part motivated by certain technical limitations of our calculations. By

In the present work we employ a R-ML-like potential model, in

show partial experimental and partial experimental data in the pro-

model (-2) + 0.19415, whereas data on 1H/1H experiments (model (-2))

and the data on 2H + 0.19415, whereas data on 1H/1H experiments (model (-2))

neutron or Z resonance in the resonance regime. The

question of the resonance has been addressed recently by the

Eisenberg + 1 H + 0.19415 (traditional) and 2H resonance) (deep)

In this paper we wish to investigate the possibilities of exact

models

which have heretofore provided the only tests and constraints on such
color confinement* in a manner consistent with the case of the ordinary baryons.

Both nuclear physics since gg pairs in nuclear matter in such a
density perturbatively Berman to effects to elucidate confinement effects in real,

particular interest to effects to elucidate confinement effects in real,

This is a

spectrum of instantons to hadron spectroscopy and dynamics. This is

actions in the QCD-limited model (QCD and potentential models) which have

action to the extent to limit and phenoamnologicaly constrain number of

tweets, while of interest in the current, can these even more revealing

gauge physics. The question of the possible existence of multi-gluon

such states represent nonperturbative evidence for the presence of multi-gluon

number cannot be produced in either gg or QCD confinement, since

interactions in this regard are the so-called exact states whose quantum

intrications differ from confinement masses and spectra of spectral

physical hadron spectroscopy contain multi-gluon states whose existence

continue (SU(3) hexagon) color degree of freedom leads to the possibility that

The motion that ordinary hadrons are composed of quarks carrying a

1. Introduction

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this we mean not so much the difficulties of the CM motion in the bag or the artificial confinement of color singlet subunits, problems which are understood and under reasonable qualitative control\textsuperscript{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 exchange\textsuperscript{7}). In making this choice we have also been influenced by Shuryak's arguments regarding the vacuum structure of QCD\textsuperscript{8}), 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 QCD\textsuperscript{9}).

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 interactions\textsuperscript{10-15}), 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\textsuperscript{8} 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\textsuperscript{8} 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 system\textsuperscript{15}). 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
To complete the evaluation of (6), one therefore needs only the pattern
of the spin, color and flavor matrix elements and overlap functions.

The traces \( \langle \bar{V}_a \bar{V}_b \rangle \) exist, and are not, in general, orthogonal to an
orthogonal set. However, the trace of the contracted gluon exchange
operator is zero, and so the traces \( \langle \bar{V}_a \bar{V}_b \rangle \) are not, in general, orthogonal to

\[
N \langle \bar{V}_a \bar{V}_b \rangle = N \langle \bar{V}_a \bar{V}_b \rangle
\]

with

\[
\langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle + \langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle
\]

operator such as \( N \langle \bar{V}_a \bar{V}_b \rangle \).

The partial contraction satisfies the symmetric
property of (13), i.e., for any permutations \( \{1,2,3,4\} \),

\[
\langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle
\]

quick contraction of the contravariant indices in \( \bar{V}_a \) will be handled by
the partial contraction property contained therein. The
\( \langle \bar{V}_a \bar{V}_b \rangle \) are the partial contraction of the partial contraction function.

In more detail, let

\[
\langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle
\]

dropped from (13). We will disregard the explicit forms of \( \bar{V}_a \) and \( \bar{V}_b \).

The basic idea is to extract the quantities of interest from the
explicit forms. Once this has been done, however, our calculations have been
non-relativistic reduction of one gluon exchange, respectively, possible
extract and second terms of the color and flavor matrix elements of the
functions included with the long-range contractions potential in \( \bar{V}_a \) and \( \bar{V}_b \).

\[
N \langle \bar{V}_a \bar{V}_b \rangle = N \langle \bar{V}_a \bar{V}_b \rangle
\]

with

\[
\langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle + \langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle
\]

depend on contravariant and covariant potentials,
where \( \lambda_0 \) depends on \( \lambda_0 \) and \( \lambda_0 \) and \( \lambda_0 \) respectively, the color

\[
N \langle \bar{V}_a \bar{V}_b \rangle = N \langle \bar{V}_a \bar{V}_b \rangle + N \langle \bar{V}_a \bar{V}_b \rangle = N \langle \bar{V}_a \bar{V}_b \rangle
\]

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

\[
\langle \bar{V}_a \bar{V}_b \rangle + \langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle + \langle \bar{V}_a \bar{V}_b \rangle = \langle \bar{V}_a \bar{V}_b \rangle
\]

ment.

