UNQUENCHING THE QUARK MODEL

E. SANTOPINTO
I.N.F.N. and Dipartimento di Fisica,
via Dodecaneso 33, Genova, I-16146, ITALY
E-mail: santopinto@ge.infn.it

R. BIJKER
Departamento de Estructura de la Materia,
Instituto de Ciencias Nucleares,
Universidad Nacional Autónoma de México,
AP 70-543, 04510 México DF, MEXICO
E-mail: bijker@nucleares.unam.mx

We present the formalism for a new generation of unquenched quark models in which quark-antiquark pair effects are taken into account in an explicit form via a microscopic QCD-inspired quark-antiquark creation mechanism. No truncation in the sum over all the big tower of states is necessary since these states are automatically generated by means of powerful group-theoretical techniques. An important check on the formalism and the numerical results is provided by the closure limit. As an application, the effect of quark-antiquark pairs on the strange content of the proton spin, $\Delta_s$, is discussed. The contributions of the up and down quarks, $\Delta_u$ and $\Delta_d$, are also calculated. This has become possible after solving the difficult problem of permutational symmetry related to quark rearrangements. Finally, we present some preliminary results in the closure limit as well as an outlook for future applications.

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I. INTRODUCTION

Since the QCD equations in the non-perturbative region are not solvable, many effective models of the hadrons have been developed, in particular many versions of the Constituent Quark Model (CQM). These are able to give good results for the static properties of the hadrons (like the spectrum and the magnetic moments), while they all fail to reproduce the dynamic ones, like the electromagnetic transition form factors at low $Q^2$ values. This seems to be a problem of degrees of freedom, since the region of low $Q^2$ means high distance, in which the creation of quark-antiquark-pair degrees has higher probability.

A priori, one would expect pair creation to be (after switching on a pair-creation mechanism) so enhanced in probability that a valence quark model would fail dramatically. On the other hand, it is known from the phenomenology that pair creation is suppressed, at least at the zeroth order in the static observables, since the observed spectrum is dominated by the valence quarks $qqq$ for the baryons and by the valence $q\bar{q}$ for the mesons. Moreover, the resonances do not have decay widths of the order of the same resonance masses. Many years ago, Isgur tackled the intriguing and important question of why the CQMs work so well at an intermediate regime. The aim of this paper is to continue from where he left off and to develop the formalism for the light baryons with the linked problem of the permutational symmetry.

II. DEGREES OF FREEDOM

Constituent Quark Models based on the effective degrees of freedom of three constituent quarks have been proposed in several versions: the Isgur-Karl model [1], the Capstick-Isgur model [2], the algebraic U(7) model [3], the hypercentral model [4], the chiral boson exchange model [5][6] and the Bonn instanton model [7]. While these models display important and peculiar differences, they all have main features in common: they are all based on the effective degrees of freedom of three constituent quarks and on the SU(6) spin-flavor symmetry; they also contain a long-range linear confining potential and an SU(6)-breaking term, even though the form and the advocated origin of this last term may be different, as in the one-gluon-exchange-inspired hyperfine interaction or the Goldston Boson Exchange SU(6)-breaking interaction or the instanton-induced-breaking term.

The photocouplings calculated by means of different constituent quark models (see for example Table 2 of Ref. [8], and references therein) have the same overall behaviour, having the same SU(6) structure in common, but in many cases they show a lack of strength.
For each of those models for which calculations of the electromagnetic transition form factors are available, there is clearly the problem of missing strength at low $Q^2$, as is well illustrated in Fig. 1 which reports the transverse electromagnetic transition form factors for the $D_{13}(1520)$ resonance. The experimental data are compared with the predictions of different CQMs. The common problem of missing strength at low $Q^2$ can be ascribed to the lack of quark-antiquark effects, which are probably more important in the outer region of the nucleon. After understanding that the main problems of the CQMs are all linked to a problem of missing-degrees of freedom, and after identifying the key degrees of freedom as $q\bar{q}$ pairs, we are faced with two possibilities: the phenomenological parametrization or the microscopic explicit description. We have focused on the letter and tried to find a QCD-inspired pair-creation mechanism that at the same time should be encoded in such a way not as not to destroy the good CQM results. Moreover, it is important to find symmetry constraints in order to help to keep all the difficult calculations under control. In this respect, the closure limit, is very helpful, as will be explained in the following sections.

III. FLUX-TUBE BREAKING MODEL

In the flux-tube model for hadrons, the quark potential model arises from an adiabatic approximation to the gluonic degrees of freedom embodied in the flux tube [11]. The impact of $q\bar{q}$ pairs in meson spectroscopy has been studied in an elementary flux-tube breaking model [12] in which the quark pair creation occurs with $^3P_0$ quantum numbers. Subsequently, it has been shown by Geiger and Isgur [13] that a "miraculous" set of cancellations between apparently uncorrelated sets of mesons occurs in such a way to compensate each other and not to destroy the good CQM results for the mesons, and in particular in such a way that (i) the OZI hierarchy is preserved and (ii) there is a near immunity of the long-range confining potential. The change in the linear potential due to pairs bubbling in the string can be reabsorbed in a new strength of the linear potential, i.e., in a new string tension; thus the net effect of the mass shifts from pair creation is smaller than the naive expectation of the order of the strong decay widths. The important point is that it is necessary to sum over very large towers of intermediate states to see that the spectrum of the mesons is only weakly perturbed, after unquenching and renormalizing. No simple truncation of the set of meson loops can reproduce such results.

