ABSTRACT
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OPTIMUM CHANNELS FOR EXOTIC Z^* BARYONS

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1. Introduction

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Abstract

We investigate the possibility of exotic, negative parity, strange-ness $+1$ ($4q\bar{s}$) baryons in the context of the constituent quark model. These multiquark states, whose valence configurations differ from those of conventional mesons and baryons, are capable of providing unambiguous evidence for the presence of multi-quark physics. We find that the $I=1$ $S=5/2$ ΔK^* resonance is the strongest candidate for such a state and an $I=0$ $S=3/2$ state is possible, although less strongly indicated. It is unlikely that any other ($4q\bar{s}$) Z^* states exist.

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The notion that ordinary hadrons are composed of quarks carrying a confined $SU(3)$ colour degree of freedom leads to the possibility that the physical hadronic spectrum contains multi-quark states whose valence contributions differ from conventional mesons and baryons. Of special interest in this regard are the so-called exotic states whose quantum numbers cannot be produced in either qqq or $q\bar{q}$ configurations, since such states represent unambiguous evidence for the presence of multi-quark physics. The question of the possible existence of multi-quark states, while of interest in its own right, is perhaps even more relevant to the desire to test and phenomenologically constrain quark-quark interactions in the QCD-inspired models (bag and potential models) which have been applied extensively to hadron spectroscopy and dynamics. This is particularly germane to efforts to elucidate constituent effects in few-body nuclear physics since qq pairs in nuclei exist in both $\bar{3}$ and 6 color configurations, in contrast to the case of the ordinary baryons which have heretofore provided the only tests and constraints of such models.

In this paper we wish to investigate the possibility of exotic, strangeness $+1$ baryons (traditionally called Z^* resonances) deeply bound with respect to available asymptotic hadron thresholds. The question of Z^* resonances has been addressed previously in the bag model¹⁻²), following Jaffe's original work on $qq\bar{q}$ states³). P -matrix⁴) analyses of existing experimental data in the $\frac{1}{2}^-$ sector show partial agreement and partial disagreement with the bag predictions⁵). In the present work we employ a QCD-like potential model, in part motivated by certain technical limitations of bag calculations. By

this we mean not so much the difficulties of the CM motion in the bag or the artificial confinement of color singlet sub-units, problems which are understood and under reasonable qualitative control^{4,6}), but rather that, 1) the boundary condition separation of inside and outside regimes in P-matrix applications of the bag represents a rather crude step-function approximation to exchange spatial matrix elements and, 2) the static spherical cavity approximation usually employed in bag model calculations neglects lower spatial symmetry configurations which are known to be favored by the color-spin structure of the quark-quark hyperfine interaction generated by one gluon exchange⁷). In making this choice we have also been influenced by Shuryak's arguments regarding the vacuum structure of QCD⁸), arguments which suggest a rather more potential-model-like than bag-like picture as the most natural phenomenological realization of QCD. Such a constituent quark picture is also supported by recent calculations using the strong coupling limit of lattice QCD⁹).

It is worth noting that potential model calculations have provided a good qualitative and, indeed, semi-quantitative understanding of the short-range features of the s-wave NN interactions¹⁰⁻¹⁵), the physics of the two channels being dominated by the repulsive character of the exchange hyperfine interaction. Since one believes the hyperfine interaction to be on reasonably firm theoretical footing, and since Z^* channels, like NN channels, are free of the presence of phenomenologically unconstrained annihilation effects, this suggests that similar calculations should provide good qualitative guidance in the Z^* sector, to the extent that effects of interest are dominated by features other than confinement.

We re-iterate at this point that our aims are primarily qualitative. We wish only to isolate those channels in which deeply bound states are expected to exist. Such states should be more accessible to experiment than those above ordinary hadronic thresholds, for which fall-apart decay modes exist, and unlike weakly bound states, whose existence or non-existence many depend on ambiguous features of the model such as effective inter-cluster mass parameters, should be unambiguously predictable. In what follows we have restricted ourselves to the negative parity, zero orbital excitation sector. While this is primarily a matter of simplicity, it allows us to ignore possible spin-orbit forces whose apparent absence in the baryon spectrum is something of a puzzle. In addition, we have included neither quark tensor force nor the mixings to p-wave color octet states analogous to those claimed to be responsible for much of the intermediate range attraction in the NN system¹⁵). This means that the masses predicted for any deeply bound states cannot be trusted to any great accuracy. Since such states are expected to have significant hidden color admixtures, the energies and mixing amplitudes of which depend strongly on unconstrained features of the confinement potential (which is, itself, presumably, only a phenomenological representation of more realistic many-body confinement effects), one does not expect to be able to predict such masses accurately, in any case.

