the production and $g(KN\Theta) = 2g(KN\Lambda_5)$). If the arguments about $L \otimes S$ coupling and fall-apart dynamics are correct, then we can expect the spin $3/2$ state to be enhanced by a factor of two relative to the spin $1/2$ counterpart. A search in $\gamma p \rightarrow K^+ \Lambda_5$ therefore seems appropriate.

If $\bar{\mathbf{10}}-\mathbf{8}_5$ mixing is ideal, then also charged Σ_d^\pm states will occur which for $J^P = 3/2^+$ should be unmixed. For $J^P = 1/2^+$, the amplitudes $g(\Theta^+ K^+ n) = \sqrt{2}g(\Sigma_5^- K^- n)$ and so the relative photoproduction cross sections should scale as $\sigma(\gamma n \rightarrow K^- \Theta^+) \sim 2 \times \sigma(\gamma n \rightarrow K^+ \Sigma_5^-)$ [17]. If the Σ_5 is mixed into the $\Sigma(1660)$ then the latter state should be photoproduced at least at the above rate and so may be a test for consistency.

7. Summary

We advocate study of decays of the Ξ states, especially the ratios of various charge modes, and searching for $\Xi^* \pi$ and ΛK as tests of the underlying dynamics that forms the states. We also stress the importance of isolating the $J^P = 3/2^+$ states that must occur in $\bar{\mathbf{10}}$ and $\mathbf{8}_5$ in pentaquark models but which have no analogue in the chiral soliton model. These states are predicted to be within a few tens of

MeV of their $1/2^+$ counterparts in highly correlated models such as those of Jaffe–Wilczek or Karliner–Lipkin. Were the mass gap to be significantly larger, then it could point to the presence of other components in the pentaquark wavefunction. By contrast, the absence of such states together with the appearance of $J^P = 3/2^+$ in higher representations such as **27** or **35** would support the chiral soliton models. The $\Lambda_5(3/2^+)$ state may be relatively light and narrow and should be produced with similar strength to Θ in photoproduction. Its confirmation could play a significant role in helping to decode the mixing between pentaquarks and conventional states.

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