This formalism set up for mesons could not be immediately extended to baryons, since there were extra problems, mainly linked with the permutation symmetry between identical quarks. In [14], Geiger and Isgur investigated the importance of $s\bar{s}$ loops in the proton by taking into account the contribution of the six different diagrams of Fig. 2 (in Fig. 3 of [14] only four diagrams are considered) and by using harmonic oscillator wave functions for the baryons and mesons.

In the present manuscript, we discuss some generalizations of this formalism for baryons. These new extensions make it possible to study the quark-antiquark contributions

- for any initial baryon resonance
- for any flavor of the quark-antiquark pair
- for any model of baryons and mesons, as long as their wave functions are expressed on the basis of the harmonic oscillator.
The problem of the permutation symmetry between identical quarks has been solved by means of group-theoretical techniques. In this way, we can evaluate the contribution of all the diagrams for any initial baryon $q_1q_2q_3$ (ground state or resonance) and for any flavor of the $q\bar{q}$ and not just for the $s\bar{s}$.

The pair-creation mechanism is inserted at the level of quarks and the one-loop diagram should be calculated, see Fig. 3 by summing over all the possible intermediate states. This is the crucial step, after solving all the problems linked with this infinite sum in the case of the baryons.

\[
\hat{M} h_{q\bar{q}} = \hat{M} (BC\vec{k}lJ) \frac{\langle BC\vec{k}lJ | h_{q\bar{q}} | A \rangle}{M_A - E_B - E_C}.
\]

where $h_{q\bar{q}}$ is the quark-antiquark pair-creation operator as in [14], $A$ is the initial baryon, $B$ ($C$) denotes the intermediate baryon (meson), $\vec{k}$ and $l$ the relative radial momentum and orbital angular momentum of $B$ and $C$, and $J$ is the total angular momentum $\vec{J} = \vec{J}_B + \vec{l} = \vec{J}_B + \vec{J}_C + \vec{l}$. In general, matrix elements of observables $\hat{O}$ are given by

\[
\langle \psi_A | \hat{O} | \psi_A \rangle = \mathcal{O}_{\text{valence}} + \mathcal{O}_{\text{sea}}
\]

where the first term denotes the contribution from the three valence quarks

\[
\mathcal{O}_{\text{valence}} = \mathcal{N}^2 \langle A | \hat{O} | A \rangle
\]

and the second term the contribution from the quark-antiquark pairs

\[
\mathcal{O}_{\text{sea}} = \mathcal{N}^2 \sum_{BCIJ} \int d\vec{k} \sum_{B'C'J'} \int d\vec{k}' \langle A | h_{q\bar{q}} | B'C'\vec{k}'l'J' \rangle \frac{\langle BC\vec{k}lJ | h_{q\bar{q}} | A \rangle}{M_A - E_B - E_C}.
\]
The sum is over a complete set of intermediate states, rather than just a few low-lying states. Not only does this have a significant impact on the numerical result, but it is necessary for consistency with the OZI-rule and the success of the CQMs in spectroscopy. This is possible since the intermediate states are generated automatically by means of group-theory techniques. It has been implemented in a computer code which is able to construct states with the correct permutation symmetry for any model and up to any shell. Therefore, the sum over intermediate states can be performed up to saturation and not just for the first few shells as in [14].

IV. CLOSURE LIMIT

In this section, we discuss the closure limit which arises when the energy denominators do not depend strongly on the quantum numbers of the intermediate states in Eq. (4). In this case, the sum over the complete set of intermediate states can be solved by closure and the contribution of the quark-antiquark pairs to the matrix element reduces to

\[ \mathcal{O}_{\text{sea}} \propto \langle A | h_{\bar{q}q} \mathcal{O} h_{\bar{q}q}^\dagger | A \rangle. \] (5)

This means for example that, at least qualitatively, we can already say that the strange content of the proton should be very small, since within the closure limit it comes out to be exactly zero. The closure limit is very important for many reasons that are closely interconnected. It is a very good stringent test that a numerical calculation should fulfill. Moreover, and even more important, it also explains, when it comes out equal to zero, the success of the CQMs. The corrections due to the \( q\bar{q} \) pairs are zero in the closure limit for many observables, while the sums over big tower of states are constrained by the closure limit in such a way that very different meson and baryon states compensate one another.

