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Dibaryons and where to find them

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Abstract

In recent years, there has been tremendous progress in the investigation of bound systems of quarks with multiplicities beyond the more usual two- and three-quark systems. Experimental and theoretical progress has been made in the four-, five- and even six-quark sectors. In this paper, we review the possible lightest six-quark states using a simple ansatz based on SU(3) symmetry and evaluate the most promising decay branches. The work will be useful to help focus future experimental searches in this six-quark sector.

Keywords: dibaryons, hexaquarks, baryon–baryon molecules, light and strange quark systems

1. Introduction

Theories of the strong interaction indicate there could be more ways to combine multiple quarks into a single object than established experimentally [1][1]. The recent experimental discoveries of four- and five-quark systems [1] has led to renewed interest in the field of multiquark states—where 'multiquark' refers to objects with 'non-standard' quark configurations beyond that of baryons and mesons. Unfortunately, for many of these multiquark systems, the internal structure is not currently established. Due to this, CERN has adopted a new scheme in referring to these states[2]. Tetraquarks refer to all objects with four quarks inside regardless of whether it is a genuine 'tetraquark' or a meson-meson molecule,

[1] See e.g. Chapter 15. Quark Model PDG overview of current situation.

[2] It is very likely that most pentaquarks, which PDG [1] refer as P(XXXX) in a new 'exotic baryons' chapter are in reality meson-baryon molecules.

Table 1. Six non-strange dibaryons.

I, J	D. and X. prediction [MeV]	Experimental results [MeV]	SU(3) multiplet	References
0, 1	1876	1876	10^*	[2]
1, 0	1876	1878	27	[12, 13]
1, 2	2160	2160	27	[62, 63]
2, 1	2160	2160	35	[69]
0, 3	2350	2380	10^*	[26, 27, 32]
3, 0	2350	2464	28	[60]

pentaquarks for five-quark objects (genuine pentaquarks and meson-baryon molecules) and hexaquarks (genuine hexaquarks and baryon–baryon molecules). We adopt this system in the current paper—but note that some of the earlier literature adopted different systems with hexaquark referring only to the non-molecular states.

The hexaquark family can be further decomposed by quark content since the six-quark system can be arranged in two possible ways with three quarks and three antiquarks (*baryonium*) and six quarks (*dibaryon*). In this paper, we consider the latter configuration and develop a simple theoretical ansatz so that we can infer the likely decay properties of such 'dibaryon' six quark states.

It is interesting to note that hexaquarks were the first multiquark states to be established. The deuteron, a trivial hexaquark composed (predominantly) of a molecular state of a proton and neutron, was discovered in 1931 [2]. The earliest paper on non-trivial hexaquarks is attributed to Dyson and Xuong [3], who predicted the existence of six non-strange hexaquarks based on the SU(6) model just half a year after the discovery of quarks [4]. All of the states predicted by Dyson and Xuong have been the subject of experimental searches in the subsequent years. In table 1 below, we summarise the experimental signals for the states of different isospin I and spin J, along with their associated SU(3) multiplet.

As shown in the table, all states in this u,d quark sector have given experimental signatures with properties broadly consistent with the predicted properties. Signatures for some of the members have only been obtained very recently. A detailed review of the progress in dibaryon searches and their observations in some of the channels, can be found in [5].

Dibaryons, like any other particles, appear in SU(3) multiplets. Since SU(3) symmetry works reasonably well, one can use this symmetry to provide expectations of the properties of the family of particles in the multiplet (spin, isospin, mass).

In their pioneering paper, Dyson and Xuong [3] focused on the u,d quark sector and did not consider the properties of strange quark containing members of dibaryon multiplets or their decay branches. However, modern experimental facilities (JLab, J-PARC, ALICE) now offer the prospect of experimental study of dibaryons in the strange quark sector. In this paper, we will review what is known about these mysterious strange quark-containing particles, point to the most promising experimental channels and, where feasible, estimate their decay branches.

The paper is structured as follows. In section 2 we review the current experimental evidence and theoretical interpretation of the deuteron anti-decuplet. In section 3 we cover the complimentary NN-FSI 27-plet. Section 4 addresses the anti-decuplet based on the $d^*(2380)$ hexaquark. Following a review we present our new theoretical estimate based on SU(3)

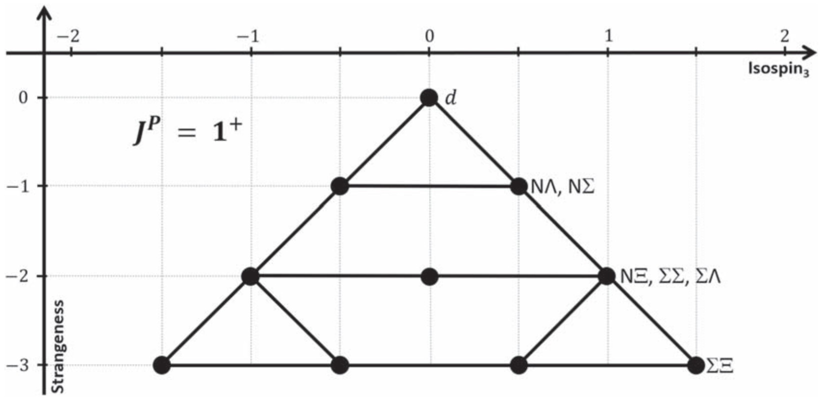

Figure 1. Deuteron multiplet.

Table 2. Deuteron SU(3) multiplet.

Strangeness	Decomposition	Masses of components[MeV]	Binding [MeV]
0	NN	NN(1878)	−2.2
−1	$\frac{1}{\sqrt{2}}(\mathrm{N}\Lambda - \mathrm{N}\Sigma)$	NΛ(2054) NΣ(2129)	+0.166
−2	$\frac{1}{\sqrt{6}}(\sqrt{2}\,\mathrm{N}\Xi + \Sigma\Sigma + \sqrt{3}\,\Sigma\Lambda)$	NΞ(2253) ΣΣ(2382) ΣΛ(2305)	+0.5
−3	ΣΞ	ΣΞ(2504)	+1.0

symmetry considerations for the decay branches of the strange quark containing partners in this anti-decuplet under assumptions of pure hexaquark and pure molecular natures of the d^* multiplet members. The SU(3) Clebsch–Gordan coefficients for these evaluations were taken from [6, 7]. We than use the same theoretical approach to investigate decay properties of the ΔΔ 28-plet (sSection 5), $N\Delta$ 27-plet (section 6) and $N\Delta$ 35-plet (section 7).

2. Deuteron antidecuplet

The deuteron SU(3) multiplet is probably the most well-studied hexaquark SU(3) family. We know that the deuteron, and hence all the other objects in this multiplet, are predominantly baryon–baryon molecules. The estimated genuine hexaquark $|6q>$ component of the deuteron wavefunction is predicted to be only around 0.15% [8]. Figure 1 shows the members of the deuteron multiplet in standard convention (x-axis—3rd projection of isospin, y-axis—strangeness). For each member of this antidecuplet we also show which $8 \oplus 8$ states can contribute. The SU(3) decomposition of this multiplet is also summarised in table 2.

Out of all multiplet members, only the deuteron appears to be bound. The $\Lambda - \mathrm{N}$ interaction is not attractive enough to make a bound state e.g. the hyperdeuteron appears to be unbound.[3] The ΛN state therefore presents as a virtual state or a final-state interaction (FSI) [9]. The $\Sigma - \mathrm{N}$ interaction is suggested to be less attractive than $\Lambda - \mathrm{N}$ [9]. Due to the large

[3] With additional particles, the nuclear binding would be expected to increase—as a result, hypertritium does become bound but by only 150 keV.

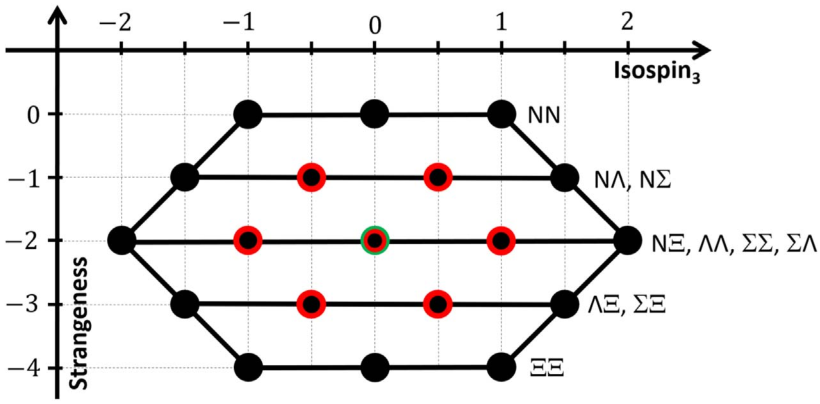

Figure 2. 27-plet $8 \oplus 8$.

Table 3. Expected decay branches of the $8 \oplus 8$ SU(3) multiplet.

S	Max Isospin	Med Isospin	Min Isospin	Mass [MeV]
0	NN			NN(1876)
−1	NΣ	$\frac{1}{\sqrt{10}}(3\mathrm{N}\Lambda + \mathrm{N}\Sigma)$		NΣ(2127) NΛ(2054)
−2	ΣΣ	$\frac{1}{\sqrt{10}}(\sqrt{6}\,\Sigma\Lambda + 2\mathrm{N}\Xi)$	$\frac{1}{\sqrt{40}}(\sqrt{27}\,\Lambda\Lambda + \sqrt{12}\,\mathrm{N}\Xi - \Sigma\Sigma)$	ΣΣ(2378)ΣΛ(2305) NΞ(2253)ΛΛ(2232)
−3	ΣΞ	$\frac{1}{\sqrt{10}}(3\Lambda\Xi + \Sigma\Xi)$		ΣΞ(2504)ΛΞ(2431)
−4	ΞΞ			ΞΞ(2630)

SU(3) mass splitting in this decuplet ($M_\Sigma - M_\Lambda \sim 77$ MeV), compared to the MeV–keV range of bound/virtual states, it is difficult to treat these states in an SU(3) limit.

Aside from the deuteron, none of the states from the multiplet form a bound state, so all of them can be observed only in elastic/quasi-elastic reactions or via the correlation observations in heavy ion collisions. There were several recent publications from STAR and ALICE about extraction of the ΛN [10], ΣN [10] and ΞN [11] scattering lengths. Due to very unfavourable decay branches in the Σ family, it is extremely difficult to access any correlation distributions where the Σ baryon is involved. However, one may get some information in $\Xi\mathrm{N} \to \Sigma\Sigma$ or $\Xi\mathrm{N} \to \Lambda\Sigma$ from KLF or JPARC facilities from tertiary Ξ beams.