2. Model and method

The idea behind the present calculation is that the qualitative short-range features of two-hadron physics are governed by the "diagonal" hadronic potential between the two color singlet hadrons (i.e. the induced quark exchange interaction). If this potential is repulsive in character then the system is shielded from mixing with available hidden

color excitations, (which, due to confinement, are localized at short range) and, as a result, is either unbound or at best weakly bound. In contrast, if the diagonal potential is attractive one expects significant mixing with hidden color states and hence, potentially deep binding, the degree of binding being a dynamical question and in the present context, somewhat poorly determined. The basic idea is motivated by the qualitative features of the NN system and has already been applied to the question of possible deeply bound di-baryon resonances¹⁶).

In table 1 we display all $4q\bar{s}$ states ($q=u$ or d) which can be constructed from qqq and $q\bar{s}$ clusters in spatial ground states and in a relative s -wave. The three quark clusters are taken to be totally antisymmetrized and the $4q\bar{s}$ states so constructed are coupled to the appropriate total spin, isospin, and zero net color in the standard manner.

We employ a non-relativistic QCD-inspired Hamiltonian of the form

$$H = \sum_{i=1}^5 (m_i + p_i^2/2m_i) + \sum_{i < j} (H_{conf}^{ij} + H_{hf}^{ij}) \quad (1)$$

where $H_{conf}^{ij}, H_{hf}^{ij}$ depend on $\vec{r}_{ij}, \vec{S}_i, \vec{S}_j$ and represent, respectively, the color dependent confinement and hyperfine potentials

$$H_{conf}^{ij} = \vec{F}_i \cdot \vec{F}_j \, h_{conf}(\vec{r}_{ij})$$

$$H_{hf}^{ij} = \vec{F}_i \cdot \vec{F}_j (S_i \cdot S_j \, h_{hf}(\vec{r}_{ij}) + S_{ij} \, h_{hf}(\vec{r}_{ij})) \quad (2)$$

with

$$F_i = \begin{cases} \frac{1}{2} \vec{\lambda}_i & i=1, \dots, 4 \\ -\frac{1}{2} \vec{\lambda}_i^* & i=5 \end{cases} \quad (3)$$

$$S_{ij} = 3S_i \cdot S_j \, S_{ij}/r_{ij}^2 - S_i \cdot S_j \quad (4)$$

The short-range spin-independent pieces of one gluon exchange are included with the long-range confining potential in H_{conf}^{ij} , while the first and second terms of H_{hf}^{ij} are the contact and tensor pieces of the non-relativistic reduction of one gluon exchange, respectively. Possible spin-orbit terms which do not, however, affect our calculation have been dropped from (1). We will discuss the explicit forms of h_{conf} and h_{hf} in more detail later.

Let us consider a particular channel in table 1 and denote by $|I\rangle$, $|J\rangle$, ... the partially antisymmetrized states contained therein. The quark content of the constituent clusters in $|I\rangle$ will be indicated by

$$|I(1jk;45)\rangle \quad (5)$$

where i, j, k, ℓ $\{1, 2, 3, 4\}$ are the labels of the light quarks and the index 5 refers to $\bar{s}$. One can then show, as a consequence of the partial antisymmetries of $|I\rangle$, $|J\rangle$ that for any permutationally symmetric operator such as H ,

$$\langle I_A | H | J_A \rangle = \langle I(123;45) | H(1-3\pi_{34}) | J(123;45) \rangle / N_I N_J \quad (6)$$

where π_{34} is the transposition operator $3 \leftrightarrow 4$, $|I_A\rangle$ is the normalized, fully antisymmetrized state corresponding to $|I\rangle$ and

$$N_I = \sqrt{1-3\langle I(123;45) | I(124;35) \rangle} \quad (7)$$

The states $|I_A\rangle$, $|J_A\rangle$, ... are not, in general, orthogonal; an orthogonal set, however, is readily constructed given the exchange overlap factors $\langle I_A | J_A \rangle$.

