## Masses of the $70^-$ Baryons in Large $N_c$ QCD

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## Abstract

The masses of the negative parity 70-plet baryons are analyzed in large  $N_c$  QCD to order  $1/N_c$ and to first order in SU(3) symmetry breaking. The existing experimental data are well reproduced and twenty new observables are predicted. The leading order SU(6) spin-flavor symmetry breaking is small and, as it occurs in the quark model, the subleading in  $1/N_c$  hyperfine interaction is the dominant source of the breaking. It is found that the  $\Lambda(1405)$  and  $\Lambda(1520)$  are well described as three-quark states and spin-orbit partners. New relations between splittings in different SU(3)multiplets are found.

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The study of excited mesons and baryons has been largely the domain of the quark model [1]. Despite the success of this model in reproducing the general features of the spectrum and decays, it is clear that in its different versions it is not a complete representation of QCD. One consequence of this incompleteness is that in those cases where the quark model does not agree with phenomenology, such as the problem of the mass splittings between spin-orbit partners in the negative parity baryons (spin-orbit puzzle), it is not clear whether the problem is due to the quark model itself or to specific dynamical properties of the states involved. In the last few years it was realized that the  $1/N_c$  expansion can provide a link between the phenomenology of excited baryons and QCD that avoids the assumptions required in the quark model. This link has the form of an effective theory that implements an expansion in  $1/N_c$ . As shown in this paper, the main features of the quark model emerge unscathed from the large  $N_c$  analysis of the masses in the 70-plet of negative parity baryons, and in addition some of the missing pieces of the model are recovered. For example, the long standing spin-orbit puzzle seems to be easily resolved by the presence of one operator of  $\mathcal{O}(N_c^0)$  that has no counterpart in the quark model.

In the  $N_c \to \infty$  limit of QCD the ground state baryons display an exact dynamical spinflavor symmetry SU(2F) (F is the number of flavors, equal to three in this Letter), which is a consequence of unitarity in pion-nucleon scattering in that limit [2]. In general SU(2F)is broken at  $\mathcal{O}(1/N_c)$  and for some observables even at  $\mathcal{O}(1/N_c^2)$  [3]. This implies that perturbation theory around the SU(2F) symmetric limit in the form of a  $1/N_c$  expansion is a powerful tool of analysis, as shown in numerous works [4]. In the context of the  $1/N_c$ expansion the sector of excited baryons is less understood. The principal reason is that even in the  $N_c \rightarrow \infty$  limit there is no exact dynamical symmetry [5]. However, it is an observational fact that excited states can be classified into multiplets of spin-flavor SU(2F). For example, most of the known baryons of negative parity seem to fit very well in the (3,70) irreducible representation (irrep) of  $O(3) \otimes SU(6)$ . As a matter of consistency it is shown in this work that the leading order spin-flavor breaking  $(\mathcal{O}(N_c^0))$  is small, thus justifying SU(2F) as an approximate symmetry. Under the assumption that the allowed mixing between states in different spin-flavor multiplets is suppressed by dynamical reasons, the  $1/N_c$  expansion for the full 70-plet can be implemented along the lines recently developed for two flavors [5, 6].

The states in the (3,70) decompose in terms of  $SU(2) \otimes SU(3)$  into two octets with

total angular momentum J = 1/2 (<sup>2S+1</sup> $d = {}^{2}8$  and <sup>4</sup>8, where S is the total spin and d the degeneracy of the SU(3) irrep), two octets with J = 3/2 ( ${}^{2}8$  and  ${}^{4}8$ ), one octet with J = 5/2 (<sup>4</sup>8), one decuplet with J = 1/2 and one with J = 3/2 (both  ${}^{2}10$ ), and two singlet As with J = 1/2 and 3/2 (both  ${}^{2}1$ ).