3. NN-FSI 27-plet

The 27-plet $8 \oplus 8$ of spin zero states are shown schematically in figure 2 and table 3 details the quark compositions. A strong attraction in the spin-zero pp-system was probably first realised by Migdal [12] and Watson [13] in their description of NN-FSI (Final State Interactions). Indeed it appears the absence of the tensor force in the NN $^1S_0$ channel implies the 'demon deuteron' would not form, as this system is unbound by 66 keV. Thus this multiplet is expected to be comprised of unbound molecular-like virtual states since the strangeness zero members are established to be unbound. As is the case for the deuteron multiplet (section 2), this SU(3) family suffers from large SU(3) mass splitting. All the members can be observed only in elastic/quasi-elastic reactions and/or correlations in heavy-ion collisions.

Although this multiplet in isolation would not be expected to produce bound states, the presence of overlapping members with the same quark content but different isospin raises the question about mixing. Effects established in the meson sector, such as $\rho - \omega - \phi$ mixing in the case of meson nonet [14], could, in principle, have corollaries here. Several members of this 27-plet can be mixed, but the most famous case proposed is the so-called H-dibaryon, a potential $\Lambda - \Lambda$ deeply bound state [15]. Experimentally this is now shown to be unbound [16–19]. Recent advances in Lattice-QCD calculations for this sector indicate it to be located in the vicinity of the NΞ-threshold, [20].

In a simplified picture of this multiplet, we have three states with isospins 2,1,0 (plus a SU(3) singlet isosingle H-dibaryon state). These states can mix, but only in an isospin-violating way, analagous to $\Lambda - \Sigma$ mixing in the baryon octet or $\rho - \omega$ mixing in meson nonet. The redistribution of wave function components, similar to the $\omega - \phi$ vs $\phi_0 - \phi_8$ in the meson nonet, is not allowed here since all states have different isospin(besides ${}^{27}d_{ss}(I = 0) - H$-dibaryon mixing).

All three states have dissimilar properties: the $I = 2$ is a pure $\Sigma\Sigma$ state, the $I = 1$ is $(\sqrt{\frac{2}{5}}|\mathrm{N}\Xi > +\sqrt{\frac{3}{5}}|\Sigma\Lambda>)$ and $I = 0$ $(\sqrt{\frac{12}{40}}|\mathrm{N}\Xi > -\sqrt{\frac{1}{40}}|\Sigma\Sigma > +\sqrt{\frac{27}{40}}|\Lambda\Lambda>)$. The $|I, I_z> = |1, 0>$ state does not couple to $\Sigma\Sigma$ and $|I, I_z> = |0, 0>$ state does not couple to $\Sigma\Lambda$. The $\Sigma\Sigma$ component of an $I = 0$ state is tiny—such that no mixing with $I = 2$ state is expected. The $I = 0$ and the $I = 1$ states can mix only via the $N\Xi$ component since it is the only common part in their wave functions. The H-dibaryon $(|H> =\sqrt{\frac{1}{2}}|\mathrm{N}\Xi > -\sqrt{\frac{3}{8}}|\Sigma\Sigma > -\sqrt{\frac{1}{8}}|\Lambda\Lambda>)$ belongs to a singlet state, but it has exactly the same quantum numbers, including isospin, as the central state of the 27-plet. One could therefore expect some mixing between these states. However, since no H-dibaryon was found in the vicinity of the $\Lambda\Lambda$ threshold, the $I = 0$ state from the 27-plet should be even less bound and should have much weaker experimental signatures.

Since, for an experimental observation of this multiplet we are limited to elastic scattering and heavy-ion correlations, there are very few states which we can currently access. An elastic scattering with either direct beams or rescattering of secondary beams within target material [21] can give us access to NN, ΛN, ΣN, ΞN scattering data. The latter reaction can be performed at the recently proposed KLF facility [22] with $K_L\mathrm{N} \to K^+\Xi$ as a first step and $\Xi\mathrm{N} \to X$ as a second step reaction within a large target volume (or with similar methodology at J-PARC).

For the heavy-ion correlation searches, all events with neutral final states need to be excluded from realisable experimental study, which essentially removes all channels with Σ baryons. The NN, ΛN, ΞN and $\Lambda\Lambda$ channels were already investigated, see [10, 11, 18, 19]. The only other channels pending analysis for this multiplet is $\Xi\Xi$ (the $\Xi\Lambda$ ALICE analysis was recently published [23], but it suffers from low statistics). We do not expect any breakthrough here—with the expectation of loosely unbound states. However, these data would help to constrain several ChPT coupling constants in the baryon–baryon sector [24] and also improve our understanding of the nuclear equation of state for a system with strangeness, essential for neutron star physics [25].

4. d^* (2380) antidecuplet

The antidecuplet of states built on the six-quark containing $d^*(2380)$ (hexaquark) are shown schematically in figure 3. Signatures of the $d^*(2380)$ have been established quite rigorously in

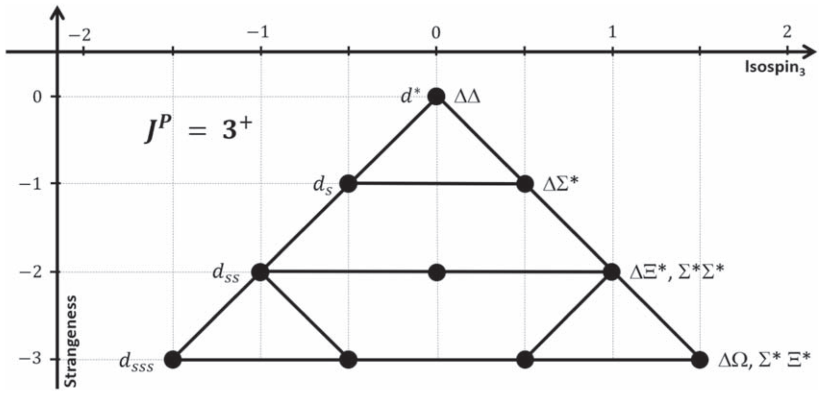

Figure 3. d^*(2380) multiplet.

recent years following its initial observation in proton-neutron scattering and pionic fusion reactions [26–34]. It has a mass of $M_{d*} = 2380$ MeV, vacuum width $\Gamma = 70$ MeV and quantum numbers $I(J^P) = 0(3^+)$. All of the strong decay branches have been identified and measured in experiment [34]. The electromagnetic properties of the d^*(2380) were also investigated recently from measurements of its photoexcitation from deuteron targets [35–38], with further programmes planned [39]. This particle has also recently been suggested to have a potential impact in astrophysics [25, 40, 41].

The search for strange SU(3) partners of the d^*(2380) is a natural next step for hexaquark studies. Characterisation of additional members of the antidecuplet would provide valuable insights into the underlying physics of hexaquark systems. All previous multiplets mentioned above have at best a very low binding suggesting its molecular nature. In contrast the d^*(2380) has a very large binding energy, which unavoidably means it should have a large $|6q\rangle$ component even if one considers it from a molecular picture. Moreover, the majority of theoretical papers in different ansatz suggest that the d^*(2380) might indeed be a genuine hexaquark-dominated object [42, 44, 45]. Its evidence in photoproduction (e.g. relative contribution of multipoles) also favours this possibility [35, 38].

However, as there is an abundance of states and associated decay channels from the many states in the d^*(2380) antidecuplet, some prioritisation of initial searches is beneficial. In this paper, we develop a simple model to identify the most appropriate final states for each member of the antidecuplet, including estimates of the partial decay widths. Due to conservation laws, the production of any of the multiplet members is expected to proceed via associated many-body reactions at rather high energies. Deriving accurate predictions of the production cross-sections are therefore challenging. In their absence, an experimental strategy to search for such states is to look for structure in the cross sections consistent with the predicted locations, widths and branching ratios derived from theoretical models. In the following sections we present the first such calculations to guide future measurement strategies. The properties of the decouplet are inferred under assumption of a hexaquark (section 4.1.1) and a molecular state (section 4.1.2).

4.1. *d^*(2380) multiplet masses*

From the unitary group theory of the strong interaction for the light quark (u, d, s) sector any strongly interacting particle, such as the d^*(2380), should be part of a SU(3) multiplet. For the

Table 4. Expected masses of the $d^*(2380)$ SU(3) multiplet in pure genuine hexaquark picture.

Particle	Mass structure	Mass value [MeV]	Binding[a] [MeV]
d^*	$6M_n + \frac{15K}{M_n^2} - B_h$	2380	84
d_s	$5M_n + M_s + \frac{10K}{M_n^2} + \frac{5K}{M_n M_s} - B_h$	2474	141
d_{ss}	$4M_n + 2M_s + \frac{6K}{M_n^2} + \frac{8K}{M_n M_s} + \frac{K}{M_s^2} - B_h$	2573	193
d_{sss}	$3M_n + 3M_s + \frac{3K}{M_n^2} + \frac{9K}{M_n M_s} + \frac{3K}{M_s^2} - B_h$	2677	234

[a] Relative to Decuplet–Decuplet pole.

case of the $d^*(2380)$ it would be expected to be a member of an antidecuplet, figure 3. The spectroscopic study of the other multiplet members would provide important new constraints on the d^* internal structure, complementary to that potentially achievable in form factor studies. In a simple molecular picture, the d^* SU(3) multiplet would derive from the coupling of two baryon decuplet members bound by long-range pion exchange: corresponding to $\Delta\Delta$ for the $d^*(2380)$ to $\Delta\Omega$ for the d_{sss}. However, since the pion does not have a coupling to strange quarks, there is an expectation for such molecular systems that the binding energy should decrease with increasing strangeness content of the state. Conversely, in a genuine hexaquark (non-molecular) picture the binding energy should *increase* with increasing strangeness content as the presence of heavier s-quarks would imply stronger binding. In these two cases, a rough evaluation can be made for the masses of the $d^*(2380)$ multiplet.

4.1.1. $d^*(2380)$ multiplet as genuine hexaquarks. The work of Gell-Mann presented a way to calculate masses within an SU(3) multiplet [46], which can be further simplified if we consider SU(6) symmetry where the baryon mass splitting originates from colour-magnetic interactions between quarks [14]. In this approach the masses originate from two factors—the masses of the constituent quarks (M_q) and a contribution from hyperfine splitting ($\frac{K}{M_{q_1} M_{q_2}}$). In this approach the mass difference between the nucleon and the Δ (which have the same quark content) can be reproduced: $M_N = 3M_q - \frac{3K}{M_q^2}$, $M_\Delta = 3M_q + \frac{3K}{M_q^2}$. Here $M_q = 363$ MeV is the masses of light, unflavored constituent quarks and the K is the parameter responsible for hyperfine splitting $\frac{K}{M_q^2} = 50$ MeV. One can immediately see that if we substitute the light quark mass in the hyperfine term with the heavier strange quark mass, the splitting gets smaller and the baryon decuplet members get lighter. A hexaquark system has a lot more permutations between quarks, which is why the splitting term is essential. Using this ansatz we have combined the evaluation of the $d^*(2380)$ multiplet masses in table 4.

For the entries in the table 4, note that $M_n = 363$ MeV is the mass of a light constituent quark, $M_s = 538$ MeV is the mass of a strange constituent quark, $\frac{K}{M_n^2} = 50$ MeV is the splitting parameter and $B_h \sim -550$ MeV is the hexaquark binding which reproduces the observed $d^*(2380)$ mass. This latter parameter is taken to be the same for all members of the multiplet (similar to a bag constant of [47]).