The spin, color and isospin matrix elements and overlaps required in (6) are straightforward to evaluate and we do not quote the results here. To complete the evaluation of (6) one, therefore, needs only the spatial

wavefunctions of the incompletely antisymmetrized states $|I\rangle, |J\rangle$. We chose the form

$$\phi(123;45) = \psi(R_{123;45})\phi_B(123)\phi_M(45) \quad (8)$$

where
$$\phi_B(123) = \frac{\alpha^3}{\pi^{3/2}} \exp(-\alpha^2(\rho_{123}^2 + \lambda_{123}^2)/2)$$

$$\phi_M(45) = \frac{\alpha^{3/2}}{\pi^{3/4}} \exp(-\alpha^2 r_{45}^2/2) \quad (9)$$

with
$$\rho_{123} = (\xi_1 - \xi_2)/\sqrt{2}$$

$$\lambda_{123} = (2\xi_3 - \xi_1 - \xi_2)/\sqrt{6}$$

$$\xi_{45} = \xi_4 - \xi_5$$

$$R_{123;45} = (\xi_1 + \xi_2 + \xi_3)/3 - (\xi_4 + \frac{m_s}{m} \xi_5)/(1+m_s/m). \quad (10)$$

ϕ_B, ϕ_M are the ground state cluster wavefunctions of the baryon and meson, respectively, and ψ is a variational wavefunction for the inter-cluster coordinate $R_{123;45}$, taken to be of the form

$$\psi(R) = \frac{1}{N} \sum_i \xi_i \exp(-\beta_i^2 R^2/2) \quad (11)$$

where ξ_i, β_i are variational parameters and N is a normalization factor relative to the measure $d\tau = d^3\rho_{123} d^3\lambda_{123} d^3r_{45} d^3R_{123;45}$.

3. Diagonal hadronic potentials

Given the lowest lying meson-baryon state in a given channel of table 1, the Hamiltonian (1) and the expression (6) for the Hamiltonian matrix element (with $I=J$), one may readily evaluate the "diagonal" meson-baryon exchange interaction. Effectively repulsive behavior is signalled by a depletion of the trial wavefunction (11) at short distance. In the absence of binding the system is placed in a weak harmonic box in order to ascertain the short distance behavior.

Let us define direct and exchange spatial matrix elements as follows:

$$\begin{aligned} (1j)_c^d &= \langle \phi(123;45) | h_{conf}(r_{1j}) | \phi(123;45) \rangle \\ (1j)_h^d &= \langle \phi(123;45) | h_{hfc}^{1j}(r_{1j}) | \phi(123;45) \rangle \\ k^d &= \langle \phi(123;45) | k | \phi(123;45) \rangle \\ (1j)_c^e &= \langle \phi(123;45) | h_{conf}(r_{1j}) | \phi(124;35) \rangle \\ (1j)_h^e &= \langle \phi(123;45) | h_{hfc}^{1j}(r_{1j}) | \phi(124;35) \rangle \\ k^e &= \langle \phi(123;45) | k | \phi(124;35) \rangle \\ B^e &= \langle \phi(123;45) | \phi(124;35) \rangle \end{aligned} \quad (12)$$

where the superscripts d,e label 'direct' and 'exchange,' respectively, and $K = \sum_i p_i^2/2m_i$ is the kinetic energy operator. One can readily show, from the permutational symmetries present in the matrix elements of (12), that

$$\begin{aligned} (12)_{c,h}^d &= (13)_{c,h}^d = (23)_{c,h}^d \\ (1k)_{c,h}^d &= (2k)_{c,h}^d = (3k)_{c,h}^d \quad k=4,5 \\ (13)_{c,h}^e &= (23)_{c,h}^e = (14)_{c,h}^e = (24)_{c,h}^e \\ (15)_{c,h}^e &= (25)_{c,h}^e \\ (35)_{c,h}^e &= (45)_{c,h}^e. \end{aligned} \quad (13)$$

The trial energy of the meson-baryon state in the channel in question can then be written in the general form

$$\begin{aligned} \langle I_A | H | I_A \rangle &= \left[k^d + \nu k^e - 2(12)_c^d - \frac{4}{3} (45)_c^d + \alpha_{h1}^d (12)_h^d \right. \\ &\quad + \alpha_{h2}^d (45)_h^d - \frac{2}{3} (\nu (12)_c^e + 4(13)_c^e - 2(15)_c^e - 2(34)_c^e + 4(35)_c^e) \\ &\quad \left. + \alpha_{h1}^e (12)_h^e + \alpha_{h2}^e (13)_h^e + \alpha_{h3}^e (15)_h^e + \alpha_{h4}^e (34)_h^e + \alpha_{h5}^e (35)_h^e \right] / (1+\nu B^e). \end{aligned} \quad (14)$$

