BARYONS IN CHIRAL CONSTITUENT QUARK MODEL

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Abstract

Beyond the spontaneous chiral symmetry breaking scale light and strange baryons should be considered as systems of three constituent quarks with an effective confining interaction and a flavor-spin chiral interaction that is mediated by the octet of Goldstone bosons (pseudoscalar mesons) between the constituent quarks. One cannot exclude, however, the possibility that this flavor-spin interaction has an appreciable vector- and higher meson exchange component.

1 Introduction

Our aim in physics is not only to calculate some observable and get a correct number but mainly to understand a physical picture responsible for the given phenomenon. It very often happens that a theory formulated in terms of fundamental degrees of freedom cannot answer such a question since it becomes overcomplicated at the related scale. Thus a main task in this case is to select those degrees of freedom which are indeed essential. For instance, the fundamental degrees of freedom in crystals are ions in the lattice, electrons and the electromagnetic field. Nevertheless, in order to understand electric conductivity, heat capacity, etc. we instead work with "heavy electrons" with dynamical mass, phonons and their interaction. In this case a complicated electromagnetic interaction of the electrons with the ions in the lattice is "hidden" in the dynamical mass of the electron and the interactions among ions in the lattice are eventually responsible for the collective excitations of the lattice - phonons, which are Goldstone bosons of the spontaneously broken translational invariance in the lattice of ions. As a result, the theory becomes rather simple - only the electron and phonon degrees of freedom and their interactions are essential for all the properties of crystals mentioned above.

Quite a similar situation takes place in QCD. One hopes that sooner or later one can solve the full nonquenched QCD on the lattice and get the correct nucleon and pion mass in terms of underlying degrees of freedom: current quarks and gluon fields. However, QCD at the scale of 1 GeV becomes too complicated, and hence it is rather difficult to say in this case what kind of physics, inherent in QCD, is relevant to the nucleon mass and its low-energy properties. I will show that it is the spontaneous breaking of chiral symmetry which is the most important QCD phenomenon in this case, and that beyond the scale of spontaneous breaking of chiral symmetry light and
strange baryons can be viewed as systems of three constituent quarks which interact
by the exchange of Goldstone bosons (pseudoscaler mesons) and are subject to con-
finement.

2 Spontaneous Chiral Symmetry Breaking and its Consequences for Low-
Energy QCD

The QCD Lagrangian with three light flavors has a global symmetry

\[ SU(3)_L \times SU(3)_R \times U(1)_V \times U(1)_A, \] (1)

if one neglects the masses of current u,d, and s quarks, which are small compared to a
typical low-energy QCD scale of 1 GeV. The \( U(1)_A \) is not a symmetry at the quantum
level due to the axial anomaly. If the \( SU(3)_L \times SU(3)_R \) chiral symmetry of the QCD
Lagrangian were intact in the vacuum state we would observe degenerate multiplets in
the particle spectrum corresponding to the above chiral group, and all hadrons would
have their degenerate partners with opposite parity. Since this does not happen the
implication is that the chiral symmetry is spontaneously broken down to \( SU(3)_V \) in
the QCD vacuum, i.e., realized in the hidden Nambu-Goldstone mode. A direct evi-
dence for the spontaneously broken chiral symmetry is a nonzero value of the quark
condensates for the light flavors \(< \text{vacuum} | \bar{q}q | \text{vacuum} > \approx -(240 \text{ – } 250 \text{MeV})^3\), which
represents the order parameter. That this is indeed so, we know from three independ-
ent sources: current algebra, QCD sum rules, and lattice gauge calculations. There are two important generic consequences of the spontaneous chiral sym-
metry breaking. The first one is an appearance of the octet of pseudoscalar mesons
of low mass, \( \pi, K, \eta \), which represent the associated approximate Goldstone bosons.
The second one is that valence quarks acquire a dynamical or constituent mass. Both
these consequences of the spontaneous chiral symmetry breaking are well illustrated
by, e.g. the \( \sigma \)-model or the Nambu and Jona-Lasinio model. We cannot say
at the moment for sure what the microscopical reason for spontaneous chiral sym-
metry breaking in the QCD vacuum is. It was suggested that this occurs when quarks
propagate through instantons in the QCD vacuum.