Since in the large  $N_c$  limit baryons consist only of valence quarks, it is natural to have an intuitive non-relativistic quark model picture of the spin-flavor composition of the states. This only means that the identification of spin-flavor states in the large  $N_c$  analysis and the quark model is the same. Thus, the wave functions are constructed by coupling an orbitally excited quark with  $\ell = 1$  to  $N_c - 1$  s-wave quarks that are in a spin-flavor symmetric core. The states have the general form

$$|\Psi\rangle = |J, J_z; S; (\lambda, \mu), Y, I, I_z \rangle$$
  
= 
$$\sum_{\alpha, \alpha', \alpha''} C\mathcal{G}_{\alpha, \alpha', \alpha''} |l\rangle_{\alpha} |q\rangle_{\alpha'} |c\rangle_{\alpha''} , \qquad (1)$$

where  $\alpha$  stands for the different projection quantum numbers and  $\mathcal{CG}$  for Clebsch-Gordan coefficients. In addition, the  $(\lambda, \mu)$  labels indicate the SU(3) irrep, Y is the hypercharge, I the isospin and  $J_z$ ,  $I_z$  the obvious projections. The (3,70) states are taken to have strangeness of order  $N_c^0$ . The excited quark and core states are given in terms of their  $SU(2) \otimes SU(3)$ quantum numbers with obvious notation:

$$|q\rangle = \begin{vmatrix} \frac{1}{2} & (1,0) \\ s_z & (y,\frac{1}{2},i_z) \end{vmatrix} , \ |c\rangle = \begin{vmatrix} S^c & (\lambda^c,\mu^c) \\ S^c_z & (Y^c,I^c,I^c_z) \end{vmatrix} ,$$
(2)

where  $S^c$  is the spin of the core. From the decomposition of the SU(6) symmetric representation into representations of  $SU(2) \otimes SU(3)$  the relations  $\lambda^c + 2\mu^c = N_c - 1$  and  $\lambda^c = 2S^c$ follow. They are the generalization of the I = J rule well known for two flavors. The total wave function is in the mixed symmetric irrep of SU(6). In the <sup>2</sup>8 representation a linear combination of two core states appears, namely  $|S^c, (\lambda^c, \mu^c)\rangle \ge |0, (0, \frac{N_c-1}{2})\rangle$  and  $|1, (2, \frac{N_c-3}{2})\rangle$ , while in the <sup>4</sup>8 and <sup>2</sup>10 representations the core state is  $|1, (2, \frac{N_c-3}{2})\rangle$ , and finally, in the <sup>2</sup>1 representation the core state is  $|0, (0, \frac{N_c-1}{2})\rangle$ .

A basis of mass operators can be built using the generators of  $O(3) \otimes SU(2F)$  [5]. A generic *n*-body mass operator has the general structure

$$O^{(n)} = \frac{1}{N_c^{n-1}} O_\ell O_q O_c , \qquad (3)$$

where the factors  $O_{\ell}$ ,  $O_q$ , and  $O_c$  can be expressed in terms of products of generators of orbital angular momentum  $(\ell_i)$ , spin-flavor of the excited quark  $(s_i, t_a \text{ and } g_{ia} \equiv s_i t_a)$  and spin-flavor of the core  $(S_i^c, T_a^c \text{ and } G_{ia}^c \equiv \sum_{m=1}^{N_c-1} s_i^{(m)} t_a^{(m)})$ , respectively. The explicit  $1/N_c$ factors originate in the n-1 gluon exchanges required to give rise to a n-body operator. The matrix elements of operators may also carry a nontrivial  $N_c$  dependence due to coherence effects [2, 3]: for the states considered,  $G_{ia}^c$  (a = 1, 2, 3) and  $T_8^c$  have matrix elements of  $\mathcal{O}(N_c)$ , while the rest of the generators have matrix elements of higher order.

For  $N_c = 3$  and in the SU(3) limit there are eleven independent quantities, namely nine masses and two mixing angles  $\theta_1$  and  $\theta_3$ , which correspond to the mixing of the <sup>2</sup>8 and <sup>4</sup>8 irreps in the J = 1/2 and J = 3/2 octets, respectively. There is, therefore, a basis of eleven SU(3)-singlet mass operators. As shown in Table I, the basis of singlet operators  $O_i$  consists of one operator of  $\mathcal{O}(N_c)$ , namely the identity operator, three operators of  $\mathcal{O}(N_c^0)$ , that include the spin-orbit operator, and seven of  $\mathcal{O}(1/N_c)$ , one of which is the very important hyperfine operator. These operators are a simple generalization of those known for two flavors [7].