The coefficients v , a_h^d and a_h^e for the lowest lying meson-baryon states of the various channels of table 1 are given in table 2. In writing (14) we have dropped an inessential constant factor $\sum_i m_i$. Note that no ground state meson-baryon configuration exists for the $I=0$ $S=5/2$ channel. Since the lowest lying $L=0$ excited meson-baryon combination lies almost a GeV above NK threshold one might hope that the hidden color state of table 1 would lie lowest of the $L=0$ states in this channel and hence be of experimental interest. We have checked this possibility and found it not to be the case — the hidden color state in fact lies well above $N\frac{3}{2}K^*$ threshold. As a result we will not discuss this channel further.

4. Results and discussion

In order to proceed we must say something about the form of h_{conf} and h_{hc}^{1j} . To begin with we employ the Isgur-Karl Hamiltonian of the baryon sector, for which

$$\begin{aligned} h_{conf}(r_{1j}) &= -\frac{1}{2} \kappa r_{1j}^2 + U(r_{1j}) \\ h_{hc}^{1j}(r_{1j}) &= -\frac{8\pi\alpha_s}{3m_1m_j} \delta^3(r_{1j}) \end{aligned} \quad (15)$$

where $U(r_{1j})$ contains departures from the harmonic limit as well as the short-range spin-independent pieces of one gluon exchange. This Hamiltonian has been applied extensively, and with considerable success, to baryon spectroscopy¹⁷ and decays¹⁸). The actual form of U is not fully determined by the baryon spectrum; we have represented it by $U(r) = C + W\delta^3(r)$, which is compatible with spectroscopic constraints. The results are, in any event, rather insensitive to the short-range piece of h_{conf} , a result of both the smearing of the potentials (15) over

clusters and the domination of short range effects by the hyperfine interaction. Note that the δ -function interaction in (15) is actually an illegal operator for the Schrödinger equation so that, as usual, the cluster parameters α , σ of (6) must be held fixed during the calculation. Spectroscopic analyses of the $SU(6)$ multiplet pattern of baryons, up to $N=2$, constrain the parameters of the model as follows:

$$\begin{aligned} \frac{2}{3} \sqrt{\frac{2}{\pi}} \frac{\alpha^3}{m^2} \alpha_s &= 260 \text{ MeV} \\ \frac{3}{2} \frac{\kappa}{\alpha^2} &= 198 \text{ MeV} \\ \frac{W\alpha^3}{(2\pi)^{3/2}} &= 176 \text{ MeV} \\ C &= 175 \text{ MeV} \end{aligned} \quad (16)$$

where $m=330$ MeV is the non-strange constituent quark mass. The cluster parameter, α , is not well-determined; the usually quoted spectroscopic value of $\alpha=320$ MeV is somewhat smaller than that favored by decay analyses. The value of σ , in this context, is adjusted to fit the $K-K^*$ mass splitting, from which we obtain $\sigma=280$ MeV. One should bear in mind here that the natural framework for potential model phenomenology is a picture of QCD in which the effective scale of chiral symmetry breaking is ~ 1 GeV¹⁹). Since this is considerably greater than the strange current quark mass one expects the kaon to be a goldstone boson of the approximate, spontaneously broken chiral $SU(3) \times SU(3)$ symmetry, and, as a result, potentially rather poorly described as a constituent $q\bar{q}$ bound state. s -wave KN scattering, however, calculated in the constituent quark framework, has proven to be in good agreement with experiment²⁰), providing a certain post facto justification for this approach. The K^* is not subject to these objections.

The results of this first phase of the calculation are as follows:

I = 0	S = 1/2	NK	weakly repulsive
I = 0	S = 3/2	NK*	attractive (bound by 11 MeV)
I = 1	S = 3/2	ΔK	attractive (bound by 9 MeV)
I = 1	S = 5/2	ΔK^*	attractive (bound by 26 MeV).

all other channels strongly repulsive. (17)