Note that under these assumptions the mass of the d_{sss} is 236 MeV lower than a $\Delta\Omega$ threshold and only 160 MeV higher than the $\Sigma\Xi$ threshold. This mass difference with the $\Sigma\Xi$

Table 5. Expected masses of the d^*(2380) SU(3) multiplet in pure molecular picture.

Particle	Binding energy structure	Mass value [MeV]	Binding[a] [MeV]
d^*	$M_{\rm Red}(\Delta\Delta)(3f\cdot 3f)^2$	2380	84
d_s	$M_{\rm Red}(\Delta\Sigma^*)(3f\cdot 2f)^2$	2578	39
d_{ss}	$2/3M_{\rm Red}(\Delta\Xi^*)(3f\cdot 1f)^2 \;+\; 1/3M_{\rm Red}(\Sigma\Sigma^*)(2f\cdot 2f)^2$	2753	13
d_{sss}	$0 + 1/2M_{\rm Red}(\Sigma^*\Xi^*)(2f\cdot 1f)^2$	2909	2

[a] Relative to the lightest member of the Decuplet–Decuplet state.

threshold is not large, especially considering the $d_{sss}\rightarrow\Sigma\Xi$ decay should proceed via a D-wave in the $\Sigma\Xi$-system[4].

Although this is a simplified picture of the antidecuplet, its flexibility does enable lower limits for multiplet masses to be evaluated. We note there are a large variety of estimates with different theoretical ansatz presented in [48–50].

4.1.2. d^*(2380) multiplet as molecules. For the case of a molecular state, the dependence of the mass of the states with strangeness content should be very different than for a genuine hexaquark. Due to the larger spatial extent of molecular states, one can assume that the binding is mainly driven by long-range pion exchange. In the SU(6) approximation, the pion couples only to light quarks, so the strength of the interaction Δ:Σ^*:Ξ^*:Ω should scale as 3:2:1:0, see [25]. Since the interaction potential (V) is proportional to the product of coupling constants and the binding energy $B\sim M_{\rm Red}\cdot V^2$, where $M_{\rm Red}$ is the reduced mass [51]. One can therefore estimate binding energies for various states under these assumptions. The $\Delta\Omega$ state is predicted have a binding energy $B=0$ MeV. However, since d_{sss} has not only $\Delta\Omega$ but also a $\Xi^*\Sigma^*$ component, we expect some binding also in this case. The results of these simplified calculations are summarised in table 5. Note that $M_{\rm Red}$ is the reduced mass and f is an effective meson baryon coupling constant fixed to reproduce the d^*(2380) mass. As was pointed out already in [52] one should expect large SU(3) breaking effects in this case. Indeed, if we compare expected binding for the d_{ss} from table 5, $B(d_{ss}\sim 13)$ MeV with the $M(\Delta\Xi^*)-M(\Sigma^*\Sigma^*)\sim 7$ MeV mass splitting, the expectations for a large SU(3) breaking become obvious. The mass difference between Σ^*+ and Σ^*- already give a 4 MeV split, so one should consider the numbers from a table 5 as a general guidance for the experimentalists, rather than exact calculations.

The observed mass of each member of the antidecuplet would be expected to lie between our estimations for a pure 'genuine hexaquark' and a pure 'molecule', figure 4. We note that group theory actually prohibits a 100% genuine hexaquark state, with the maximum purity state expected to contain 20% of molecular admixture [42, 43]. More elaborate recent calculations based on hidden colour quark models predict a 30% molecular contribution to the d^* [53], table 1. (within the group theory limits). The Resonating Group Method allows to incorporate both extremes and all intermediate states consistently, see e.g. [54]. It was recently applied to get the d_s binding energy [55], where the authors reported $B(d_s)\sim 52$ MeV, which is within our $B(d_s)\in[141, 39]$ MeV estimation range.

[4] Since $J^P(d_{sss})=3^+$ and both Σ and Ξ has $J^P=1/2^+$ quantum numbers, to conserve parity only even partial waves are allowed. A S-wave (L = 0) does not conserve total angular momentum. D- and G- waves (L = 2 or 4) allow to conserve both total angular momentum and parity. Due to the momentum dependence of various partial waves, D-wave is expected to be dominant, similar to $d^*\rightarrow pn$ decay [32, 33].

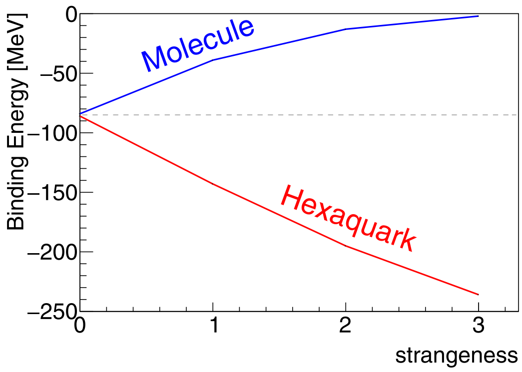

Figure 4. $d^{*}(2380)$ multiplet binding under assumption of pure genuine hexaquark (red) and pure molecule (blue).

In the following sections, we evaluate the widths and branching ratios as a function of binding energy. As a benchmark case, we take the same binding for all hexaquark members which, in this simplified ansatz, would correspond to nearly equal molecular/hexaquark contributions.

4.2. $d^{*}(2380)$ multiplet decays

Regardless of the inferred structure, the 10^{*} antidecuplet decay properties are largely driven by the SU(3) symmetry and available phase space for the decay products. Being a 10^{*} SU(3) antidecuplet, the decay couplings are limited to either octet+octet, $8 \oplus 8$ or decuplet +decuplet, $10 \oplus 10$ baryons. The possible decay branches, deriving from the established $8 \oplus 8$ and $10 \oplus 10$ baryonic members are summarised in table 6.

4.2.1. Formalism. In our calculations, we followed the PDG prescription on the parametrisation of resonances in a Breit–Wigner form [1] under the assumption that the width of all the resonances is energy-dependent and that the coupling constants for the decay into $8 \oplus 8$ are independent of the decaying particle's hypercharge/strangeness. This enables the total width of the decaying state to be expressed as

$$\Gamma_{tot} = \Gamma_8 + \Gamma_{10}, \tag{1}$$

with Γ_8 corresponding to the partial width for hexaquark decay into $8 \oplus 8$ and Γ_{10} for $10 \oplus 10$. In case of $8 \oplus 8$ decay, the final state particles are stable against strong decay so the Γ_8 decay width can be expressed as

$$\Gamma_8 = g_8^2 p^{2L+1} F_8(p) \tag{2}$$

with

$$F_8(p) = \frac{R^{2L}}{1 + R^{2L} p^{2L}}, \tag{3}$$

where g_8 is the coupling constant of the hexaquark decay into $8 \oplus 8$, p is the momentum of the ejectile in the hexaquark rest frame, L is the angular momentum in the system and $F(p)$ is a form factor, usually introduced to account for potential barriers with $R = 6.3\ (GeV/c)^{-1}$, similar to [44, 56, 57]. For a $J^P = 3^+$ particle decaying into two $J^P = 1/2^+$ baryons there are two possible scenarios—a $^3D_3$ partial wave with $L = 2$ and $^3G_3$ partial wave with $L = 4$. From

Table 6. Expected decay branches of the $d^*(2380)$ SU(3) multiplet.

Particle	$8 \oplus 8$	$10 \oplus 10$	Mass ($8 \oplus 8$) [MeV]	Mass ($10 \oplus 10$) [MeV]
d^*	pn	$\Delta\Delta$	p–n(1878)	$\Delta\Delta(2464)$
d_s	$\frac{1}{\sqrt{2}}(\mathrm{N}\Lambda - \mathrm{N}\Sigma)$	$\Delta\Sigma^*$	$N\Lambda(2054)$ $N\Sigma(2129)$	$\Delta\Sigma^*(2615)$
d_{ss}	$\frac{1}{\sqrt{6}}(\sqrt{2}\,\mathrm{N}\Xi + \Sigma\Sigma + \sqrt{3}\,\Sigma\Lambda)$	$\frac{-1}{\sqrt{3}}(\sqrt{2}\,\Delta\Xi^* + \Sigma^*\Sigma^*)$	$N\Xi(2253)$ $\Sigma\Sigma(2382)$ $\Sigma\Lambda(2499)$	$\Delta\Xi^*(2764)$ $\Sigma^*\Sigma^*(2767)$
d_{sss}	$\Sigma\Xi$	$\frac{1}{\sqrt{2}}(\Delta\Omega + \Sigma^*\Xi^*)$	$\Sigma\Xi(2504)$	$\Delta\Omega(2904)$ $\Sigma^*\Xi^*(2915)$

partial wave analysis of the $d^*(2380) \to pn$ reaction we know that the majority ($\sim$90%) of the $8 \oplus 8$ decay proceeds through the $^3D_3$ partial wave [32, 33, 58]. In our calculation, we will assume that all $8 \oplus 8$ decays proceed with $L = 2$. This assumption leads to minor corrections but allows a significant reduction in the required number of coupling constants. Since all the 3^+-hexaquarks lie far above the $8 \oplus 8$ threshold the Γ_8 width is expected to be nearly constant throughout the resonance. The g_8 constant is the same for all members of the antidecuplet. It is fixed to the value extracted from $\Gamma(d^* \to pn) = 8$ MeV.

The $10 \oplus 10$ channel is much more challenging since the baryon decuplet contains resonant states with a rather sizeable width and an associated strong energy dependence. The Γ_{10} width is expected to vary significantly between different resonances. We have calculated the Γ_{10} as follows

$$\Gamma_{10} = \gamma_{10} \int \mathrm{d}m_1^2 dm_2^2 F^2(q_{10}) |D_{D_1}(m_1^2) D_{D_2}(m_2^2)|^2 \tag{4}$$

$$F(q_{10}) = \frac{\Lambda^2}{\Lambda^2 + q_{10}^2/4}, \quad \Lambda = 0.16\ GeV/c \tag{5}$$

$$D_D = \frac{\sqrt{m_D \Gamma_D(q_M)/q_M}}{M_{BM}^2 - m_D^2 + i m_D \Gamma_D^{tot}(q_M)} \tag{6}$$

$$\Gamma_D = \gamma (q_M)^3 \frac{R^2}{1 + R^2 (q_M)^2} \tag{7}$$

here $F(q_{10})$ is a form factor which depends on the relative momentum (q_{10}) between the two decuplet baryons in the hexaquark decay. We follow the prescription of [56], and parameterise this dependence in a monopole form with a cut-off parameter Λ[5]. D_{D_1}/D_{D_2} are the propagators for the decuplet of baryons, with m_1/m_2 the baryon masses (or correspondingly the invariant mass of the meson-baryon system M_{BM} from the decuplet baryon decay into the octet of mesons and the octet of baryons) and m_D being their nominal Breit–Wigner masses. The energy-dependent width of the baryon decuplet decay, Γ_D, is parameterised in a standard form taking a P-wave decay resonance with Blatt–Weisskopf barrier factors of $R = 6.3$ GeV/c and $\gamma = 0.74$, 0.28, 0.13 for the Δ, Σ^* and Ξ^* respectively [59] [6]. Γ_D^{tot} is the sum of all partial widths. The width of the Ω is considered to be zero since it is stable with respect to strong decays. The form factor parameters are fixed based on the $d^* \to \Delta\Delta \to d\pi\pi$ invariant mass distributions of [27]. The γ_{10} is taken as a normalisation factor and is fixed to reproduce the width at zero binding energy, e.g $\Gamma(d^* \to \Delta\Delta) = 2 \cdot \Gamma_\Delta$ for $M_{d*} = 2 \cdot M_\Delta$.

$$\Gamma_{10}(B = 0) = \Gamma_{D_1} + \Gamma_{D_2} \tag{8}$$

4.2.2. d(2380) multiplet decays in a molecular picture.* If one would consider the d^* multiplet to be a purely molecular state one can expect some difference compared to the calculations above. The $10 \oplus 10$ decay would stay unchanged since it is a fall-apart decay and it is only driven by available phase space and the energy dependence of the widths. However,

[5] The appropriateness of the value for this cutoff parameter is explored in [27], where it was shown to provide agreement with measured invariant mass distributions in the region of the so-called ABC-effect.
[6] The $\gamma = 0.74$ for the Δ case can be taken directly from [59], while all the others can be recalculated from known width $\Gamma(\Sigma^*) = 36.2$ MeV, $\Gamma(\Xi^*) = 9.1$ MeV [1].

