# THE $N^* \to N\eta$ BRANCHING RATIOS AND THE QUARK MODEL*

D.O. Riska

Department of Physics, University of Helsinki
Helsinki 00014, Finland

(Received October 9, 1996)

The selectivity of the $\eta$-decay branching ratios of the negative parity baryon resonances is discussed from a quark model perspective. It is shown that the selectivity can be understood as a consequence of the selective clustering into quark-diquark configurations with different quantum numbers that is caused by the pseudoscalar octet chiral meson exchange interaction between the constituent quarks. This leads to the prediction that the resonances $N(1535)$, $\Lambda(1670)$, $\Sigma(1750)$ with mixed flavour-spin symmetry $[21]_{\mathrm{FS}}[21]_{\mathrm{F}}[21]_{\mathrm{S}}$ in their lowest order wavefunctions should have large $\eta$ branching ratios, whereas the $N(1650)$ and the $\Lambda(1800)$ should not.

PACS numbers: 13.30. -a, 13.30. Eg

## 1. Introduction

The $\eta$-decay branching ratios of the negative parity baryon resonances show a remarkable selectivity. While the lowest $\frac{1}{2}^-$ nucleon resonance, the $N(1535)$ has a 30–50% $N\eta$ branching ratio, the following $\frac{1}{2}^-$ resonance, the $N(1650)$ has at most a 1% branching ratio for $N\eta$ decay. A similar feature is revealed in the spectrum of the $\Lambda$-hyperon, where the lowest $\frac{1}{2}^-$ resonance that lies above the $\Lambda\eta$ threshold, the $\Lambda(1670)$, has a large 15–35% $\Lambda\eta$-decay branch, whereas the following $\frac{1}{2}^-$ resonance, the $\Lambda(1800)$ has no or at most a very small branching ratio for $\Lambda\eta$ decay. Finally, the lowest negative parity $\Sigma$-hyperon resonance, the $\Sigma(1750)$, which lies above the $\Sigma\eta$ threshold, again has a large (15–55%) branching ratio for $\Sigma\eta$ decay. These resonances show no comparable selectivity in their corresponding $N\pi$ and $N\bar{K}$ decays. It is therefore natural to assume that this $\eta$-decay selectivity, which seems to be a common feature of the different flavour sectors of the baryon spectrum,

* Presented at the "Meson 96" Workshop, Cracow, Poland, May 10–14, 1996.

derives from some universal feature of the hyperfine interaction between the constituent quarks that form these resonances. As the $\eta$-decay selectivity appears for states that have equal quantum numbers, and which differ only in mass, this selectivity would be hard to explain without consideration of the substructure of these resonances.

When the baryons are described as formed of 3 constituent quarks that are confined by a flavour- and spin-independent confining interaction (*e.g.* a linear quark-quark potential) the baryon spectrum is organized into successive bands (shells) of degenerate states with alternating parity. A peculiar feature of the empirical baryon spectrum is that with exception of the $\Lambda$-hyperon spectrum, the lowest positive parity resonances that belong to the *sd*-shell lie below the lowest negative (or *p*-shell) states. This shows that the hyperfine interaction has to be flavour dependent as well as strong enough to reverse the "normal" ordering of the lowest positive and negative parity resonances. Such a strong state-dependent hyperfine interaction will cause a clustering of the 3-quark state into quark-diquark configurations with state dependent quantum numbers. It has recently been pointed out that if the main hyperfine interaction between the constituent quarks is due to the exchange of the octet of light pseudoscalar mesons, which are the Goldstone bosons of the approximate chiral symmetry of QCD, this clustering effect provides a qualitative explanation of the $\eta$-decay selectivity of the negative parity (*p*-shell) baryon resonances [1].

The attempt to describe the large hyperfine splitting of the baryon spectrum in terms of exchange of the light $(\pi, K, \eta)$ pseudoscalar octet mesons fits within the two-scale scenario of Manohar and Georgi [2], according to which the appropriate effective degrees of freedom for describing the structure of the baryons at intermediate range/energy are constituent quarks and the Goldstone bosons. It was in fact found in [3] that the spin-flavour dependence of the pseudoscalar meson exchange interaction between the constituent quarks has precisely the right spin-flavour dependence for reversing the "normal" ordering of the lowest *p* and *sd*-shell states in the spectra of the nucleon and the $\Delta$-resonance thus bringing the spectral ordering in conformity with the empirical one. Numerically exact solution of the Faddeev equations for the 3-quark system with a linear confining potential and a pseudoscalar meson exchange interaction with the expected strength bear this out quantitatively [4].

Below the application of this chiral version of the constituent quark model to the problem of the selectivity of the $\eta$-decay branches of the *p*-shell baryon resonances is reviewed in Section 2. A concluding discussion is given in Section 3.