When SU(3) breaking is included with isospin conservation, the number of independent observables raises up to fifty, of which thirty are masses and twenty are mixing angles. However, if SU(3) symmetry breaking is restricted to linear order in quark masses only isosinglet octet operators can appear, and the number of independent observables is reduced to thirty five (twenty one masses and fourteen mixing angles) implying twenty four linearly independent octet mass operators. As a consequence of this reduction several mass relations exist, among them there is a Gell-Mann Okubo relation for each octet and an equal spacing rule for each decuplet. The octet contributions are proportional to  $\epsilon \propto (m_s - m_{u,d})/\nu_H$ where  $\nu_H$  is a typical hadronic mass scale, for instance  $m_{\rho}$ ; for  $N_c = 3$  the quantity  $\epsilon$  counts as of the same order as  $1/N_c$ . Explicit construction shows that up to order  $\mathcal{O}(\epsilon N_c^0)$  only a small subset of independent octet operators  $B_i$  appears. Since such octet operators are isospin singlets, it is possible to modify them by adding singlet operators so that the resulting operators vanish in the subspace of non-strange baryons. This procedure of improving the flavor breaking operators may change the  $1/N_c$  counting: for instance, after improving  $T_8$ with the identity operator  $O_1$  the resulting operator is of order  $N_c^0$ . Indeed, the improved operators give the splitting due to SU(3) breaking with respect to the non-strange baryons in each multiplet, and they must be of zeroth order or higher in  $1/N_c$  for states with strangeness of order  $N_c^0$ . The four improved flavor breaking operators  $\bar{B}_1$  through  $\bar{B}_4$  that remain at  $\mathcal{O}(\epsilon N_c^0)$  when  $N_c = 3$  are shown in Table I.

As a result of the above analysis the 70-plet mass operator up to  $\mathcal{O}(1/N_c)$  and  $\mathcal{O}(\epsilon N_c^0)$  has the most general form:

$$M_{70} = \sum_{i=1}^{11} c_i O_i + \sum_{i=1}^{4} d_i \bar{B}_i \quad , \qquad (4)$$

where  $c_i$  and  $d_i$  are numerical coefficients which can be determined by fitting the available empirical masses and mixing angles. For this purpose it is necessary to have the expressions of the matrix elements of the  $O_i$  and  $\bar{B}_i$  operators between a basis of states belonging to the 70-plet. Their analytic expressions are obtained using standard techniques and will be given elsewhere.

Because at  $\mathcal{O}(\epsilon N_c^0)$  there are only four flavor breaking operators, it is possible to find new mass splitting relations which are independent of the coefficients  $d_i$ . These relations involve states in different SU(3) multiplets. Of particular interest are the following five relations that result when the operator  $\bar{B}_3$  is neglected (from the fit below it is apparent that  $\bar{B}_3$  gives very small contributions):

$$9(s_{\Sigma_{1/2}} + s_{\Sigma_{1/2}'}) + 21s_{\Lambda_{5/2}} = 17(s_{\Lambda_{1/2}} + s_{\Lambda_{1/2}'}) + 5s_{\Sigma_{5/2}},$$

$$2(s_{\Lambda_{3/2}} + s_{\Lambda_{3/2}'}) = 3s_{\Lambda_{5/2}} + s_{\Sigma_{5/2}},$$

$$18(s_{\Sigma_{3/2}} + s_{\Sigma_{3/2}'}) + 33s_{\Lambda_{5/2}} = 28(s_{\Lambda_{1/2}} + s_{\Lambda_{1/2}'}) + 13s_{\Sigma_{5/2}},$$

$$9s_{\Sigma_{1/2}''} = s_{\Lambda_{1/2}} + s_{\Lambda_{1/2}'} + 3s_{\Lambda_{5/2}} + 4s_{\Sigma_{5/2}},$$

$$18s_{\Sigma_{3/2}''} + 3s_{\Lambda_{5/2}} = 8(s_{\Lambda_{1/2}} + s_{\Lambda_{1/2}'}) + 5s_{\Sigma_{5/2}}.$$
(5)

Here  $s_{\mathcal{B}_i}$  is the mass splitting between the baryon  $\mathcal{B}_i$  and the non-strange baryons in the SU(3) multiplet to which it belongs. These relations are independent of mixings because they result from relations among traces of the octet operators. If  $\bar{B}_3$  is not neglected there are instead four relations.