For the low-energy baryon properties it is only essential that beyond the sponta-
neous chiral symmetry breaking scale new dynamical degrees of freedom appear -
constituent quarks and chiral fields. The low-energy baryon properties are mainly
determined by these dynamical degrees of freedom and the confining interaction.

3 The Chiral Boson Exchange Interaction

In an effective chiral description of the baryon structure, based on the con-
stituent quark model, the coupling of the quarks and the pseudoscalar Goldstone
bosons will (in the \( SU(3)_F \) symmetric approximation) have the form \( ig \bar{\psi} \gamma_5 \vec{\lambda} \cdot \vec{\phi} \psi \) (or
\( g/(2m) \bar{\psi} \gamma_\mu \gamma_5 \vec{\lambda} \psi \partial^\mu \vec{\phi} \)), where \( \psi \) is the fermion constituent quark field operator, \( \vec{\phi} \)
the octet boson field operator, and \( g \) is a coupling constant. A coupling of this form, in a nonrelativistic reduction for the constituent quark spinors, will – to lowest order – give rise to a Yukawa interaction between the constituent quarks, the spin-spin component of which has the form

\[
V_Y(r_{ij}) = \frac{g^2}{4\pi^2} \frac{1}{4m_i m_j} \bar{\sigma}_i \cdot \bar{\sigma}_j \hat{\lambda}_i^F \cdot \hat{\lambda}_j^F \left\{ \mu^2 e^{-\mu r_{ij}} - 4\pi \delta(r_{ij}) \right\}.
\]

Here \( m_i \) and \( m_j \) denote the masses of the interacting quarks, and \( \mu \) that of the meson. There will also be an associated tensor component, which is discussed in ref. [2].

At short range the simple form (2) of the chiral boson exchange interaction cannot be expected to be realistic and should only be taken to be suggestive. Because of the finite spatial extent of both the constituent quarks and the pseudoscalar mesons the delta function in (2) should be replaced by a finite function, with a range of 0.6-0.7 fm, as suggested by the spatial extent of the mesons. In addition, the radial behaviour of the Yukawa potential (2) is valid only if the boson field satisfies a linear Klein-Gordon equation. The implications of the underlying chiral symmetry of QCD for the effective chiral Lagrangian (which in fact is not known), which contains constituent quarks as well as boson fields, are that these boson fields cannot be described by linear equations near their source. For a clarification on this important issue see [14].

At this stage the proper procedure should be to avoid further specific assumptions about the short range behavior of \( V(r) \) in (3), to extract instead the required matrix elements of it from the baryon spectrum, and to reconstruct by this an approximate radial form of \( V(r) \). The overall minus sign in the effective chiral boson interaction in (3) corresponds to that of the short range term in the Yukawa interaction. It is the latter that is of crucial importance in baryon physics.

The flavor structure of the pseudoscalar octet exchange interaction between two quarks \( i \) and \( j \) should be understood as follows:

\[
- V(r_{ij}) \hat{\lambda}^F_i \cdot \hat{\lambda}^F_j \bar{\sigma}_i \cdot \bar{\sigma}_j = \left( \sum_{a=1}^3 V_{\pi}(r_{ij}) \lambda^a_i \lambda^a_j + \sum_{a=4}^7 V_K(r_{ij}) \lambda^a_i \lambda^a_j + V_\eta(r_{ij}) \lambda^8_i \lambda^8_j \right) \bar{\sigma}_i \cdot \bar{\sigma}_j.
\]

The first term in (3) represents the pion-exchange interaction, which acts only between light quarks. The second term represents the Kaon exchange interaction, which takes place in u-s and d-s pair states. The \( \eta \)-exchange, which is represented by the third term, is allowed in all quark pair states.