## 2. The selectivity of the $\eta$-decays

In Table I we list the possible flavour-spin symmetries of the 3-quark configurations for the $p$-shell ($L = 1$) baryon resonances. The permutational symmetries of the flavour-spin (FS), flavour (F) and spin (S) components of the wavefunctions in the absence of hyperfine splitting are denoted by the young patterns $[f]_{\rm FS}[f]_{\rm F}[f]_{\rm S}$. Here $[f]$ represents a Young pattern described by a sequence of integers $f$, which lists the number of boxes in the successive rows in the corresponding Young pattern. By the Pauli principle the permutational symmetry of the spatial state is then also [21] in the $p$-shell. In the table the configurations of the zero order wave functions of the known $N, \Delta, \Lambda$ and $\Sigma$ states with $L = 1$ have also been indicated.

TABLE I

The configurations of the lowest order wavefunctions of the $p$-shell baryon resonances. The resonances with large $\eta$-decay branches are marked with an asterisk.

| $[f]_{\rm FS}[f]_{\rm F}[f]_{\rm S}$ | $N, \Delta$ | $\Lambda$ | $\Sigma$ |
|---|---|---|---|
| $[21]_{\rm FS}[111]_{\rm F}[21]_{\rm S}$ | - | $\Lambda(1405), \Lambda(1520)$ | - |
| $[21]_{\rm FS}[21]_{\rm F}[21]_{\rm S}$ | $N(1535)^*, N(1520)$ | $\Lambda(1670)^*, \Lambda(1690)$ | $\Sigma(1750)^*, ?$ |
| $[21]_{\rm FS}[3]_{\rm F}[21]_{\rm S}$ | $\Delta(1620), \Delta(1700)$ | - | $\Sigma(1620), \Sigma(1670)?$ |
| $[21]_{\rm FS}[21]_{\rm F}[3]_{\rm S}$ | $N(1650), N(1700),$ $N(1675)$ | $\Lambda(1800), ?,$ $\Lambda(1830)$ | $?, \Sigma(1670)?,$ $\Sigma(1775)$ |

The $\eta$-decays involve transtions between the different $p$-shell states and the ground state, which has the flavour-spin configuration $[3]_{\rm FS}[21]_{\rm F}[21]_{\rm S}$. The transition operator in the case of $\eta$-decay will be a sum of $\eta$-quark transition operators of the form

$$\hat{0} = \frac{g}{2m_q}\lambda^8 \vec{\sigma} \cdot \vec{q} e^{i\vec{q}\cdot\vec{r}}, \tag{1}$$

where $g$ is the $\eta$-quark pseudoscalar coupling constant and $m_q$ is the constituent quark mass. Note that the $\lambda^8$ operator does not change the quark flavour.

Consider now the percentage of diquark components with different flavour spin symmetry in the $p$-shell baryons. The nucleon state is formed of two components of equal weight with the flavour-spin symmetries $[11]_{\rm F}[11]_{\rm S}$, $[2]_{\rm F}[2]_{\rm F}$ for the diquark (formed say of quarks 1 and 2). The states of the $p$-shell baryons with the configuration $[21]_{FS}[21]_F[21]_S$ contain in contrast all possible diquark flavour-spin combinations $[11]_F[11]_S$, $[2]_F[11]_S$, $[11]_F[2]_S$ and $[2]_F[2]_S$ with equal weight. The states of the $p$-shell baryons with the configuration $[21]_{FS}[21]_F[3]_S$ contains only the diquark combinations $[11]_F[2]_S$ and $[2]_F[2]_S$, again with equal weight. Finally the decouplet baryons with the configurations $[21]_{FS}[3]_F[21]_S$ contain only the diquark combinations $[2]_F[2]_S$ and $[2]_F[11]_S$.

As both of the $p$-shell baryon states $[21]_{FS}[21]_F[21]_S$ and $[21]_{FS}[21]_F[3]_S$, *e.g.* the $N(1535)$ and $N(1650)$, contain diquark configurations present in the nucleon state, the selectivity of their $\eta$-decays obviously cannot be explained at the level of these zero-order wavefunctions, without account of the hyperfine interaction [1]. That the effect of the latter will be important follows from the fact that it has to be strong enough to reverse the ordering of the low lying states in the $p$ and $sd$-shells.