Discussion of the fit and conclusions – The experimental masses shown in Table II (three or more stars status in the the Particle Data listing [8]) together with the two leading order mixing angles  $\theta_1 = 0.61$ ,  $\theta_3 = 3.04$  [9, 10] are the 19 empirical quantities to be fitted. One three-star state is not included, namely the  $\Sigma(1940)$  which does not consistently fit into the 70-plet, and the two-star states  $\Sigma(1580)$  and  $\Sigma(1620)$  are not included as inputs. The errors in mass inputs are taken to be equal to the experimental errors if these are larger than the magnitude of the NLO corrections estimated at  $\pm 15$  MeV, otherwise they are taken to be equal to the latter. The fifteen coefficients  $c_i$ ,  $d_i$  are obtained from the fit, and the resulting  $\chi^2$  per degree of freedom of the fit turns out to be  $\chi^2/4 = 1.29$ . The results for the coefficients are displayed in Table I, while the best fit masses and state compositions are displayed in Table II. Note that the natural size of coefficients associated with singlet operators is set by the coefficient of  $O_1$ , and is about 500 MeV, while the natural size for the coefficients associated with octet operators is roughly  $\epsilon$  times 500 MeV.

There are a number of important points that emerge from this analysis.

Although spin flavor symmetry is broken at  $\mathcal{O}(N_c^0)$ , it is evident that the  $\mathcal{O}(N_c^0)$  operators are dynamically suppressed as their coefficients are substantially smaller than the natural size. It turns out that the chief contribution to spin-flavor breaking stems from the  $\mathcal{O}(1/N_c)$ hyperfine operator  $O_6$ , as in the ground state baryons. Since  $O_6$  is purely a core operator, it turns out that the gross spin-flavor structure of levels is determined by the two possible core states. This observation is in agreement with the findings of the quark model [10] where the strength of the hyperfine interaction is determined by the one-gluon exchange. In particular, the two singlet As are not affected by  $O_6$ , while the other states are moved upwards, explaining in a transparent way the lightness of these two states. Indeed, by keeping only  $O_1$  and  $O_6$  the <sup>2</sup>8 masses are 1510 MeV, the <sup>4</sup>8 and <sup>2</sup>10 masses are 1670 MeV, and the <sup>2</sup>1 masses are left at the bottom with 1350 MeV. This clearly shows the dominant pattern of spin-flavor breaking observed in the 70-plet.

The long standing problem in the quark model of reconciling the large  $\Lambda(1520) - \Lambda(1405)$ splitting with the splittings between the other spin-orbit partners in the 70-plet is resolved in the large  $N_c$  analysis. The singlet  $\Lambda$ s receive contributions to their masses from  $O_1$  and  $\ell$ .s while the rest of the operators give vanishing contributions because the core of the singlets carries zero spin and isospin. The splitting between the singlets is, therefore, a clear display of the spin-orbit coupling. The problem with the splittings between spin-orbit partners in the non-singlet sector, illustrated by the fact that the  $\ell$ .s operator gives a contribution to the  $\Delta_{1/2} - \Delta_{3/2}$  splitting that is of opposite sign of what is observed, is now solved by the presence of the operators  $O_4$ ,  $O_5$ ,  $O_9$  and  $O_{11}$ , with the contribution from  $O_4$  being the dominant one in accordance with the  $1/N_c$  counting. One important consequence of this result is that the interpretation of the singlet  $\Lambda$ s as three-quark states is consistent with the masses of the rest of the 70-plet. This further supports a similar claim drawn from scaling down to the strange sector the mass splitting between the  $\Lambda_c(2593)$  and the  $\Lambda_c(2625)$  [11].