Note that the results for the two NK channels, I=0, S=1/2 and I=1 S=1/2, are in agreement with those of ref. 20), which serves as a check on our calculations. One should also note that, while $\sigma=280$ MeV reproduces the K-K* splitting, it produces masses which are too low by over 100 MeV for both states. If we instead adjust σ to give the correct K mass the splitting is too small. The effect of this change on our results however, is negligible. All qualitative conclusions are unchanged. The bindings of the last three states in (17) become 14, 13 and 30 MeV, respectively. The same is true of allowing σ to vary independently for K, K* so as to fit both masses. In addition, the binding in the I=0 S=3/2 and I=1 S=5/2 channels is stable with respect to decreasing meson and baryon cluster sizes toward the smaller values suggested by the regularized alternate Hamiltonian discussed below. The same is not true of the ΔK state of the I=1 S=3/2 channel, which is rather sensitive to such changes. Since the K is very likely a smaller object than suggested by the above treatment, we believe that this state is unlikely to exist. The other bound states, at this stage, appear to be solid candidates and, being bound already by diagonal interactions, are expected to have large hidden color mixings. However, since the evaluation of such mixings, as discussed earlier, is subject to large and unknown uncertainties, we have not attempted this calculation. It is worth noting that the I=1 S=3/2

NK* state, owing to its proximity to ΔK threshold, might be expected to mix strongly with ΔK . However, it can be shown to have a strongly repulsive diagonal interaction so that such mixing will, in fact, be highly suppressed.

While the above results are interesting, suggesting, as they do, the possible existence of at least two Z* resonances, we consider them only suggestive. The fact that the K and K* masses cannot be simultaneously fit with a common value of σ means that, with its perturbative prescription for handling the δ -function singularity of the hyperfine potential, the Isgur-Karl model is not truly applicable to mesons. This, in itself, is not particularly surprising since the hyperfine "perturbations" are rather large, especially in the meson sector. Of more importance is a recent analysis of meson spectroscopy and decays, in which the singularities of the Isgur-Karl Hamiltonian are regulated and attempts are made to simulate relativistic effects²¹⁾. It suggests that 1) the strength of the hyperfine interaction has been overestimated in the Isgur-Karl Hamiltonian and 2) the cluster sizes of K, K* mesons are smaller than those obtained above, considerably so in the case of the K. Preliminary results indicate that this modified Hamiltonian may be capable of describing the baryon spectrum, with the exception of the $\Lambda(1405)^{22)}$. Note that the considerably reduced hyperfine strength will greatly alleviate the old baryon spin-orbit problem. Since, from the point of view of ref. 21), the Isgur-Karl Hamiltonian overestimates both exchange and hyperfine effects, augmenting, as a result, precisely those effects which drive bound state formation, it provides a useful criterion for isolating channels of interest for the second phase of our calculation. In this phase, as a measure of the model sensitivity of our results, we

apply the following simplified form of the meson Hamiltonian of ref. 21) to the states of interest in (17)

$$h_{conf}(r) = -br + \frac{a_s(r)}{r} + C$$

$$h_{hf}^{1/2}(r) = \frac{-8\pi}{3m_1 m_j} \sum_k a_k \frac{\gamma_k^3}{\pi^{3/2}} \exp(-\gamma_k^2 r^2) \quad (18)$$

$$a_s(r) = \sum_k a_k \text{Erf}(\gamma_k r) \quad (19)$$

The parameters of a_k , γ_k are meant to incorporate the one-loop asymptotic freedom result for the variation of the strong coupling constant, $a_s(Q^2)$, with Q^2 , supplemented by a smoothly implemented saturation at low Q^2 . The parameters of the model are

$$\begin{aligned} a_1 &= .25 & a_2 &= .15 & a_3 &= .2 \\ \gamma_1 &= .5 \text{ GeV} & \gamma_2 &= \sqrt{2.5} \text{ GeV} & \gamma_3 &= \sqrt{250} \text{ GeV} \\ b &= .135 \text{ GeV}^2 \\ C &= 580 \text{ MeV} \end{aligned} \quad (20)$$

a_N , a_Δ , a_K and a_{K^*} are then obtained by optimizing the N, Δ , K and K^* energies. The single Gaussian approximation provides a good description of the K^* but is less successful for the N, Δ , and K. The latter, in particular, is known to require short distance enhancements to account for the effects of the strong hyperfine interaction. Since, however, the K charge radius is in good agreement with that of the full K wavefunction, and since the various exchange effects are primarily sensitive to the tails of the hadronic wavefunctions, we believe this approximation to be adequate for our purposes. The resulting cluster parameters are

$$\begin{aligned} a_N &= 420 \text{ MeV} \\ a_\Delta &= 375 \text{ MeV} \\ a_K &= 515 \text{ MeV} \\ a_{K^*} &= 320 \text{ MeV} \end{aligned} \quad (21)$$

Applying the Hamiltonian (18) to the channels of interest we find that the bound states of (17) are no longer bound, although the I=1 S=5/2 ΔK^* diagonal interaction remains attractive in character. We have further checked this result by allowing the lowest excited state of each channel to mix with the ground state. The I=0 S=3/2 and I=1 S=3/2 channels remain unbound, while the I=1 S=5/2 channel is now bound by 56 MeV relative to the K^* threshold.