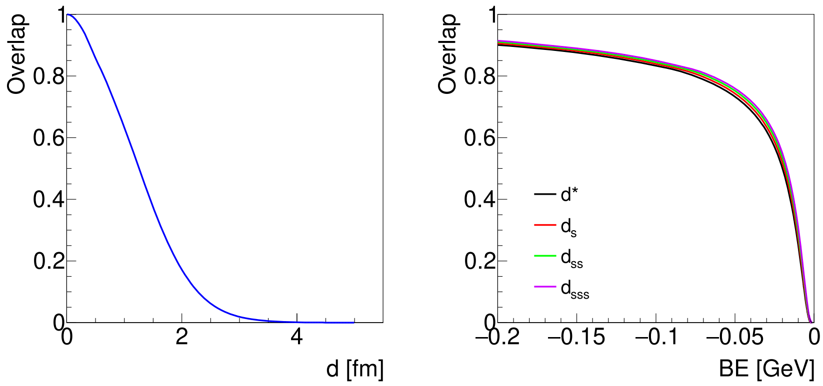

Figure 5. The density overlap for two diffuse balls in a molecule as a function of the distance (left) and binding energy (right).

to get a $8 \oplus 8$ decay, one needs to get a quark rearrangement, so this decay would be, besides other things, dependent on the wave function overlap. In the calculations above, the $8 \oplus 8$ decay width stayed essentially constant over the large range of binding energies, while in the molecular picture the wave function overlap at $B = 0$ should be zero, hence the $8 \oplus 8$ decay width should be also zero. To account for this effect we have modified equation (2) with

$$\Gamma_8 = \tilde{g}_8{}^2 p^{2L+1} F_8(p) \cdot P(B) \tag{9}$$

here $P(B)$ is a wave function overlap and $\tilde{g}_8$ is a modified g_8 constant to reproduce the $d^*(2380)$ width.

We have evaluated $P(B)$ in the following assumptions: (i) all particles assumed to be spherical with a distribution, similar to the proton charge distribution $\rho(r) = \exp(-a \cdot r)$, with a standard $a = 4.27\ \mathrm{fm}^{-1}$ (a coordinate space analogue of the proton dipole form factor with a cut-off parameter $\Lambda = 0.72\ \mathrm{GeV}/c$). The distance between molecular components was taken as

$$d = \frac{1}{\sqrt{2 \cdot M_{\mathrm{Red}}|B|}} \tag{10}$$

with M_{Red} being reduced mass and B is the binding energy. The wave function overlap $P(B)$ derived under these assumptions is shown on the figure below, figure 5. The $\tilde{g}_8$ is again fixed to reproduce the value extracted from $\Gamma(d^* \rightarrow pn) = 8\ \mathrm{MeV}$.

With these modified $8 \oplus 8$ decay width both the branching ratios and decay widths would change.

4.2.3. Results for the $d^*(2380)$ decuplet under assumption of a genuine hexaquark. We first explored the validity of the adopted cut-off parameter $\Lambda = 0.16\ \mathrm{GeV}/c$ in the model. The form factor of the form equation (5). with a cut-off parameter $\Lambda = 0.16\ \mathrm{GeV}/c$ in the $d^* \rightarrow \Delta\Delta \rightarrow d\pi\pi$ was first introduced in a [27] to explain the so-called ABC-effect, an enhancement in $M_{\pi\pi}$ close to the threshold. Indeed, for the case where the nucleons in the deuteron have a very small relative momentum, there is a correspondence in the relative momentum between the Δ's and the relative momentum between the pions (and hence with the pion invariant mass). It was later speculated by A. Gal, that reduction of the $\Delta\Delta$ system size within the d^* can lead to a further reduction of the d^* width [59]. To clarify the situation we have studied the predicted $d^* \rightarrow \Delta\Delta$ width dependence as a function of the cut of

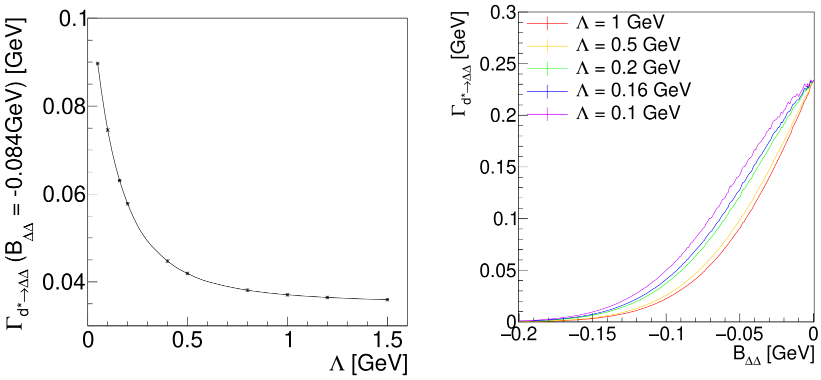

Figure 6. $d^*(2380)$ width as a function of the Form Factor cut-off parameter Λ (left) and partial width $\Gamma_{d*\to\Delta\Delta}$ as a function of binding energy (right) for various Λ values, $\Lambda = 0.1, 0.16, 0.2, 0.5, 1\ \mathrm{GeV}/c$ in rainbow order from violet to red.

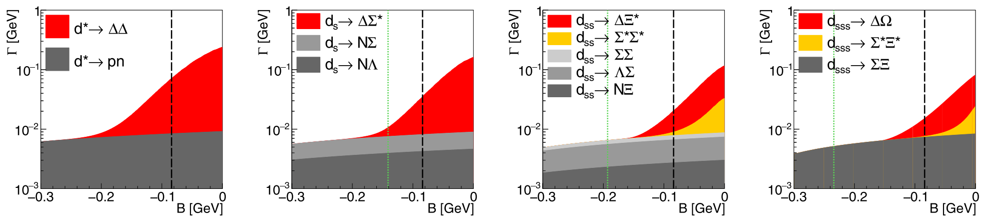

Figure 7. d^* multiplet total width as a function of binding energy (relative to the lightest member of the Decuplet–Decuplet pole) for the d^*, d_s, d_{ss}, d_{sss} from left to right split into major decay branches (note the log scale). The vertical line (common to all figures) shows the nominal expected mass, obtained under the assumption of the same binding (84 MeV) for all multiplet members(black dashed) and for specific binding (green dotted) as specified in table 4.

parameter Λ, figure 6. The results indicate that the adopted $\Lambda = 0.16\ \mathrm{GeV}/c$ cannot only reproduce the ABC effect but also gives agreement with the measured $\Gamma(d^* \to \Delta\Delta) = 62$ MeV. As evident in figure 6 we also reproduce the trend in which a higher value adopted for the cut-off parameter leads to a smaller width.

We, therefore, adopt $\Lambda = 0.16\ \mathrm{GeV}/c$ for subsequent calculations. The predicted width for all decuplet members is shown in figure 7. In figure 8 we show the predicted branching ratios as a function of binding energy. The results are summarised in tables 7, 8.

One can clearly see that the $8 \oplus 8$ decays are predicted to be increasingly important for the higher strangeness states, while both partial and total widths reduce substantially. For all hexaquark members only $10 \oplus 10$ with Δ in a final state are important. For the d_{ss} non-Δ channels ($\Sigma^*\Sigma^* \to$ anything) covers only 6%. For the d_{sss} the non-Δ ($\Xi^*\Sigma^*$) decay has a 2% branching ratio. It is interesting to note that the semi-electromagnetic decay branch $d_{sss} \to \Omega\Delta \to \Omega N\gamma$ is predicted to have a higher probability than the purely hadronic $d_{sss} \to \Xi^*\Sigma^* \to \Xi\Sigma\pi\pi$ decay.

Hexaquarks are expected to be produced copiously in heavy ion collisions, however, our estimations indicate the width for all d^* multiplet states is rather large. Unfortunately, this indicates their clean identification in the tough background conditions in typical heavy ion collisions may be challenging. The most feasible channel for such studies, having both a large

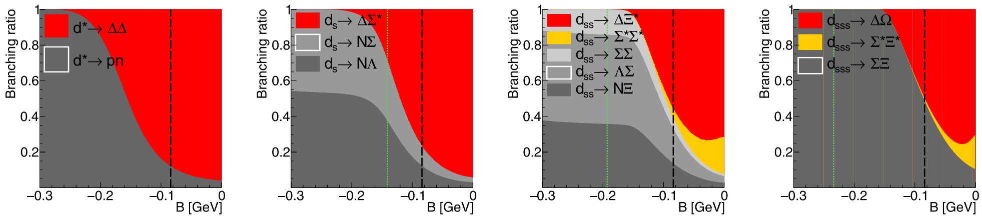

Figure 8. d^* multiplet branching ratio as a function of binding energy (relative to Decuplet–Decuplet pole) for the d^*, d_s, d_{ss}, d_{sss} from left to right. The vertical dashed line (common to all figures) shows the nominal expected mass, obtained under the assumption of the same binding (84 MeV) for all multiplet members (black dashed) and for specific binding (green dotted) as specified in table 4.

partial width and a convenient isospin $3/2$ separation, is the $\Xi\Sigma$ branch in d_{sss} decay. This could potentially be tested using $\Xi - \Sigma$ correlation functions.

For all channels the most prominent final states have high particle multiplicity, indicating hermetic detector apparatus and exclusivity conditions would need to be established. The d_{sss} is the only member with a dominant octet-octet decay, however, the necessity of associated kaon production to conserve strangeness in the production mechanism also provides challenges. The requirements may be more easily achieved using beams containing intrinsic strangeness (e.g. Kaon beams as proposed in [22]).