There is a hierarchy of mixing effects. As already mentioned, at  $\mathcal{O}(N_c^0)$  there are two mixing angles, namely  $\theta_1$  and  $\theta_3$  that mix the octets with same J. These mixing angles are inputs and are obtained from an analysis of the  $N^*$  decays [9, 10]. All  $\mathcal{O}(N_c^0)$  operators in principle contribute to these mixings, but the  $\ell$ .s and  $O_4$  contributions tend to cancel each other leaving the  $O_3$  as the dominant one. Indeed, the coefficient of  $O_3$  is largely determined by mixing as this operator gives only modest contributions to the masses [12]. The rest of the mixings are of higher order because they are due to SU(3) breaking, and expected to be small. Table II shows this in the composition of states. A good example is that the  $\Lambda(1405)$  and  $\Lambda(1520)$  remain largely singlet states. In some cases, however, due to close degeneracy the SU(3) breaking can induce a larger than expected mixing angle which cannot be predicted with precision from an analysis of the masses alone. This occurs in the  $J = 3/2 \Sigma s$  and  $\Xi s$ , where for that reason the corresponding amplitudes in Table II are shown between parentheses.

The hyperfine operator between the excited  $\ell = 1$  quark and the core,  $O_7$ , is suppressed with respect to  $O_6$ , indicating that the hyperfine interaction is of short range in agreement with the quark model. The large errors that make the coefficients compatible with zero show that the operators  $O_8$  and  $\bar{B}_3$  are largely irrelevant. On the other hand, the collective effects of the three-body operators  $O_9$ ,  $O_{10}$  and  $O_{11}$  amount to mass shifts of modest magnitude (50 MeV or less).

The first relation in equation (6) predicts the  $\Sigma_{1/2}$  to be 103 MeV above the  $N_{1/2}$ , consistent with the  $\Sigma(1620)$ , a two star state that is not included as input to the fit. Each of the remaining relations makes a similar prediction for other states but requires further experimental input to be tested.

The analysis of this Letter shows that the  $1/N_c$  expansion provides a systematic approach to the spectroscopy of the negative parity baryons. It successfully describes the existent data and to the order considered it also leads to numerous predictions yet to be tested. In addition to the well known Gell-Mann-Okubo and equal spacing relations, new splitting relations between different multiplets that follow from the spin flavor symmetry have been found. The  $\Lambda(1405)$  is well described as a three quark state and the spin orbit partner of the  $\Lambda(1520)$ . Finally, effective interactions that correspond to flavor quantum number exchanges, such as the ones mediated by the operators  $O_3$  and  $O_4$ , are apparently needed. Although the corresponding coefficients seem to be dynamically suppressed their relevance shows up in the well established finer effects, namely mixings and splittings between nonsinglet spin-orbit partners. These interactions are not accounted for in the standard quark model based on one gluon exchange.

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[12] Although the numerical results obtained in the present analysis are similar to those in the two-flavor case [6], our conclusions about the relevance of the  $O_3$  operator is different.

## Table captions

**TABLE I:** Operator list and best fit coefficients.

**TABLE II:** Masses and spin-flavor content as predicted by our large  $N_c$  analysis. Also given are the empirical masses and those obtained in the quark model (QM) calculation of Ref. [10].

Operator	Fitted coef. [MeV]
$O_1 = N_c \ 1$	$c_1 = 449 \pm 2$
$O_2 = l_h s_h$	$c_2 = 52 \pm 15$
$O_3 = \frac{3}{N_c} \ l_{hk}^{(2)} \ g_{ha} \ G_{ka}^c$	$c_3 = 116 \pm 44$
$O_4 = \frac{4}{N_c + 1} \ l_h \ t_a \ G_{ha}^c$	$c_4 = 110 \pm 16$
$O_5 = \frac{1}{N_c} l_h S_h^c$	$c_5 = 74 \pm 30$
$O_6 = \frac{1}{N_c} S_h^c S_h^c$	$c_6 = 480 \pm 15$
$O_7 = \frac{1}{N_c} \ s_h \ S_h^c$	$c_7 = -159 \pm 50$
$O_8 = \frac{1}{N_c} l_{hk}^{(2)} s_h S_k^c$	$c_8 = 6 \pm 110$
$O_9 = \frac{1}{N_c^2} l_h g_{ka} \{ S_k^c, G_{ha}^c \}$	$c_9 = 213 \pm 153$
$O_{10} = \frac{1}{N_c^2} t_a \{ S_h^c, G_{ha}^c \}$	$c_{10} = -168 \pm 56$
$O_{11} = \frac{1}{N_c^2} \ l_h \ g_{ha} \{ S_k^c, G_{ka}^c \}$	$c_{11} = -133 \pm 130$
$\bar{B}_1 = t_8 - \frac{1}{2\sqrt{3}N_c}O_1$	$d_1 = -81 \pm 36$
$\bar{B}_2 = T_8^c - \frac{N_c - 1}{2\sqrt{3}N_c}O_1$	$d_2 = -194 \pm 17$
$\bar{B}_3 = \frac{1}{N_c} d_{8ab} g_{ha} G^c_{hb} + \frac{N_c^2 - 9}{16\sqrt{3}N_c^2(N_c - 1)}O_1 + \frac{N_c^2}{16\sqrt{3}N_c^2(N_c - 1)}O_1 + \frac{N_c^2}{16\sqrt{3}$	
$+\frac{1}{4\sqrt{3}(N_c-1)}O_6 + \frac{1}{12\sqrt{3}}O_7$	$d_3 = -150 \pm 301$
$\bar{B}_4 = l_h g_{h8} - \frac{1}{2\sqrt{3}}O_2$	$d_4 = -82 \pm 57$