The interaction (3) should be contrasted with the gluon-exchange one [15], which has been used in numerous earlier attempts to describe the baryon spectra with the constituent quark model (see e.g. [16]). The gluon-exchange model fails to explain the following outstanding features of baryon spectroscopy: (i) the different ordering of lowest positive and negative parity states in the spectra of nucleon and \( \Lambda \)-hyperon (the gluon exchange is sensitive to the spin and colour degrees of freedom of quarks only, so the \( N \) and \( \Lambda \) spectra should be very similar as they differ only in their flavor
content); (ii) the fact that some of the two-quantum excitations of positive parity in all spectra (e.g. \(N(1440), \Delta(1600), \Lambda(1600), \Sigma(1660)\)) lie below the one-quantum excitations of negative parity; (iii) an absence in the empirical spectra of the strong spin-orbit splittings implied by the gluon-exchange interaction. The gluon-exchange model is self-contradictory in principle: an introduction into the theory of the constituent quark mass (to be contrasted with the small current quark mass) implies that the underlying chiral symmetry of QCD is spontaneously broken. If so, according to the Goldstone theorem there must be Goldstone bosons which should participate in the baryon structure on the same footing as constituent quarks.

4 The Structure of the Baryon Spectrum

The two-quark matrix elements of the interaction (3) are:

\[
\langle [f_{ij}]_F \times [f_{ij}]_S : [f_{ij}]_{FS} \mid -V(r_{ij}) \vec{\lambda}_i^F \cdot \vec{\lambda}_j^F \vec{\sigma}_i \cdot \vec{\sigma}_j \mid [f_{ij}]_F \times [f_{ij}]_S : [f_{ij}]_{FS} \rangle = \begin{cases} \\
\frac{4}{3} V(r_{ij}) [2]_F, [2]_S : [2]_{FS} \\
-8 V(r_{ij}) [11]_F, [11]_S : [2]_{FS} \\
4 V(r_{ij}) [2]_F, [11]_S : [11]_{FS} \\
\frac{8}{3} V(r_{ij}) [11]_F, [2]_S : [11]_{FS} \\
\end{cases}
\]

(4)

From these the following important properties may be inferred:

(i) At short range, where \(V(r_{ij})\) is positive, the chiral interaction (3) is attractive in the symmetric FS pairs and repulsive in the antisymmetric ones. At large distances the potential function \(V(r_{ij})\) becomes negative and the situation is reversed.

(ii) At short range, among the FS-symmetrical pairs, the flavor antisymmetric pairs experience a much larger attractive interaction than the flavor-symmetric ones, and among the FS-antisymmetric pairs the strength of the repulsion in flavor-antisymmetric pairs is considerably weaker than in the symmetric ones.

Given these properties we conclude, that with the given flavor symmetry, the more symmetrical the FS Young pattern is for a baryon the more attractive contribution at short range comes from the interaction (3). For two identical flavor-spin Young patterns \([f]_FS\) the attractive contribution at short range is larger for the more antisymmetrical flavor Young pattern \([f]_F\).

Consider first, for the purposes of illustration, a schematic model which neglects the radial dependence of the potential function \(V(r)\) in (3), and assume a harmonic confinement among quarks as well as \(m_u = m_d = m_s\). In this model

\[
H_\chi \sim - \sum_{i<j} C_\chi \vec{\lambda}_i^F \cdot \vec{\lambda}_j^F \vec{\sigma}_i \cdot \vec{\sigma}_j.
\]

(5)

If the only interaction between the quarks were the flavor- and spin-independent harmonic confining interaction, the baryon spectrum would be organized in multiplets.
of the symmetry group $SU(6)_{FS} \times U(6)_{conf}$. In this case the baryon masses would be determined solely by the orbital structure, and the spectrum would be organized in an alternative sequence of positive and negative parity states. The Hamiltonian (5), within a first order perturbation theory, reduces the $SU(6)_{FS} \times U(6)_{conf}$ symmetry down to $SU(3)_F \times SU(2)_S \times U(6)_{conf}$, which automatically implies a splitting between the octet and decuplet baryons.