The basic idea is to extract the quantities of interest from the
explicit forms. Once this has been done, however, our calculations have been
non-relativistic reduction of one gluon exchange, respectively, possible
extract and second terms of the color and flavor matrix elements of the
functions included with the long-range contractions potential in \( \bar{V}_a \) and \( \bar{V}_b \).

\[
N \langle \bar{V}_a \bar{V}_b \rangle = N \langle \bar{V}_a \bar{V}_b \rangle + N \langle \bar{V}_a \bar{V}_b \rangle = N \langle \bar{V}_a \bar{V}_b \rangle
\]

the color exchange potentials, which, due to contravariant, are localized at short
range (and, as a result, are either unbound or at best weakly bound).
wavefunctions of the incompletely antisymmetrized states \(|I\rangle, |J\rangle\). We chose the form

$$\Phi(123;45) = \Phi(B_{123;45}) \Phi(B_{123}) \Phi(N_{45}) \quad (8)$$

where

$$\Phi(B_{123}) = \frac{a^3}{2^{7/2}} \exp(-a^2 (\rho_{123^2} + \lambda_{123^2})/2)$$
$$\Phi(N_{45}) = \frac{a^{3/2}}{2^{3/2}} \exp(-a^2 r_{45}^2/2) \quad (9)$$

with

$$E_{123} = (E_1 + E_2)/2$$
$$\lambda_{123} = (2E_3 - E_1 - E_2)/6$$
$$E_{45} = E_4 + E_5$$

$$R_{123;45} = (E_1 + E_2 + E_3)/3 - (E_4 + E_5)/(1 + \frac{m_a}{m_a} E_5). \quad (10)$$

$$\Phi_B, \Phi_M$$ are the ground state cluster wavefunctions of the baryon and meson, respectively, and \(\Phi\) is a variational wavefunction for the inter-cluster coordinate \(B_{123;45}\), taken to be of the form

$$\Phi(R) = \frac{1}{N} \sum \xi_{45} \exp(-\xi_{45}^2 R^2/2) \quad (11)$$

where \(\xi_{45}\) are variational parameters and \(N\) is a normalization factor relative to the measure \(d\tau = d^4p_{123}d^3r_{45}d^3R_{123;45}\).

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:

$$\langle 1j \rangle^c_k = \langle 1j \rangle^d_k = \langle 1j \rangle^e_k = \langle 1j \rangle^d_k = \langle 1j \rangle^e_k = \langle 1j \rangle^d_k = \langle 1j \rangle^e_k = \langle 1j \rangle^d_k = \langle 1j \rangle^e_k \rangle$$

where the superscripts \(d, e\) label 'direct' and 'exchange,' respectively, and \(K = \sum_1 p_i^2/2m_i\) is the kinetic energy operator. One can readily show, from the permutational symmetries present in the matrix elements of (12), that

$$\langle 12 \rangle^d_{c,h} = \langle 13 \rangle^d_{c,h} = \langle 23 \rangle^d_{c,h}$$
$$\langle 1k \rangle^d_{c,h} = \langle 2k \rangle^d_{c,h} = \langle 3k \rangle^d_{c,h} \quad k=4,5$$
$$\langle 13 \rangle^e_{c,h} = \langle 23 \rangle^e_{c,h} = \langle 14 \rangle^e_{c,h} = \langle 24 \rangle^e_{c,h}$$
$$\langle 15 \rangle^e_{c,h} = \langle 25 \rangle^e_{c,h}$$
$$\langle 35 \rangle^e_{c,h} = \langle 45 \rangle^e_{c,h} \quad (13)$$