4.2.4. Results for d(2380) decuplet in a molecular picture.* The strong variation of the $8 \oplus 8$ decay in the molecular picture leads to very different behaviour in total width and branching ratios, see table 9, figures 9, 10. One also should keep in mind that for this case we generally would expect much smaller binding energies for all strange members than $B = -84$ MeV of a $d^*(2380)$. While at $B = -84$ MeV binding the decay branchings are nearly identical for the two scenarios, in a purely molecular picture with reduced binding the importance of the $8 \oplus 8$ decays is also reduced. For the d_{sss} case with a tiny ($B = 2$ MeV) binding the $8 \oplus 8$, $\Sigma\Xi$ decay is $\Gamma(\Sigma\Xi) = 1.5$ MeV, or only 2% of the total decay width. It is also interesting to note that in a molecular picture, the state with zero strangeness, $d^*(2380)$ should have the smallest width. One can see that even the stability of the Ω baryon for the case of d_{sss} state cannot compensate the massive reduction of the binding energy, which leads to the prediction of a much larger width for d_{sss} compared to the $d^*(2380)$.

5. ΔΔ 28-plet

The $\Delta\Delta$ 28-plet(figure 11) couples only to a $10 \oplus 10$, hence it is extremely difficult to access it even in a low strangeness state, see table 10. Indeed, the zero strangeness state has isospin equal to 3. Due to isospin conservation, it does not couple to NN, hence cannot be directly accessed in nucleon-nucleon collisions. It does not couple to NNπ, hence cannot be probed with pion beams on nuclei. (One can still use pion double-charge exchange reactions on nuclei, e.g. $\pi^+ A \to \pi^- X$, but not with direct S-channel production).

The simplest way to produce even the lowest lying $\Delta\Delta$ state is a nucleon-nucleon collision with two associated pions, $pp \to (\pi\pi)(\Delta\Delta)$. A six-body final state and rather high energy required for this reaction make it extremely challenging to predict the background from conventional reactions, which can also potentially produce structures in the measured cross sections. As a result bump-hunting techniques become extremely unreliable. If the

Table 7. d^* multiplet width results for the $B = -84$ MeV.

$d^*(2380)$			$d_s(2531)$			$d_{ss}(2670)$			$d_{sss}(2820)$		
decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]
$\Delta\Delta$	62.3	88	$\Delta\Sigma^*$	26.9	77	$\Delta\Xi^*$	11.3	56	$\Delta\Omega$	7.6	50
						$\Sigma^*\Sigma^*$	1.1	6	$\Sigma^*\Xi^*$	0.2	2
pn	8.4	12	$\mathrm{N}\Lambda$	4.2	12	$\mathrm{N}\Xi$	2.7	13	$\Sigma\Xi$	7.4	48
			$\mathrm{N}\Sigma$	3.9	11	$\Lambda\Sigma$	3.9	19			
						$\Sigma\Sigma$	1.2	6			
total	**70.7**		total	**35.0**		total	**20.2**		total	**15.2**	

Table 8. d^* multiplet width results for the expected binding energies from table 4.

$d^*(2380)$, $B = 84$ MeV			$d_s(2474)$, $B = 141$ MeV			$d_{ss}(2571)$, $B = 193$ MeV			$d_{sss}(2670)$, $B = 234$ MeV		
decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]
$\Delta\Delta$	62.3	88	$\Delta\Sigma^*$	3.1	29	$\Delta\Xi^*$	<0.1	<1	$\Delta\Omega$	0	0
						$\Sigma^*\Sigma^*$	<0.1	<1	$\Sigma^*\Xi^*$	0	0
pn	8.4	12	NΛ	3.9	37	NΞ	2.3	36	$\Sigma\Xi$	5.1	100
			NΣ	3.6	34	$\Lambda\Sigma$	3.3	50			
						$\Sigma\Sigma$	0.9	14			
total	**70.7**		total	**10.6**		total	**6.5**		total	**5.1**	

Table 9. d^* multiplet width results for the expected binding energies from table 5. (molecule).

$d^*(2380)$			$d_s(2576)$, B = 39 MeV			$d_{ss}(2751)$,B = 13 MeV			$d_{sss}(2902)$, B = 2 MeV		
decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]
ΔΔ	62.3	88	$\Delta\Sigma^*$	82.9	91	$\Delta\Xi^*$	69.9	74	ΔΩ	56.3	78
						$\Sigma^*\Sigma^*$	17.3	18	$\Sigma^*\Xi^*$	14.4	20
p–n	8.4	12	NΛ	4.2	5	NΞ	2.3	3	ΣΞ	1.5	2
			NΣ	3.9	4	ΛΣ	3.3	4			
						ΣΣ	1.0	1			
total	**70.7**		total	**91.0**		total	**94.0**		total	**72.2**	

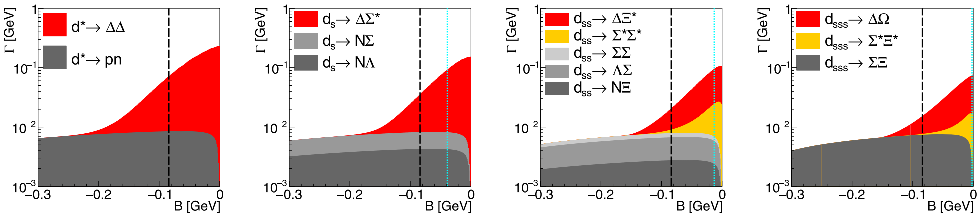

Figure 9. Same as figure 7, but for a molecular picture with nominal molecular masses shown by the cyan dotted line as in table 5.

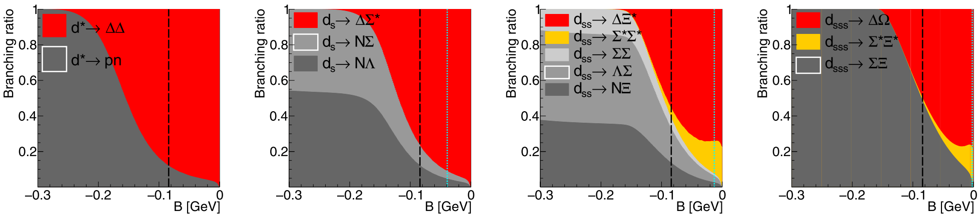

Figure 10. same as figure 8, but for a molecular picture with nominal molecular masses shown by cyan dotted line as in table 5.

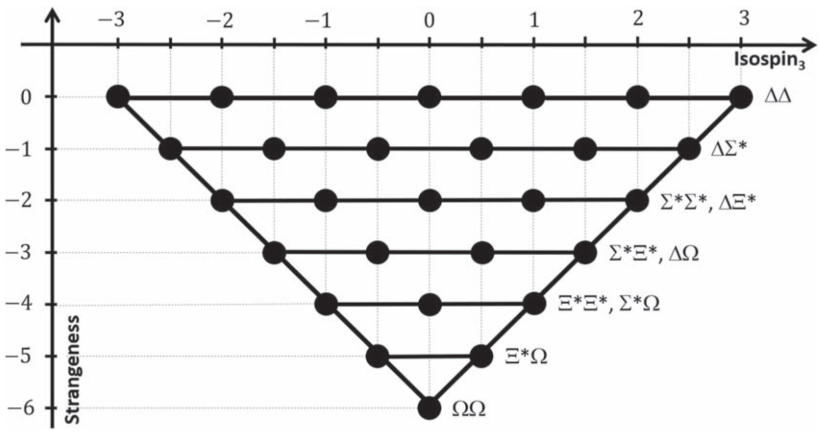

Figure 11. $\Delta\Delta$ 28-plet $10 \oplus 10$.

$\Delta\Delta$ 28-plet would be strongly bound ($\sim$80 MeV or higher, as the d^* antidecuplet) it members would have an even smaller width than corresponding d^* multiplet particles, since the $8 \oplus 8$ decay is not allowed here. Unfortunately, in the only measurement from [60] , where the authors tried to extract the maximum isospin projection state, $pp \to \pi^-\pi^-\Delta^{++}\Delta^{++}$, it was demonstrated that this multiplet is either loosely bound or loosely unbound. That unavoidably means that the widths of all states would be similar to the combined width of constituent particles, e.g. in $\Delta\Delta$ case it would be $\Gamma \sim 240$ MeV. Under such conditions, the only state which can be experimentally accessed is the strong-decay-free $\Omega\Omega$ dibaryon. There is a hope to get the $\Omega\Omega$ scattering length from heavy ion collisions. Lattice calculations also support the

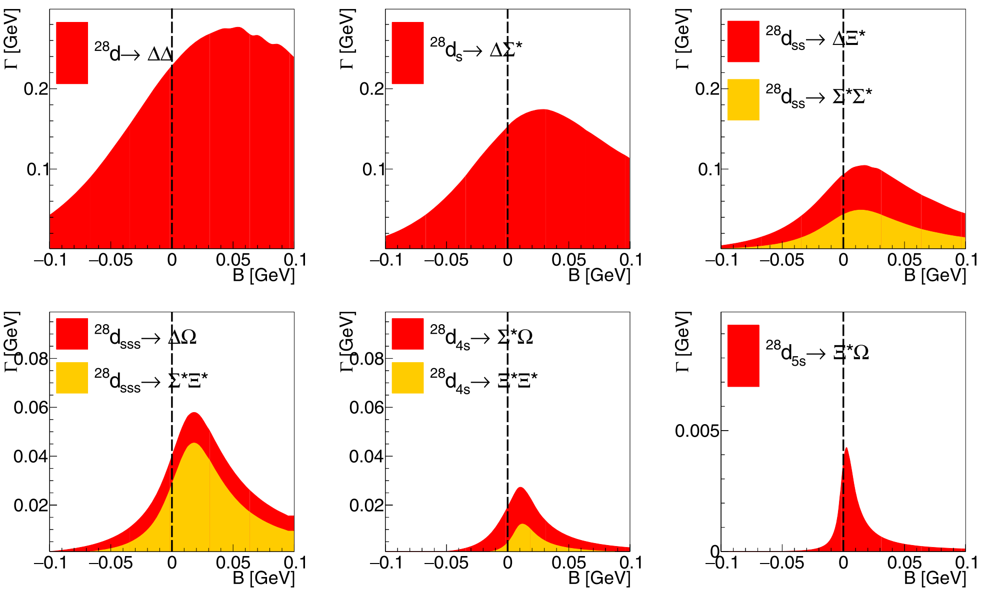

Figure 12. 28-plet total width as a function of binding energy (relative to the Decuplet–Decuplet pole) for the various states with increased strangeness from top left ($S = 0$) to bottom right ($S = -5$) split into major decay branches.

Table 10. Expected decay branches of the spin $J = 0$ SU(3) 28-plet.