For the octet states $N$, $\Lambda$, $\Sigma$, $\Xi$ ($N = 0$ shell, $N$ is the number of harmonic oscillator excitations in a 3-quark state) as well as for their first radial excitations of positive parity $N(1440)$, $\Lambda(1600)$, $\Sigma(1660)$, $\Xi(?)$ ($N = 2$ shell) the flavor and spin symmetries are $[3]_{FS}[21]_F[21]_S$, and the contribution of the Hamiltonian (5) is $-14C_\chi$. For the decuplet states $\Delta$, $\Sigma(1385)$, $\Xi(1530)$, $\Omega$ ($N = 0$ shell) the flavor and spin symmetries, as well as the corresponding matrix element, are $[3]_{FS}[3]_F[3]_S$ and $-4C_\chi$, respectively. The first negative parity excitations ($N = 1$ shell) in the $N$ and $\Sigma$ spectra $N(1535) - N(1520)$ and $\Sigma(1750) - \Sigma(?)$ are described by the $[21]_{FS}[21]_F[21]_S$ symmetries, and the contribution of the interaction (5) in this case is $-2C_\chi$. The first negative parity excitation in the $\Lambda$ spectrum ($N = 1$ shell) $\Lambda(1405) - \Lambda(1520)$ is flavor singlet $[21]_{FS}[11]_F[21]_S$, and, in this case, the corresponding matrix element is $-8C_\chi$.

These matrix elements alone suffice to prove that the ordering of the lowest positive and negative parity states in the baryon spectrum will be correctly predicted by the chiral boson exchange interaction (5). The constant $C_\chi$ may be determined from the $N-\Delta$ splitting to be 29.3 MeV. The oscillator parameter $\hbar\omega$, which characterizes the effective confining interaction, may be determined as one half of the mass differences between the first excited $\frac{1}{2}^+$ states and the ground states of the baryons, which have the same flavor-spin, flavor and spin symmetries (e.g. $N(1440)$ - $N$, $\Lambda(1600)$ - $\Lambda$, $\Sigma(1660)$ - $\Sigma$), to be $\hbar\omega \approx 250$ MeV. Thus the two free parameters of this simple model are fixed and we can make now predictions. In the $N$ and $\Sigma$ sectors the mass difference between the lowest excited $\frac{1}{2}^+$ states ($N(1440)$ and $\Sigma(1660)$) and $\frac{1}{2}^- - \frac{3}{2}^-$ negative parity pairs ($N(1535) - N(1520)$ and $\Sigma(1750) - \Sigma(?)$) will then be

$$N, \Sigma: \quad m(\frac{1}{2}^+)^N - m(\frac{1}{2}^- - \frac{3}{2}^-) = 250 \text{ MeV} - C_\chi(14 - 2) = -102 \text{ MeV},$$

whereas for the $\Lambda$ system ($\Lambda(1600)$, $\Lambda(1405)$ - $\Lambda(1520)$) it should be

$$\Lambda: \quad m(\frac{1}{2}^+)^\Lambda - m(\frac{1}{2}^- - \frac{3}{2}^-) = 250 \text{ MeV} - C_\chi(14 - 8) = 74 \text{ MeV}.$$