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

$$\langle E_A | H | E_A \rangle = \left[ a_{h1}^d (12)_{c,h}^d - \frac{6}{3} (45)_{c,h}^d + \frac{4}{3} (12)_{c,h}^d \right]$$
$$+ a_{h2}^d (45)_{h}^d + \frac{2}{3} (v_{h1} (12)_{c,h}^e (13)_{c,h}^e - (15)_{c,h}^e (34)_{c,h}^e + (34)_{c,h}^e (35)_{c,h}^e)$$
$$+ a_{h3}^d (12)_{h}^e (13)_{h}^e (15)_{h}^e (34)_{h}^e (35)_{h}^e \right]/(1 + \nu_{h}^e). \quad (14)$$
subject to these objections.

Certain parts of this paper deal with certain aspects of the potential

function, as shown in Fig. 1, which are not discussed in detail. However, the

question of the potential function is dealt with in more detail in the

concluding remarks. If you are interested in further details, you are

referred to the conclusion of this paper.

The potential function is also discussed in some detail in the

concluding remarks. The conclusions are based on the potential function

and the conclusions suggested in this paper are due to the potential

function.

The conclusions are based on the potential function and the conclusions

suggested in this paper are due to the potential function.
The results of this first phase of the calculation are as follows:

<table>
<thead>
<tr>
<th>I = 0</th>
<th>S = 1/2</th>
<th>NK</th>
<th>weakly repulsive</th>
</tr>
</thead>
<tbody>
<tr>
<td>I = 0</td>
<td>S = 3/2</td>
<td>NK*</td>
<td>attractive (bound by 11 MeV)</td>
</tr>
<tr>
<td>I = 1</td>
<td>S = 3/2</td>
<td>AK</td>
<td>attractive (bound by 9 MeV)</td>
</tr>
<tr>
<td>I = 1</td>
<td>S = 5/2</td>
<td>AK*</td>
<td>attractive (bound by 26 MeV).</td>
</tr>
</tbody>
</table>

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 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 $\sigma$ 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 $\sigma$ 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 AK 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 $\Delta K$ threshold, might be expected to mix strongly with $\Delta K$. 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 $\sigma$ means that, with its perturbative prescription for handling the 6-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 the A(1405)22. Note that the considerably reduced hyperfine strength will readily 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
channel XX. "N" scattering calculation of some interest in this
context. The scattering of the antiferromagnetic phase above the threshold, making a coupled
not, external to at least an "N" bound state, although it is possible that
not externally due to the existence of an "N" bound state, it is highly unlikely at this level of
descriptions to be limited to the existence that the Hamiltonian (19) becomes

\[ \frac{Z}{X} \approx \sqrt{5/2} \approx \text{spin-flip state is completely depleted} \]

The agreement of this result with the exact value is

\[ \frac{Z}{X} = \frac{1}{2} \approx \text{spin-flip state is completely depleted} \]

should be noted that in the approximation made here, the

XX and YY scattering will be determined as a mean of excitation, and one

for scattering, a small tensor decay width for this reason, as a result,

less than that of the nonperturbative interaction, one expects, as was the case.

The strength of the nonperturbative interaction is shown in another]

\[ Z = \frac{1}{2} \approx \text{spin-flip state is completely depleted} \]

in comparison, as predicted by our results. Note then, the

5. Conclusions

5.5. New result to the XX threshold:

the channel XX channel is now bound to the XX channel to mix with the XX channel. The model suggests, therefore, that

the bound states of (17) are no longer bound, although the XX channel to interest in this

applying the Hamiltonian (18) to the channels of interest, we find

\[ (12) \]

\[ A_X = \frac{1}{2}, A = 0 \]

\[ A_X = \frac{1}{2}, A = 0 \]

\[ A_X = \frac{1}{2}, A = 0 \]

\[ A_X = \frac{1}{2}, A = 0 \]

\[ A_X = \frac{1}{2}, A = 0 \]

\[ A_X = \frac{1}{2}, A = 0 \]

\[ A_X = \frac{1}{2}, A = 0 \]
channel. Finally, the $I=1$ $S=3/2$ $\Delta K$ 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 $\Delta K$, $N K^*$ diagonal interactions induced by the Hamiltonian (18) also make it unlikely that any resonant enhancement will be present above threshold.