5. Conclusions

In conclusion, we believe the existence of an I=1 S=5/2 Z^* resonance is strongly indicated by our results. Note that, since the strength of the quark tensor interaction is almost an order of magnitude less than that of the hyperfine interaction, one expects, as was the case for di-baryons, a small tensor decay width for this state. As a result, NK elastic scattering will be disfavored as a means of detection, and one should instead look in processes such as $N\pi \rightarrow X\bar{K}$ or $NK \rightarrow X\pi$ where polarization measurements might allow one to identify an I=1 $J^P = 5^-$ state recoiling against $\bar{K}$ or π . The I=0 S=3/2 state is considerably less strongly indicated. To the extent that the Hamiltonian (18) seems to us a more reasonable one, we believe it is likely that this state does not exist, at least as an NK^* bound state, although it is possible that a residue of the attraction may persist above threshold, making a coupled channel NK, NK^* scattering calculation of some interest in this

channel. Finally, the $I=1$ $S=3/2$ ΔK state suggested by the Isgur-Karl Hamiltonian seems to us, almost certainly, an artifact of the approximations of the model. The repulsive character of both ΔK , NK^* diagonal interactions induced by the Hamiltonian (18) also make it unlikely that any resonant enhancement will be present above threshold.

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Table 1: Available States Constructed From qqq and $q\bar{s}$ Clusters In Spatial Gound States and In a Relative s-wave*

S	1/2			3/2			5/2		
	I			I			I		
2	ΔK^*			ΔK			ΔK^*		
	$\Delta_c(1/2)K_c$			ΔK^*					
	$\Delta_c(1/2)K_c^*$			$\Delta_c(1/2)K_c^*$					
1	NK			ΔK			ΔK^*		
	NK*			NK*			$N_c(3/2)K_c^*$		
	ΔK^*			ΔK^*					
	$N_c(1/2)K_c$			$N_c(3/2)K_c$					
	$N_c(1/2)K_c^*$			$N_c(1/2)K_c^*$					
0	$N_c(3/2)K_c^*$			$N_c(3/2)K_c^*$					
	$\Delta_c(1/2)K_c$			$\Delta_c(1/2)K_c^*$					
	$\Delta_c(1/2)K_c^*$								
	NK			NK*			$N_c(3/2)K_c^*$		
0	NK*			$N_c(3/2)K_c$					
	$N_c(1/2)K_c$			$N_c(1/2)K_c^*$					
	$N_c(1/2)K_c^*$			$N_c(3/2)K_c^*$					
	$N_c(3/2)K_c^*$								

*The subscript 'c' denotes a cluster in a color octet configuration. All states are coupled to the appropriate value of I, S and to zero net color.

Table 2: Coefficients for Hamiltonian Matrix Elements of Lowest Lying Meson-Baryon States

Channel I S	lowest state	ν	α_{n1}^d	α_{n2}^d	α_{n1}^e	α_{n2}^e	α_{n3}^e	α_{n4}^e	α_{n5}^e
0 1/2	NK	0	1/2	1	1/12	-1/6	-1/12	0	0
0 3/2	NK*	1/3	1/2	-1/3	5/18	-2/9	0	5/54	-13/27
0 5/2	--	--	--	--	--	--	--	--	--
1 1/2	NK	-1/3	1/2	1	-1/18	-1/9	-1/18	-1/3	-2/3
1 3/2	ΔK	1/6	-1/2	1	-1/36	-1/9	-1/18	1/6	1/3
1 5/2	ΔK^*	1/3	-1/2	-1/3	-1/18	-2/9	1/9	1/9	-2/9
2 1/2	ΔK^*	--	-1/2	-1/3	-1/18	-2/9	-1/3	-7/9	-2/9
2 3/2	ΔK	-1/2	-1/2	1	1/12	1/3	1/6	-1/2	-1
2 5/2	ΔK^*	-1	-1/2	-1/3	1/6	2/3	-1/3	-1/3	2/3