This simple example shows how the chiral interaction (5) provides different ordering of the lowest positive and negative parity excited states in the spectra of the nucleon and the $\Lambda$-hyperon. This is a direct consequence of the symmetry properties of the boson-exchange interaction discussed at the beginning of this section. Namely, the $[3]_{FS}$ state in the $N(1440)$, $\Delta(1600)$ and $\Sigma(1660)$ positive parity resonances from the $N = 2$ band feels a much stronger attractive interaction than the mixed symmetry state $[21]_{FS}$ in the $N(1535) - N(1520)$, $\Delta(1620) - \Delta(1700)$ and $\Sigma(1750) - \Sigma(?)$
resonances of negative parity \((N = 1\) shell). Consequently the masses of the positive parity states \(N(1440), \Delta(1600)\) and \(\Sigma(1660)\) are shifted down relative to the other ones, which explains the reversal of the otherwise expected "normal ordering". The situation is different for \(\Lambda(1405) - \Lambda(1520)\) and \(\Lambda(1600)\), as the flavor state of \(\Lambda(1405) - \Lambda(1520)\) is totally antisymmetric. Because of this the \(\Lambda(1405) - \Lambda(1520)\) gains an attractive energy, which is comparable to that of the \(\Lambda(1600)\), and thus the ordering suggested by the confining oscillator interaction is maintained.

Consider now, in addition, the radial dependence of the potential with the \(SU(3)\) invariant version \((3)\) of the chiral boson exchange interaction (i.e., \(V_\pi(r) = V_K(r) = V_\eta(r)\)). If the confining interaction in each quark pair is taken to have the harmonic oscillator form as above, the exact eigenvalues and eigenstates to the confining 3q Hamiltonian are \(E = (N + 3)\hbar\omega + 3V_0\), \(\Psi = |N(\lambda\mu)L[f]_X[f]_FS[f]_F[S] >\), where \(N\) is the number of quanta in the state, the Elliott symbol \((\lambda\mu)\) characterizes the \(SU(3)\) harmonic oscillator symmetry, and \(L\) is the orbital angular momentum. The spatial (X), flavor-spin (FS), flavor (F), and spin (S) permutational symmetries are indicated by corresponding Young patterns \((\lambda\mu)\).

When the boson exchange interaction \((3)\) is treated in first order perturbation theory, the mass of the baryon states takes the form \(M = M_0 + N\hbar\omega + \delta M_X\), where the chiral interaction contribution is \(\delta M_X = <\Psi|H_X|\Psi\rangle\), and \(M_0 = \sum_{i=1}^3 m_i + 3(V_0 + \hbar\omega)\). The contribution from the chiral interaction to each baryon is a linear combination of the matrix elements of the two-body potential \(V(r_{12})\), defined as \(P_{nl} = <\varphi_{nlm}(\vec{r}_{12})|V(r_{12})|\varphi_{nlm}(\vec{r}_{12})\rangle\).

The oscillator parameter \(\hbar\omega\) and the four integrals are extracted from the mass differences between the nucleon and the \(\Delta(1232)\), the \(\Delta(1600)\) and the \(N(1440)\), as well as the splittings between the nucleon and the average mass of the two pairs of states \(N(1535) - N(1520)\) and \(N(1720) - N(1680)\). This procedure yields the parameter values \(\hbar\omega=157.4\) MeV, \(P_{00}=29.3\) MeV, \(P_{11}=45.2\) MeV, \(P_{20}=2.7\) MeV and \(P_{22}=-34.7\) MeV. Given these values, all other excitation energies (i.e., differences between the masses of given resonances and the corresponding ground states) of the nucleon, \(\Delta\)- and \(\Lambda\)-hyperon spectra are predicted to within \(\sim 15\%\) of the empirical values where known, and are well within the uncertainty limits of those values. Note that these matrix elements provide a quantitatively satisfactory description of the \(\Lambda\)-spectrum (Table 1) even though they are extracted from the \(N - \Delta\) spectrum.