TABLE I

	Masses [MeV]			Spin-flavor content			
State	Expt.	Large $N_c$	QM	$^{2}1$	$^{2}8$	<sup>4</sup> 8	$^{2}10$
$N_{1/2}$	$1538 \pm 18$	1541	1490		0.82	0.57	
$\Lambda_{1/2}$	$1670\pm10$	1667	1650	-0.21	0.90	0.37	
$\Sigma_{1/2}$	(1620)	1637	1650		0.52	0.81	0.27
$\Xi_{1/2}$		1779	1780		0.85	0.44	0.29
$N_{3/2}$	$1523\pm8$	1532	1535		-0.99	0.10	
$\Lambda_{3/2}$	$1690\pm5$	1676	1690	0.18	-0.98	0.09	
$\Sigma_{3/2}$	$1675\pm10$	1667	1675		-0.98	-0.01	-0.19
$\Xi_{3/2}$	$1823\pm5$	1815	1800		-0.98	0.03	-0.19
$N_{1/2}'$	$1660\pm20$	1660	1655		-0.57	0.82	
$\Lambda'_{1/2}$	$1785\pm65$	1806	1800	0.10	-0.38	0.92	
$\Sigma_{1/2}'$	$1765\pm35$	1755	1750		-0.83	0.54	0.17
$\Xi_{1/2}'$		1927	1900		-0.46	0.87	0.18
$N_{3/2}^{\prime}$	$1700\pm50$	1699	1745		-0.10	-0.99	
$\Lambda'_{3/2}$		1864	1880	0.01	-0.09	-0.99	
$\Sigma_{3/2}^{\prime}$		1769	1815		0.01	(-0.57)	(-0.82)
$\Xi_{3/2}'$		1980	1985		-0.02	(-0.57)	(-0.82)
$N_{5/2}$	$1678\pm8$	1671	1670			1.00	
$\Lambda_{5/2}$	$1820\pm10$	1836	1815			1.00	
$\Sigma_{5/2}$	$1775\pm5$	1784	1760			1.00	
$\Xi_{5/2}$		1974	1930			1.00	
$\Delta_{1/2}$	$1645\pm30$	1645	1685				1.00
$\Sigma_{1/2}''$		1784	1810		-0.14	-0.31	0.94
$\Xi_{1/2}^{\prime\prime}$		1922	1930		-0.14	-0.31	0.94
$\Omega_{1/2}$		2061	2020				1.00
$\Delta_{3/2}$	$1720\pm50$	1720	1685				1.00
$\Sigma_{3/2}^{\prime\prime}$		1847	1805		-0.19	(-0.80)	(0.57)
$\Xi_{3/2}^{\prime\prime}$		1973	1920		-0.19	(-0.80)	(0.57)
$\Omega_{3/2}$		2100	2020				1.00
$\Lambda_{1/2}''$	$1407 \pm 4$	1407	1490	0.97	0.23	0.04	
$\Lambda_{3/2}''$	$1520 \pm 1$	1520	1490	0.98	0.18	-0.01	

TABLE II