The main component of the chiral boson exchange hyperfine interaction between the constituent quarks has the form of a flavour dependent spin-spin interaction

$$H_\chi \sim -\sum_{i<j} V(r_{ij})\vec{\lambda}_i^F \cdot \vec{\lambda}_j^F \vec{\sigma}_i \cdot \vec{\sigma}_j. \qquad (2)$$

The potential function $V(r)$ will at long range have the usual Yukawa function behaviour, whereas at short range it will qualitatively behave as the $\delta$-function term in the pseudoscalar exchange interaction smeared over the size of the constituent quarks. As an example the explicit form used in Ref. [4] for $V(r)$ in combination with a linear confining potential is

$$V(r) = \frac{g^2}{4\pi}\frac{1}{4m_1m_2}\{\mu^2\frac{e^{-\mu r}}{r}H(r) - \frac{4}{\sqrt{\pi}}\alpha^3 e^{-\alpha^2(r-r_0)^2}\}, \qquad (3)$$

where $g$ is the quark-meson coupling constant ($g^2/4\pi = 0.67$), $m_1$ and $m_2$ are the constituent masses of the interacting quarks and $\mu$ is the meson mass. In (3) $H(r)$ is a smooth short-range cut-off function and $\alpha^{-1}$ and $r_0$ are "smearing" parameters of the order $0.4 fm$. It is important to note, that if the hyperfine interaction (3) is treated in first order perturbation theory, the $s$- and $p$-shell matrix elements of the hyperfine interaction potential $V(r)$ are of the same order of magnitude.

The last point makes it appropriate to compare the relative strengths (and signs) of the hyperfine interaction (2) in quark pair states with different flavour and spin symmetry $[f]_F[f']_S$. One readily finds that

$$\langle [f]_{\rm F}[f']_{\rm S}| - V(r)\vec{\lambda}^1 \cdot \vec{\lambda}^2 \vec{\sigma}^1 \cdot \vec{\sigma}^2 | [f]_{\rm F}[f']_{\rm S}\rangle = \begin{cases} -\frac{4}{3}V(r), & [2]_{\rm F}[2]_{\rm S} \\ -8V(r), & [11]_{\rm F}[11]_{\rm S} \\ +4V(r), & [2]_{\rm F}[11]_{\rm S} \\ +\frac{8}{3}V(r), & [11]_{\rm F}[2]_{\rm S}. \end{cases} \quad (4)$$

This shows that the hyperfine interaction is 6 times more attractive in diquark states with antisymmetric $[11]_{\rm F}[11]_{\rm S}$ flavour and spin states than in those with symmetric $[2]_{\rm F}, [2]_{\rm S}$ flavour and spin states. Moreover, as shown in [1] the baryon spectrum indicates that the matrix element of the hyperfine interaction is $\sim 3.3$ times larger than the matrix element of the effective confining interaction in the ground state. This implies that the 3 quarks in the ground states will cluster into a quark-diquark configuration, the latter being very compact and having $[11]_{\rm F}[11]_{\rm S}$ flavour and spin symmetries ($S = T = 0$ in the case of the nucleon). As the decuplet baryons lack diquark configurations with antisymmetric flavour states, there will be no corresponding clusterization in their ground states.

In the $p$-shell states with the zero-order state configuration $[21]_{\rm FS}[21]_{\rm F}$ $[3]_{\rm S}$ there will be a corresponding clusterization into a quark-diquark configuration, but in this case the compact diquark will have symmetric flavour and spin states: $[2]_{\rm F}[2]_{\rm S}$. The reason is the combined effect of the weak attractive hyperfine interaction in the $[2]_{\rm F}[2]_{\rm S}$ diquark state and the stronger repulsive interaction in the $[11]_{\rm F}[2]_{\rm S}$ state. The net effect of the weak "pull" in the former and strong "push" in the latter will cause a clustering into the $[2]_{\rm F}[2]_{\rm S}$ diquark configuration. Now as the transition operator (1) for $\eta$ decay cannot connect the $[11]_{\rm F}[11]_{\rm S}$ and $[2]_{\rm F}[2]_{\rm F}$ diquark states in the ground state and the excited state with $[21]_{\rm FS}[21]_{\rm F}[3]_{\rm F}$ (or with total spin $\frac{3}{2}$), $\eta$-decay of these $p$-shell resonances should be suppressed. This provides a qualitative explanation of the empirical fact that the $N(1650)$ and $\Lambda(1800)$ $\frac{1}{2}^-$ resonances have very small $\eta$-decay branches. An immediate consequence of this is also that the corresponding $\frac{3}{2}^-$ and $\frac{5}{2}^-$ (*e.g.* the $\Lambda(1830)$) resonances should have very small $\eta$-decay branches. There will be no corresponding suppression of the $\pi$ and $\bar{K}$ decays of these resonances as the transition operators for those decays contain flavour operators $\lambda_1 ... \lambda_7$, that can connect diquark states with different flavour symmetry.