5 Three-Body Faddeev Calculations

In the previous section we have shown how the Goldstone boson exchange (GBE), taken to first order perturbation theory and without explicit parameterizing the radial dependence, can explain the correct level ordering of positive and negative parity states in light and strange baryon spectra, as well as the splittings in those spectra. A question, however, arises about what will happen beyond first order perturbation theory. In order to check this we have numerically solved three-body Faddeev equa-
Table 1: The structure of the Λ-hyperon states up to $N = 2$, including predicted unobserved or nonconfirmed states, indicated by question marks. The predicted energies (in MeV) are given in the brackets under the empirical values.

| $N(\lambda\mu)L[f|x|f][f|f][f]|S$ | LS multiplet | average energy | $\delta M_x$ |
|-----------------------------|--------------|----------------|--------------|
| 0(00)[3]x[3]f[1]s[1]f[1]s[1]s | $\frac{1}{2}^+, \Lambda$ | 1115 | $-14P_{00}$ |
| 1(10)[2]x[2]f[1]s[1]f[1]s[1]s | $\frac{1}{2}^-, \Lambda(1405); \frac{3}{2}^-, \Lambda(1520)$ | 1462 (1512) | $-12P_{00} + 4P_{11}$ |
| 2(20)[0]x[3]f[2]s[2]f[2]s[2]s | $\frac{1}{2}^+, \Lambda(1600)$ | 1600 (1616) | $-7P_{00} - 7P_{20}$ |
| 1(10)[2]x[2]f[2]s[2]f[2]s[2]s | $\frac{1}{2}^-, \Lambda(1670); \frac{3}{2}^-, \Lambda(1690)$ | 1680 (1703) | $-7P_{00} + 5P_{11}$ |
| 1(10)[1]x[2]f[2]s[2]f[2]s[2]s | $\frac{1}{2}^-, \Lambda(1800); \frac{3}{2}^-, \Lambda(?)$; $\frac{5}{2}^-, \Lambda(1830)$ | 1815 (1805) | $-2P_{00} + 4P_{11}$ |
| 2(20)[0]x[2]f[1]s[1]f[1]s[1]s | $\frac{1}{2}^+, \Lambda(1810)$ | 1810 (1829) | $-6P_{00} - 6P_{20} + 4P_{11}$ |
| 2(20)[2]x[3]f[2]s[2]f[2]s[2]s | $\frac{3}{2}^+, \Lambda(1890); \frac{5}{2}^+, \Lambda(1820)$ | 1855 (1878) | $-7P_{00} - 7P_{22}$ |
| 2(20)[0]x[2]f[2]s[2]f[2]s[2]s | $\frac{1}{2}^+, \Lambda(?)$ | ? (1954) | $-\frac{7}{2}P_{00} - \frac{7}{2}P_{20} + 5P_{11}$ |
| 2(20)[0]x[2]f[2]s[2]f[2]s[2]s | $\frac{3}{2}^+, \Lambda(?)$ | ? (1989) | $-P_{00} - P_{20} + 4P_{11}$ |
| 2(20)[2]x[2]f[2]s[2]f[2]s[2]s | $\frac{1}{2}^+, \Lambda(?)$; $\frac{3}{2}^+, \Lambda(?)$; $\frac{5}{2}^+, \Lambda(2020?)$ | 2020? (2026) | $-P_{00} - P_{22} + 4P_{11}$ |
| 2(20)[2]x[2]f[1]s[1]f[1]s[1]s | $\frac{3}{2}^+, \Lambda(?)$; $\frac{5}{2}^+, \Lambda(2110)$ | 2110? (2085) | $-6P_{00} - 6P_{22} + 4P_{11}$ |

Besides the confinement potential, which is now taken in linear form, the GBE interaction between the constituent quarks is now included to all orders. These results further support the adequacy of the GBE for baryon spectroscopy.

In addition to the octet-exchange interaction we include here also the flavor-singlet ($\eta'$) exchange. In the large $N_C$ limit the axial anomaly becomes suppressed [17], and the $\eta'$ becomes the ninth Goldstone boson of the spontaneously broken $U(3)_L \times U(3)_R$ chiral symmetry.