Strangeness	$10 \oplus 10$	Mass [MeV]
0	$\Delta\Delta$	$\Delta\Delta(2464)$
-1	$\Delta\Sigma^*$	$\Delta\Sigma^*(2615)$
-2	$\frac{1}{\sqrt{5}}(\sqrt{2}\,\Delta\Xi^* + \sqrt{3}\,\Sigma^*\Sigma^*)$	$\Sigma^*\Sigma^*(2766)$ $\Delta\Xi^*(2764)$
-3	$\frac{1}{\sqrt{10}}(\Delta\Omega + 3\Sigma^*\Xi^*)$	$\Sigma^*\Xi^*(2915)$ $\Delta\Omega(2904)$
-4	$\frac{1}{\sqrt{5}}(\sqrt{2}\,\Sigma^*\Omega + \sqrt{3}\,\Xi^*\Xi^*)$	$\Xi^*\Xi^*(3064)$ $\Sigma^*\Omega(3055)$
-5	$\Xi^*\Omega$	$\Xi^*\Omega(3204)$
-6	$\Omega\Omega$	$\Omega\Omega(3344)$

expectation of loosely bound 28-plet. For example, [61] predict $\Omega\Omega$ state to be either bound by less than an MeV or unbound.

As a first approximation, the 28-plet decay widths can be evaluated by taking the $10 \oplus 10$ decay width from the d^* antidecuplet. The only difference is the binding energy dependence of the width for a 28-plet compared to an antidecuplet would be the size of a cut-off parameter Λ. As one can see in figure 6, for a given binding energy the width gets larger with an increase of the cut-off and for a small binding energy, close to $B = 0$ MeV the effect of cut-off is minimal due to the normalisation condition. Since we expect the 28-plet Λ to be larger than for the antidecuplet, the use of $\Lambda_{\overline{10}} \sim 0.16$ GeV give us an upper limit for an expected width, figure 12. The $\Omega\Omega$ state is stable against hadronic decays if bound so is therefore not shown in figure 12.

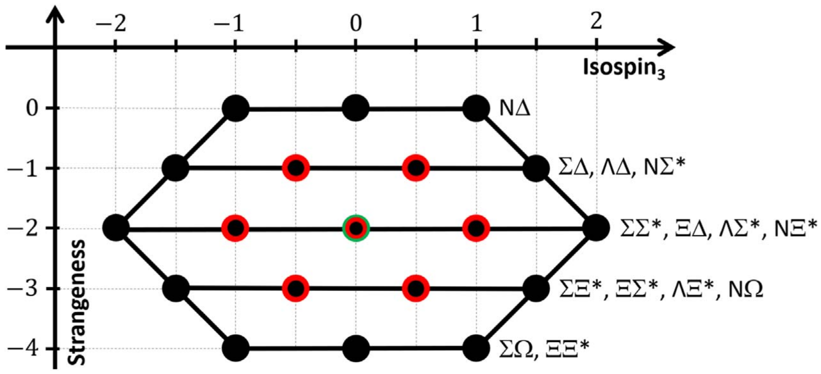

Figure 13. NΔ 27-plet $8 \oplus 10$.

6. NΔ 27-plet

This multiplet (figure 13) is very interesting. The zero strangeness NΔ state is known since the '60s and is often referred to as an R Arndt resonance [62, 63]. The best way to observe it is a $\pi d \to pp$ or $pp \to \pi d$ reaction, where it reveals itself as a perfect loop on the Argand plot in a $^1D_2$ partial wave. There are plenty of data collected on this resonance experimentally—including differential cross-sections, single- and double-polarisation observables [63]. There are even some data on deuteron tensor polarisation for this reaction [64]. While experimental evidence for this state is settled, the theoretical interpretation of these data is challenging. We have a very long list of proposed explanations, including NΔ-FSI, extra attraction due to ρ-meson exchange, the effect of box diagrams and it being a genuine dibaryon state. The latest analysis of Hoshizaki tends to explain it with a dibaryon state [65, 66]. Also, modern Fadeev calculations from Gal and Garcilazo tend to explain this state with NΔ molecule [67]. Very recently Niskanen put a paper explaining this state by scattering dynamics in presence of attractive $N-\Delta$ potential [68]. So after 50 years of research the field is not settled. It is also unclear if its various descriptions derive from semantics due to the different theoretical approaches or if we are really discussing different phenomena. Regardless, under the assumption of a bound state with 20 MeV binding energy, it is hard to expect anything other than a loosely bound molecules, inline with the prediction of [67]. We therefore take this as an assumption in our theoretical analysis.

There are two important questions to be addressed—what are the expected binding energies for the other states and what are their branching ratios? Flavour configurations for all members of multiplet for both $8 \oplus 10$ and $8 \oplus 8$ decays together with their nominal masses are shown in tables 11, 12, 13. Making parallels to the deuteron antidecuplet, where the first state (deuteron) is bound by 2.2 MeV and the second (NΛ) is unbound by 166 keV only, we can expect for the 27-plet with the first state bound by 20 MeV that higher strangeness bound states might also exist. Since the mass splitting is comparable or even larger than expected binding energies, the use of SU(6) formalism from section 4.1.2 is not applicable. However, we can use similar arguments to make qualitative evaluations. One can say that the binding should be proportional to the number of light quarks squared, so if a state with all 6 light quarks has a binding energy of 20 MeV, the state with 3 light and 3 strange quarks should have a binding energy of ~1.3 MeV (table 14). The latter state was recently explored both theoretically and experimentally with the observation of $p\Omega$ state in heavy ion collisions [70] and in lattice QCD calculations [71], which predicted it to be bound by 2.5 MeV (1.5 MeV

Table 11. Expected $8 \oplus 10$ decay branches of the spin $J = 2$ 27-plet.

S	Max Isospin	Med Isospin	Min Isospin
0	$\mathrm{N}\Delta$		
−1	$\frac{1}{4}(\sqrt{5}\,\Sigma\Delta + 3\Lambda\Delta - \sqrt{2}\,\mathrm{N}\Sigma^*)$	$\frac{1}{\sqrt{5}}(2\mathrm{N}\Sigma^* - \Sigma\Delta)$	
−2	$\frac{1}{2}(\sqrt{3}\,\Xi\Delta - \Sigma\Sigma^*)$	$\frac{\sqrt{5}}{10}(3\Sigma\Sigma^* + \sqrt{6}\,\Lambda\Sigma^* - 2\mathrm{N}\Xi^* - \Xi\Delta)$	$\frac{1}{\sqrt{5}}(\sqrt{3}\,\mathrm{N}\Xi^* - \sqrt{2}\,\Sigma\Sigma^*)$
−3	$\frac{1}{\sqrt{2}}(\Xi\Sigma^* - \Sigma\Xi^*)$	$\frac{\sqrt{5}}{20}(7\Sigma\Xi^* + 3\Lambda\Xi^* - 3\sqrt{2}\,\mathrm{N}\Omega - 2\Xi\Sigma^*)$	
−4	$\frac{1}{2}(\sqrt{3}\,\Sigma\Omega - \Xi\Xi^*)$		

Table 12. Expected $8 \oplus 8$ decay branches of the spin $J = 2$ 27-plet.

S	Max Isospin	Med Isospin	Min Isospin
0	NN		
−1	$\mathrm{N}\Sigma$	$\frac{1}{\sqrt{10}}(3\mathrm{N}\Lambda\ - \mathrm{N}\Sigma)$	
−2	$\Sigma\Sigma$	$\frac{1}{\sqrt{10}}(2\mathrm{N}\Xi + \sqrt{6}\,\Sigma\Lambda)$	$\frac{1}{\sqrt{10}}(\frac{3\sqrt{3}}{2}\Lambda\Lambda - \sqrt{3}\,\mathrm{N}\Xi - \frac{1}{2}\Sigma\Sigma)$
−3	$\Sigma\Xi$	$\frac{1}{\sqrt{10}}(3\Lambda\Xi - \Sigma\Xi)$	
−4	$\Xi\Xi$		

Table 13. Masses of $8 \oplus 10$ and $8 \oplus 8$ decay branches of the spin $J = 2$ 27-plet.

S	$8 \oplus 10$ Mass [MeV]	$8 \oplus 8$ Mass [MeV]
0	$\mathrm{N}\Delta(2170)$	NN(1876)
−1	$\Sigma\Delta(2421)\ \Lambda\Delta(2348)\ \mathrm{N}\Sigma^{*(2321)}$	$\mathrm{N}\Lambda(2054)\ \mathrm{N}\Sigma(2127)$
−2	$\Sigma\Sigma^{*(2572)\ \Xi\Delta(2547)\ \Lambda\Sigma^*(2499)\ \mathrm{N}\Xi^*(2470)}$	$\mathrm{N}\Xi(2253)\ \Lambda\Lambda(2232)\ \Sigma\Sigma(2378)$ $\Sigma\Lambda(2305)$
−3	$\Sigma\Xi^{*(2721)\ \Xi\Sigma^*(2698)\ \Lambda\Xi^*(2648)\ \mathrm{N}\Omega(2610)}$	$\Lambda\Xi(2431)\Sigma\Xi(2504)$
−4	$\Sigma\Omega(2861)\ \Xi\Xi^*(2847)$	$\Xi\Xi(2630)$

Table 14. Expected masses of the spin $J = 2$ d-27 SU(3) multiplet in pure molecular picture.

Particle	Binding energy structure	Mass value [MeV]	Binding[a] [MeV]
^{27}d	$M_{\mathrm{Red}}(\mathrm{N}\Delta)(3f \cdot 3f)^2$	2151	20.0
$^{27}d_s$	$[\frac{2}{16}M_{\mathrm{Red}}(\mathrm{N}\Sigma^*) + \frac{9}{16}M_{\mathrm{Red}}(\Delta\Lambda) + \frac{5}{16}M_{\mathrm{Red}}(\Delta\Sigma)](3f \cdot 2f)^2$	2311	10.0
$^{27}d_{ss}$	$3/4M_{\mathrm{Red}}(\Xi\Delta)(3f \cdot 1f)^2 + 1/4M_{\mathrm{Red}}(\Sigma\Sigma^*)(2f \cdot 2f)^2$	2551	3.2
$^{27}d_{sss}$	$1/2[M_{\mathrm{Red}}(\Sigma\Xi^*) + M_{\mathrm{Red}}(\Xi\Sigma^*)](2f \cdot 1f)^2$	2704	1.3
$^{27}d_{4s}$	$3/4M_{\mathrm{Red}}(\Sigma\Omega)(2f \cdot 0f)^2 + 1/4M_{\mathrm{Red}}(\Xi\Xi^*)(1f \cdot 1f)^2$	2853	0.3

[a] Relative to the lightest Octet-Decuplet pole.

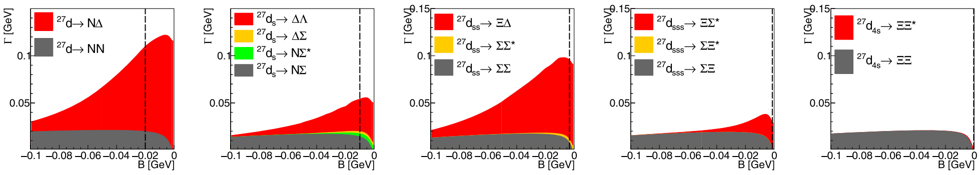

Figure 14. *d*-27 multiplet total width as a function of binding energy (relative to the lightest member of the Octet-Decuplet pole) for the states with various strangeness increasing from left to right split into major decay branches. The vertical dashed line shows the expected mass as specified in table 14.