In contrast the $p$-shell state with the zero-order configuration $[21]_{\rm FS}[21]_{\rm F}$ $[21]_{\rm F}$ will experience a clusterization into a quark-diquark structure, where the diquark has the same symmetric flavour-spin states $[2]_{\rm F}[2]_{\rm S}$ as the nucleon. There is therefore no flavour change (or isospin flip) required in the meson decay, and therefore the $\eta$-transition operator can connect the diquarks of the resonance and the nucleon. This is the reason for the large

$\eta$-decay branches of *e.g.* the $N(1535)$ and $\Lambda(1670)$ resonances. While the $\frac{3}{2}^-$ resonances $N(1520)$ and $\Lambda(1690)$ have the same zero order configurations, their small $\eta$-decay branches can be understood as a consequence of the kinematic suppression of the required $D$-wave amplitude for $\eta$-decay near the threshold.

Finally the $\frac{1}{2}^+$ $N(1710)$ resonance in the *sd* shell also has the flavour-spin symmetry $[21]_{FS}[21]_F[21]_S$. This should by the above argument also have a large $N\eta$ decay branch, prediction that conforms well with the empirical $N\eta$ branching ratio of 20–40%. For the same reason the $\frac{1}{2}^+$ $\Sigma(1880)$ resonance should also have a large $\Sigma\eta$ decay branch.

The present model, according to which the *p*-shell baryon resonances with zero order wavefunctions with the $[21]_{FS}[21]_F[21]_S$ configuration, have a strong $\eta$ decay branch implies that the $\Sigma(1750)$ $\frac{1}{2}^-$ resonance that has a large $\Sigma\eta$ decay branch should also have this zero order configuration. As a corollary this should be the lowest negative parity $\Sigma$-hyperon resonance, and then the 2-star $\Sigma(1580)$ and $\Sigma(1620)$ resonances should be spurious. It should be interesting to have this issue settled experimentally.

## 3. Discussion

The argument above ties the $\eta$-decay selectivity of the negative parity baryon resonances to the strength and spin-flavour properties of the chiral meson exchange mediated model for the hyperfine interaction (2) between the constituent quarks. This model assumes that the coupling between the chiral boson octet and the constituent quarks is SU(3) symmetric, and has the form $g_q\gamma_5\lambda^a$, as *e.g.* in an extended linear $\sigma$-model. It is worth noting that this form of the quark-octet coupling implies the following relation between the $\pi$-nucleon and $\eta$-nucleon (pseudoscalar) coupling constants:

$$g_{\eta NN} = \frac{\sqrt{3}}{5} g_{\pi NN}. \tag{5}$$

This provides a natural explanation for the relative smallness of the $\eta NN$ coupling constant as seen *e.g.* in phenomenological boson exchange models for the nucleon-nucleon interaction [5].

The pseudoscalar meson exchange model for the constituent quarks involves charge exchange between the quarks. This by the continuity equation implies the presence of a meson exchange current operator between the quarks, which leads to a magnetic moment operator with the form

$$\mu^{ex} = \mu_N\{V_\pi(r_{ij})(\tau^i \times \tau^j)_3 + V_K(r_{ij})(\lambda_4^i\lambda_5^j - \lambda_5^i\lambda_4^j)\}(\vec{\sigma}_i \times \vec{\sigma}_j). \tag{6}$$

This operator plays an important role in compensating the large relativistic corrections to the magnetic moments of the baryons predicted by the static quark model [6]. Without the exchange current contribution the relativistic corrections destroy the otherwise satisfactory predictions of the static quark model. The relativistic corrections on the other hand provide a simple and natural explanation of why the prediction of the static quark model for the axial coupling constants of the baryons are systematically larger than the empirical values [6].

The pseudoscalar meson exchange model for the hyperfine interaction between the constituent quarks thus, besides providing a explanation for the ordering of the positive and negative parity states in the baryon spectrum, also provides a simple and at least qualitative explanation for both the $\eta$-decay selectivity of the negative parity baryon resonances and for the magnetic moments and the axial coupling constants of the baryons.

This presentation is based on research conducted in collaboration with Dr. L. Ya. Glozman.

## REFERENCES

[1] L.Ya. Glozman, D.O. Riska, *Phys. Lett.* **B366**, 305 (1996).
[2] A. Manohar, H Georgi, *Nucl. Phys.* **B234**, 189 (1984).
[3] L.Ya. Glozman, D.O. Riska, *Phys. Rep.* (1996).
[4] L.Ya. Glozman, Z. Papp, W. Plessas, Light Baryons in a Constituent Quark Model with Chiral Dynamics, Preprint UNIGRAZ-UTP 27-01-96 (1996).
[5] R. Machleidt, K. Holinde, Ch. Elster, *Phys. Rep.* **149**, 1 (1987).
[6] K. Dannbom, L.Ya. Glozman, C. Helminen, D.O. Riska, in preparation.