We show our results in Fig. 1 It is well seen that the whole set of lowest N and Δ states is reproduced quite correctly. In the most unfavourable cases deviations from the experimental values do not exceed 3%! In addition all level orderings are correct.
In particular, the positive-parity state \( N(1440) \) (Roper resonance) lies below the pair of negative-parity states \( N(1535) - N(1520) \). The same is true in the \( \Delta \) spectrum with \( \Delta(1600) \) and the pair \( \Delta(1620) - \Delta(1700) \).

![Energy levels for the 14 lowest non-strange baryons with total angular momentum and parity \( J^P \). The shadowed boxes represent experimental uncertainties.](image)

6 Binding of Quarks and the \( \pi N \sigma \)-Term

The pion-nucleon \( \sigma \)-term is a measure of the explicit chiral symmetry breaking effects in the nucleon. The additive quark ansatz, where the nucleon is considered as a system of three weakly interacting constituent quarks, leads to a much smaller value than the empirical result extracted from pion-nucleon scattering data \[18\]. This indicates that some essential piece of physics is absent within the additive quark ansatz. It is shown in ref. \[19\] that the contribution to \( \sigma_{\pi N} \) that arises from the short range part of GBE between the constituent quarks is crucial for the explanation of its empirical value.

7 Instead of a Conclusion
Instead of a conclusion we discuss some important recent lattice QCD results in this last section. It was shown already a few years ago that one can obtain a qualitatively correct splitting between $\Delta$ and $N$ already within a quenched approximation (for a review and references see [20]). In the quenched approximation for baryons one takes into account only 3 continuous valence quark lines and full gluodynamics. This quenched approximation contains, however, part of antiquark effects related to the $Z$ graphs formed of valence quark lines. One can even construct diagrams within the quenched approximation which correspond to the exchange of the color-singlet isospin 1 or 0 $q\bar{q}$ pairs between valence quark lines [21]. It is also important that these diagrams contribute to the baryon mass to leading order ($\sim N_C$) in a $1/N_C$ expansion (their contribution to the $\Delta - N$ splitting appears, however, to subleading orders).

From the quenched measurements [20] it is not clear what were the physical reason for the $\Delta - N$ splitting: gluon exchanges, instantons, or something else. To clarify this question, Liu and Dong have recently measured the $\Delta - N$ splitting in the quenched and a further so-called "valence approximation" [22]. In the valence approximation the quarks are limited to propagating only forward in time (i.e., $Z$ graphs and related quark-antiquark pairs are removed). The gluon exchange and all other possible gluon configurations, including instantons, are exactly the same in both approximations. The striking result is that the $\Delta - N$ splitting is observed only in the quenched approximation but not in the valence approximation, in which the $N$ and the $\Delta$ levels are degenerate within error bars. Consequently the $\Delta - N$ splitting must receive a considerable contribution from the diagrams with $q\bar{q}$ excitations, which correspond to the meson exchanges, but not from the gluon exchange or instanton-induced interaction between quarks (the instanton-induced interaction could be rather important for the interactions between quarks and antiquarks as it is strongly attractive in the $q\bar{q}$ pseudoscalar channel while it is weak in $qq$ pairs).

Finally, the flavor-spin structure of the interaction (3) is also compatible with the vector meson ($\rho, K^*, \omega$) and axial-vector meson exchanges [23]. The radial behaviour of the spin-spin force associated with the vector meson exchange is similar to that of in (2), while in the axial-vector meson exchange case the contact term is absent and the required sign of the interaction (3) at short range comes from the Yukawa tail. Thus one cannot exclude a possibility that in reality the flavor-spin interaction (3) is some superposition of all possible meson exchanges. It is well known that the spin-spin components of the pseudoscalar and vector meson exchange interactions have the same sign while their tensor components tend to cancel. This could be an additional reason for why the tensor force is not so important for baryons.

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