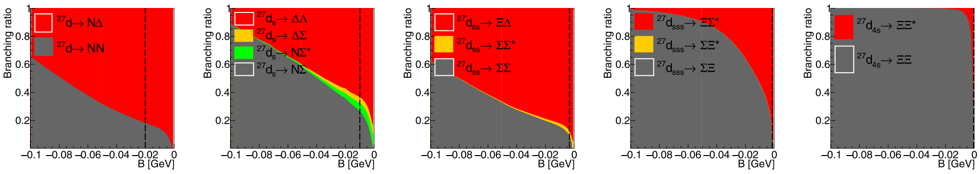

Figure 15. Same as figure 14, but for branching ratios.

strong binding plus 1.0 MeV electromagnetic binding). Hence we can expect a single strangeness state to be bound by about 10 MeV and a double strangeness state by 3.2 MeV[7].

In all states with a Δ component, the width is largely dominated by the width of the Δ and resulting states appeared to have similar widths e.g. $\Gamma(^{27}d) = 111$ MeV and $\Gamma(^{27}d_{ss}) = 96$ MeV. All of these states show interesting, and rather counterintuitive behaviour. An increase of binding energy first leads to an increase of the width. However, the increase of the binding leads to a decrease of particle separation in the state and therefore an increase of the wave function overlap. This unavoidably means an increase of the phase Space available for the octet-octet decay branch. That is why all members of this multiplet have wider widths than free states.

A $^{27}d_s$ state shows unusually small width of $\Gamma(^{27}d_s) \sim 40$ MeV. The lightest single strangeness component, $N\Sigma^*$ has a mass which is nearly 30 MeV lighter than the next lightest $\Lambda\Delta$ branch. An additional 10 MeV $^{27}d_s$ binding means that this state is bound by 40 MeV relative to $\Lambda\Delta$ pole. Such large binding leads to a sizeable reduction of a Δ width. On top of which, the $\Lambda\Delta$ branch is also suppressed by Clebsh–Gordan coefficients (see Table 11), making $^{27}d_s$ sufficiently narrow, but still wider than a free Σ^* state. The experimental search of this state might be challenging. One could search for the $\Lambda\Delta$ branch, however the necessity of a partial wave analysis of a 4 body reaction(Λ, N, π + an associated particle) makes this channel challenging from a theoretical point of view. An $8 \oplus 8$ decay, where the state can reveal itself as a Flatte [1] shape with sub-threshold (relative to the $8 \oplus 10$ pole) peaking or as a cusp-like behaviour looks more appealing theoretically, but very inconvenient from an experimental point of view due to necessity to measure Σ in the final state.

Strangeness -2 state with an isospin $I = 2$ looks very unfavourable due to the large $\Xi\Delta$ branch as can be seen in figures 14, 15. The isospin $I = 1$ state also has a $\Xi\Delta$ component, somewhat suppressed by the CG coefficients. Only the isospin $I = 0$ state does not have Δ components. This state also has rather convenient $8 \oplus 8$ decays, $N\Xi$ and $\Lambda\Lambda$ which can be explored experimentally. As discussed earlier for another 27-plet case one may expect some mixing between states. Due to this uncertainty, we do not show any width/branching calculations for the inner multiplet states, where SU(3) mixing can occur, table 15.

[7] the strangeness 4 state is likely to be unbound.

Table 15. ^{27}d multiplet width results for the maximum isospin states and expected binding energies from table 14.

$^{27}d(2151)$, $B = 20$ MeV			$^{27}d_s(2311)$, $B = 10$ MeV			$d_{ss}(2551)$, $B = 3$ MeV			$d_{sss}(2704)$, $B = 1.3$ MeV			$d_{4s}(2853)$, $B = 0.3$ MeV		
decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]	decay	Γ, [MeV]	BR [%]
NΔ	91	82	ΛΔ	34.4	64	ΞΔ	84.5	89	$\Xi\Sigma^*$	27.1	91	$\Xi\Xi^*$	2.2	81
			ΣΔ	2.2	4	$\Sigma\Sigma^*$	3.3	3	$\Sigma\Xi^*$	0.2	1	ΣΩ	0	0
			$N\Sigma^*$	3.2	6									
NN	20	18	NΣ	14.2	26	ΣΣ	7.7	8	ΣΞ	2.3	8	ΞΞ	0.5	19
total	**111**		total	**54.0**		total	**95.5**		total	**29.6**		total	**2.7**	

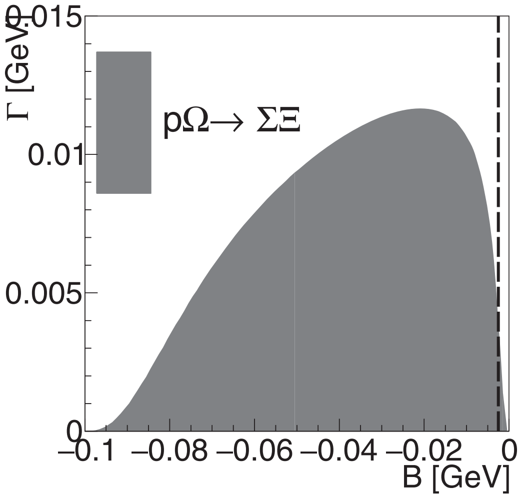

Figure 16. $p\Omega \rightarrow \Sigma\Xi$ decay width.

It is very unlikely that strangeness -3 isospin $I = 3/2$ state will be observed. A tiny binding of about 1 MeV would be really difficult to see in a 30 MeV wide state. The inner $I = 1/2$ state is a lot more interesting. We do not know the composition of that state due to unknown isospin mixing and enormous symmetry breaking due to the mass difference, however, we do know that the $p\Omega$ component is by far the lightest one among all ($M(p\Omega) = 2611$ MeV, $M(\Lambda\Xi^*) = 2648$ MeV, $M(\Xi\Sigma^*) = 2705$ MeV, $M(\Sigma\Xi^*) = 2721$ MeV). Both lattice QCD [71] and the early heavy ion correlation function analysis [70] claimed the existence of this state with a binding energy of about 2.5 MeV, where half of the binding originates from the electromagnetic attraction. So this state, if it exists, is a mixture of an atomic and a hadronic molecule. The latest heavy ion results with increased statistics tend to suggest this state is unbound. If bound, this state is stable against $8 \oplus 10$ decays, since Ω is stable against strong decays, however, it is still allowed to decay into $8 \oplus 8$ via the most prominent $\Lambda\Xi$ channel due to both phase space and the large, $9/10$, associated CG coefficient. For a $\sim$2.5 MeV binding energy one can expect a measurable $\Gamma(p\Omega \rightarrow \Sigma\Xi) \sim 4$ MeV width, see figure 16. This coupled channel effect also needs to be taken into account in the analysis of heavy ion correlation data.

The strangeness -4 state is very interesting. According to CG coefficients, it is mainly made of $\Sigma\Omega$ with $\Xi\Xi^*$ occupying only a quarter of the wavefunction. However, the mass of a $\Sigma\Omega$ configuration is nearly 10 MeV higher than the one of $\Xi\Xi^*$. So, if it exists, this state will be bound relative to the $\Xi\Xi^*$ threshold and with a $\Sigma\Omega$ branch being exactly zero. Also due to the very small binding, and hence the large size of this molecule, the wave function overlap will be tiny. That is why the $8 \oplus 8$ decay is strongly suppressed. All these factors would lead to an extremely small width of this state, provided it is bound. It should also be mentioned that an atomic $\Sigma^+\Omega^-$ state could be possible. It probably will be bound by about 1 MeV and have a width in the order of MeV with two possible decay branches $\Xi\Xi^*$ and $\Xi\Xi$. Both branches require quark rearrangements hence both will be small. The $\Xi\Xi$ has large phase space, but it requires additional spin-flip and angular momentum $L = 2$. The $\Xi\Xi^*$ happens in $S-$wave, does not require spin flips, but has tiny available phase space $<$10 MeV. One can search this state in $\Xi\Xi$ pair correlations in heavy ion collisions. The study of this atomic state

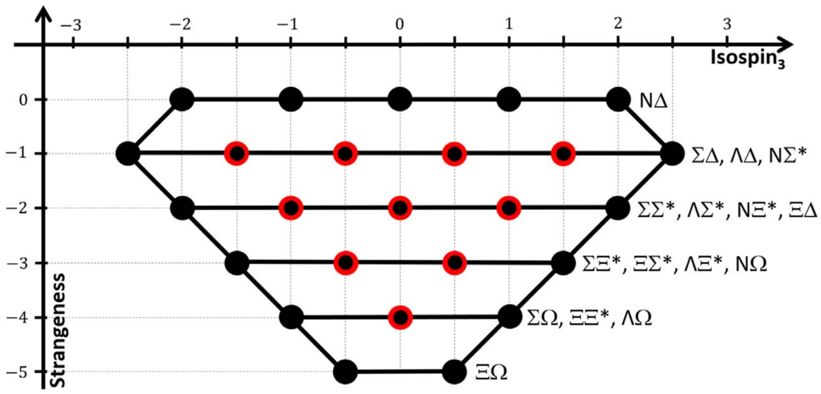

Figure 17. NΔ 35-plet $8 \oplus 10$.

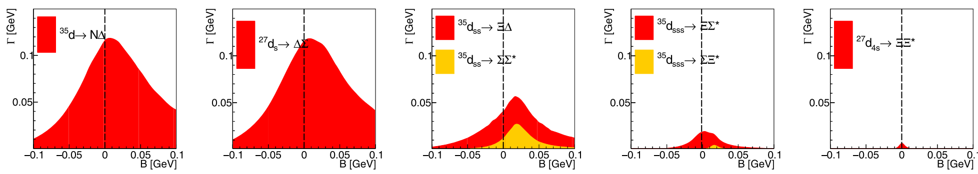

Figure 18. d-35 multiplet total width as a function of binding energy (relative to the lightest member of the Octet-Decuplet pole) for the states with various strangeness increasing from left to right split into major decay branches. The vertical dashed line shows zero.

can give a lot of information about hadronic $\Sigma^+\Omega^-$ interactions which would appear as a small addition to the larger electromagnetic binding.

7. NΔ 35-plet

In general, the analogy between NΔ 27- and 35-plets (figure 17) should be similar to that between the $10^* -$ and $27 -$ NN multiplets. Since the 27-plet is less bound compared to 10^*-, the 35-plet should be less bound compare to NΔ 27-plet. Accessing the members of these multiplet is challenging—they all have large isospin, they do not couple to $8 \oplus 8$, so production of any of these states require associated particles and all decays are fall-apart many body decays. By analogy to NN-multiplets, one can even expect that all states of this multiplet should be unbound (deuteron is bound, while nn/pp-states of 'Demon deuteron' are not). However, recently it was found experimentaly that a spin $S = 1$, isospin $I = 2$ NΔ state appears to bound by nearly 20 MeV—only slightly smaller compared to the NΔ state of a 27-plet [69]. If true, it implies that the other states of this multiplet may also be bound. That is why we decided to calculate possible decay width of these states as well, see figure 18, in a similar fashion as a 27-plet. The flavour configurations for these states are presented in table 16.