Acknowledgements

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References

Table I: Available States Constructed From gg and gp Clusters In
Spatial Condensate and a Relative-Swave

<table>
<thead>
<tr>
<th>s</th>
<th>I/2</th>
<th>3/2</th>
<th>1/2</th>
<th>1</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>N^+</td>
<td>N^+</td>
<td>N^+</td>
<td>N^+</td>
</tr>
<tr>
<td>1</td>
<td>N^+</td>
<td>N^+</td>
<td>N^+</td>
<td>N^+</td>
</tr>
<tr>
<td>2</td>
<td>N^+</td>
<td>N^+</td>
<td>N^+</td>
<td>N^+</td>
</tr>
</tbody>
</table>

---

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

<table>
<thead>
<tr>
<th>Channel</th>
<th>lowest state</th>
<th>( v )</th>
<th>( a_{n1}^d )</th>
<th>( a_{n2}^d )</th>
<th>( a_{n1}^e )</th>
<th>( a_{n2}^e )</th>
<th>( a_{n3}^e )</th>
<th>( a_{n4}^e )</th>
<th>( a_{n5}^e )</th>
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</thead>
<tbody>
<tr>
<td>0 1/2</td>
<td>NK</td>
<td>0</td>
<td>1/2</td>
<td>1</td>
<td>1/12</td>
<td>-1/6</td>
<td>-1/12</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td>0 3/2</td>
<td>NK*</td>
<td>1/3</td>
<td>1/2</td>
<td>-1/3</td>
<td>5/18</td>
<td>-2/9</td>
<td>0</td>
<td>5/54</td>
<td>-13/27</td>
</tr>
<tr>
<td>0 5/2</td>
<td>--</td>
<td>--</td>
<td>--</td>
<td>--</td>
<td>--</td>
<td>--</td>
<td>--</td>
<td>--</td>
<td>--</td>
</tr>
<tr>
<td>1 1/2</td>
<td>NK</td>
<td>-1/3</td>
<td>1/2</td>
<td>1</td>
<td>-1/18</td>
<td>-1/9</td>
<td>-1/18</td>
<td>-1/3</td>
<td>-2/3</td>
</tr>
<tr>
<td>1 3/2</td>
<td>( \Delta K )</td>
<td>1/6</td>
<td>-1/2</td>
<td>1</td>
<td>-1/36</td>
<td>-1/9</td>
<td>-1/18</td>
<td>1/6</td>
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<tr>
<td>1 5/2</td>
<td>( \Delta K^* )</td>
<td>1/3</td>
<td>-1/2</td>
<td>-1/3</td>
<td>-1/18</td>
<td>-2/9</td>
<td>1/9</td>
<td>1/9</td>
<td>-2/9</td>
</tr>
<tr>
<td>2 1/2</td>
<td>( \Delta K^* )</td>
<td>--</td>
<td>--</td>
<td>-1/2</td>
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<td>-2/9</td>
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<tr>
<td>2 3/2</td>
<td>( \Delta K )</td>
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<td>-1/2</td>
<td>1</td>
<td>1/12</td>
<td>1/3</td>
<td>1/6</td>
<td>-1/2</td>
<td>-1</td>
</tr>
<tr>
<td>2 5/2</td>
<td>( \Delta K^* )</td>
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<td>-1/2</td>
<td>-1/3</td>
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<td>2/3</td>
<td>-1/3</td>
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