V. PRELIMINARY RESULTS

In this section, we discuss some preliminary results in the closure limit for the operator \( \Delta q \) to determine the fraction of the baryon’s spin carried by each one of the flavors \( u, d \) and \( s \)

\[ \Delta q = 2(S_z(q) + \bar{S}_z(\bar{q})). \] (6)

In the closure limit the relative contribution of the quark flavors from the quark-antiquark pairs to the baryon spin is the same as that from the valence quarks

\[ \Delta u_{\text{sea}} : \Delta d_{\text{sea}} : \Delta s_{\text{sea}} = \Delta u_{\text{valence}} : \Delta d_{\text{valence}} : \Delta s_{\text{valence}}. \] (7)

Table I shows the relative contributions of \( \Delta u, \Delta d \) and \( \Delta s \) to the spin of the ground state octet baryons in the closure limit. The closure limit, when combined with symmetries, imposes other strong limits, as for example can be seen for the ground state decuplet baryons in Fig. [4]. For the \( \Delta \) resonances whose three-quark configuration does not contain strange quarks, the contribution of the \( s\bar{s} \) pairs to the spin \( \Delta s \) and the magnetic moment \( \mu_s \) vanishes in the closure limit. The same holds for the contribution of \( d\bar{d} \) pairs to the \( \Delta^{++}, \Sigma^{*+}, \Xi^{*0} \) and \( \Omega^- \) resonances, and that of \( u\bar{u} \) pairs to the \( \Delta^-, \Sigma^{-}, \Xi^{-} \) and \( \Omega^- \) resonances. These are very stringent requirements since each one involves the sum over a complete set of states of the intermediate baryon-meson system and provides a nontrivial test which involves the spin-flavor section, the permutation symmetry and the correct implementation of large towers of intermediate states.

| \( \Delta u \) : \( \Delta d \) : \( \Delta s \) |
| --- |
| \( p \) | 4 : -1 : 0 |
| \( n \) | 1 : 4 : 0 |
| \( \Sigma^+ \) | 4 : 0 : -1 |
| \( \Sigma^0 \) | 2 : 2 : -1 |
| \( \Sigma^+ \) | 0 : 4 : -1 |
| \( \Lambda \) | 0 : 0 : 3 |
| \( \Xi^0 \) | -1 : 0 : 4 |
| \( \Xi^- \) | 0 : -1 : 4 |
VI. CONCLUSIONS AND OUTLOOK

The formalism described in this paper will be systematically used in order to study the still open problems regarding light baryons, such as the sea quark content and the spin problem of the proton \cite{15} and the $q\bar{q}$ effects in the electromagnetic elastic and transition form factors \cite{16}.

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\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig4}
\caption{Ground state decuplet baryons}
\end{figure}

[1] N. Isgur and G. Karl, \textit{Phys. Rev. D} \textbf{20}, 1191 (1979).
[2] S. Capstick and N. Isgur, \textit{Phys. Rev. D} \textbf{34}, 2809 (1986).
[3] R. Bijker, F. Iachello, and A. Leviatan, \textit{Ann. Phys. (N.Y.)} \textbf{236}, 69 (1994); \textit{ibid.} \textbf{284}, 89 (2000).
[4] M. Ferraris, M. M. Giannini, M. Pizzo, E. Santopinto, and L. Tiator, \textit{Phys. Lett. B} \textbf{364}, 231 (1995).
[5] L. Ya. Glozman and D.O. Riska, \textit{Phys. Rep. C} \textbf{268}, 263 (1996).
[6] L. Ya. Glozman, Z. Papp, W. Plessas, K. Varga, and R. F. Wagenbrunn, \textit{Phys. Rev. C} \textbf{57}, 3406 (1998); L. Ya. Glozman, W. Plessas, K. Varga, and R. F. Wagenbrunn, \textit{Phys. Rev. D} \textbf{58}, 094030 (1998).
[7] U. L"oring, K. Kretzschmar, B. Ch. Metsch, and H. R. Petry, \textit{Eur. Phys. J. A} \textbf{10}, 309 (2001); U. L"oring, B.Ch. Metsch, and H. R. Petry, \textit{Eur. Phys. J. A} \textbf{10}, 395 (2001); \textit{ibid.} \textbf{447} (1998).
[8] M. Aiello, M. Ferraris, M. M. Giannini, M. Pizzo, and E. Santopinto, \textit{Phys. Lett. B} \textbf{387}, 215 (1996).
[9] M. Aiello, M. M. Giannini, and E. Santopinto, \textit{J. Phys. G} \textbf{24}, 753 (1998).
[10] R. Bijker, F. Iachello, and A. Leviatan, \textit{Phys. Rev. C} \textbf{54}, 1935 (1996).
[11] N. Isgur and J. Paton, \textit{Phys. Rev. D} \textbf{31}, 2910 (1985).
[12] R. Kokoski and N. Isgur, \textit{Phys. Rev. D} \textbf{35}, 907 (1987).
[13] P. Geiger and N. Isgur, \textit{Phys. Rev. Lett.} \textbf{67}, 1066 (1991); \textit{Phys. Rev. D} \textbf{44}, 799 (1991); \textit{ibid.} \textbf{47}, 5050 (1993).
[14] P. Geiger and N. Isgur, \textit{Phys. Rev. D} \textbf{55}, 299 (1997).
[15] R. Bijker and E. Santopinto, to be published.
[16] R. Bijker and E. Santopinto, work in progress.