In all cases we expect the widths of these states to be close to the nominal width of their constituents, with not much chance for an experiment to reliably detect any of these states. The two exceptions are ΛΩ and ΞΩ, which would be stable against strong decays, if bound. These states can be potentially accessed via heavy-ion correlation functions studies. However, accumulating sufficient statistics for the strangeness −4 or −5 states may be extremely challenging.

Table 16. Expected decay branches of the spin $J = 1$ SU(3) 35-plet.

Strangeness	Max Isospin	Min Isospin	Mass [MeV]
0	$N\Delta$		$N\Delta(2170)$
-1	$\Sigma\Delta$	$\frac{1}{4}(\sqrt{10}N\Sigma^* - \Sigma\Delta + \sqrt{5}\Lambda\Delta)$	$\Sigma\Delta(2421)$ $\Lambda\Delta(2348)$ $N\Sigma^*(2321)$
-2	$\frac{1}{2}(\sqrt{3}\Sigma\Sigma^* + \Xi\Delta)$	$\frac{1}{2\sqrt{3}}(2N\Xi^* - \Sigma\Sigma^* + \sqrt{6}\Lambda\Sigma^* - \Xi\Delta)$	$\Sigma\Sigma^*(2572)$ $\Xi\Delta(2547)$ $\Lambda\Sigma^*(2499)$ $N\Xi^*(2470)$
-3	$\frac{1}{\sqrt{2}}(\Sigma\Xi^* + \Xi\Sigma^*)$	$\frac{1}{4}(\sqrt{2}N\Omega - \Sigma\Xi^* + 3\Lambda\Xi^* - 2\Xi\Sigma^*)$	$\Sigma\Xi^*(2721)$ $\Xi\Sigma^*(2698)$ $\Lambda\Xi^*(2648)$ $N\Omega(2610)$
-4	$\frac{1}{2}(\Sigma\Omega + \sqrt{3}\Xi\Xi^*)$	$\frac{1}{\sqrt{2}}(\Lambda\Omega - \Xi\Xi^*)$	$\Sigma\Omega(2861)$ $\Xi\Xi^*(2847)$ $\Lambda\Omega(2788)$
-5	$\Xi\Omega$		$\Xi\Omega(2987)$

8. Summary

We have developed a theoretical model, employing experimentally constrained parameters, to predict the possible decay branches and partial widths for all members of the d^* hexaquark antidecuplet. For all strange-quark containing members of the antidecuplet, the predicted widths are rather large, with the most promising decay channels including the broad Δ resonance in the final state. We demonstrated that a d^* form factor, which was first introduced to explain peculiarities in the $d^* \to d\pi\pi$ decay, can also explain the smallness of the d^* width, in agreement with the qualitative dimensional arguments of A. Gal. The results of the paper will be an important guide for the ongoing search for the d^* anti-decuplet members employing photon-, pion-, and kaon-induced reactions, as well as in high-energy collider experiments. We have also extended our model to systematically study all other light and strange dibaryons which can appear as a ground state of $8 \oplus 8$, $8 \oplus 10$ and $10 \oplus 10$ baryon multiplets configurations. Several interesting states were identified for experimental searches and strategies in search of these states were proposed. From our evaluations the most promising channels in the d^* SU(3) multiplet appeared to be the d_s state with polarisation observables studies of the $N\Lambda$ branch and heavy ion collisions correlation analysis of the $N\Xi$ and $\Sigma\Xi$ channels in search for d_{ss} and d_{sss} states. From the other multiplets the most accessible dibaryons appear in $N\Delta$ 27-plet with the possible $N\Sigma$ cusp at the vicinity of $N\Sigma^*$ threshold and $N\Omega$ state studied with correlation analysis.

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Data availability statement

The data that support the findings of this study will be openly available following an embargo at the following URL/DOI: https://doi.org/10.15124/7e27c95e-63f5-43de-9cf7-d1aeb52f6d02. Data will be available from 1 January 2024.

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References

[1] Workman R L *et al* (Particle Data Group) 2022 *Prog. Theor. Exp. Phys.* **083C01** 2022
[2] Urey H and Murphy G 1932 *Phys. Rev.* **39** 164–5
[3] Dyson F J and Xuong N-H 1964 *Phys. Rev. Lett.* **13** 815
[4] Gell-Mann M 1964 *Phys. Lett.* **8** 214–5
[5] Clement H 2017 *Prog. Part. Nucl. Phys.* **93** 195
[6] de Swart J J 1963 *Rev. Mod. Phys.* **35** 916
[7] McNamee P S J and Chilton F 1964 *Rev. Mod. Phys.* **36** 1005
[8] Miller G A 1864 *Few-Body Syst.* **56** 319–24
[9] Haidenbauer J, Meiner U-G and Nogga A 2020 *Eur. Phys. J.* A **56** 91
[10] (ALICE collaboration) 2022 *Phys. Lett.* B **833** 137272
[11] Acharya S *et al* 2019 *Phys. Rev. Lett.* **123** 112002
[12] Migdal A B 1954 *Sov. Phys.* **1** 2
[13] Watson K M 1952 *Phys. Rev.* **88** 1163
[14] Perkins D 2001 *Introduction to High-Energy Physics* (Cambridge University press)
[15] Jaffe R 1977 *Phys. Rev. Lett.* **38** 617
[16] Takahashi H *et al* 2001 *Phys. Rev. Lett.* **87** 212502
[17] Ahn J K *et al* 2013 *Phys. Rev.* C **88** 014003
[18] Adamczyk L *et al* 2015 *Phys. Rev. Lett.* **114** 022301
[19] Acharya S *et al* 2019 *Phys. Lett.* B **797** 134822
[20] Sasaki K *et al* 2020 *Nucl. Phys.* A **998** 121737
[21] Rowley J *et al* 2021 *Phys. Rev. Lett.* **127** 272303
[22] Amaryan M *et al* 2008 arXiv:2008.08215
[23] (ALICE collaboration) 2023 *Phys. Lett.*B **844** 137223
[24] Haidenbauer J, Meiner U-G and Nogga A 2020 *Eur. Phys. J.* A **56** 91
[25] Mantziris A, Pastore A, Vidaña I, Watts D P, Bashkanov M and Romero A 2020 *Astron. Astrophys.* **638** A40
[26] Bashkanov M *et al* 2009 *Phys. Rev. Lett.* **102** 052301
[27] Adlarson P *et al* 2011 *Phys. Rev. Lett.* **106** 242302
[28] Adlarson P *et al* 2013 *Phys. Lett.*B **721** 229
[29] Adlarson P *et al* 2013 *Phys. Rev.* C **88** 055208
[30] Adlarson P *et al* 2015 *Phys. Lett.* B **743** 325
[31] Adlarson P *et al* 2016 *Eur. Phys. J.* A **52** 147
[32] Adlarson P *et al* 2014 *Phys. Rev. Lett.* **112** 202301
[33] Adlarson P *et al* 2014 *Phys. Rev.* C **90** 035204
[34] Bashkanov M, Clement H and Skorodko T 2015 *Eur. Phys. J.* A **51** 87
[35] Bashkanov M *et al* 2019 *Phys. Lett.* B **789** 7
[36] Bashkanov M *et al* 2020 *Phys. Rev. Lett.* **124** 132001
[37] Bashkanov M *et al* 2023 *Phys. Lett.* B **844** 138080
[38] Bashkanov M, Watts D P and Pastore A 2020 *Phys. Rev.* C **100** 012201
[39] Watts D P, Bashkanov M and Ostrick M 2020 Photodsintegration of transversely polarised deuterons, A2 proposal
[40] Vidaña I, Bashkanov M, Watts D P and Pastore A 2018 *Phys. Lett.* B **781** 112–6
[41] Bashkanov M and Watts D P 2020 *J. Phys.* G **47** 03LT01
[42] Bashkanov M, Brodsky S J and Clement H 2013 *Phys. Lett.* B **727** 438–42
[43] Harvey M 1979 *Nucl. Phys.* A **352** 301
[44] Kukulin V *et al* 2020 *Eur. Phys. J.* A **56** 229
[45] Dong Y, Shen P and Zhang Z 2023 *Prog. Part. Nucl. Phys.* **131** 104045
[46] Gell-Mann M The eightfold way: a theory of strong interaction symmetry *Synchrotron Laboratory* (https://doi.org/10.2172/4008239)
[47] Th A, Aerts M, Mulders P J G and de Swart J J 1978 *Phys. Rev.* D **17** 260

[48] Goldman J T, Maltman K, Stephenson G J, Schmidt K E and Wang F 1989 *Phys. Rev.* C **39** 1889–95
[49] Mulders P J and Thomas A W 1983 *J. Phys.* G **9** 1159
[50] Maltman K 1985 *Nucl. Phys.* A **438** 669–84
[51] Krane K S *Introductory Nuclear Physics*
[52] Buchmann A J and Henley E M 2000 *Phys. Lett.* B **484** 255
[53] Dong Y, Shen P and Zhang Z 2018 *Phys. Rev.* D **97** 114002
[54] Buchmann A J, Wagner G and Faessler A 1998 *Phys. Rev.* C **57** 3340
[55] Zhang T G, Dai L R, Cai X J, Chen L N and Wang Y H 2022 *Eur. Phys. J.* A **58** 200
[56] Pricking A, Bashkanov M and Clement H arXiv:1310.5532
[57] Bashkanov M, Clement H and Skorodko T 2017 *Nucl. Phys.* A **958** 129
[58] Adlarson P *et al* 2020 *Phys. Rev.* C **102** 015204
[59] Gal A 2017 *Phys. Lett.* B **769** 436
[60] Adlarson P *et al* 2016 *Phys. Lett.* B **762** 455–61
[61] Gongyo S *et al* 2018 *Phys. Rev. Lett.* **120** 212001
[62] Arndt R A 1968 *Phys. Rev.* **165** 1834
[63] Arndt R A, Roper L D, Workman R L and McNaughton M W 1992 *Phys. Rev.* C **45** 399
[64] Ungricht E *et al* 1985 *Phys. Rev.* C **31** 934
[65] Hoshizaki N 1992 *Phys. Rev.* C **45** R1414
[66] Hoshizaki N 1993 *Prog. Theor. Phys.* **89** 245
Hoshizaki N 1993 *Prog. Theor. Phys.* **251** 569
[67] Gal A 2017 *Phys. Lett.* B **769** 436–40
Gal A and Garcilazo H 2013 *Phys. Rev. Lett.* **111** 172301
[68] Niskanen J A arXiv: 2305.08647
[69] Adlarson P *et al* 2018 *Phys. Rev. Lett.* **121** 052001
[70] Acharya S *et al* 2020 *Nature* **588** 232–8
[71] Iritani T *et al* 2019 *Phys. Lett.* B **792** 284–9