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Meson-loop contributions in the quark model

Roelof Bijker

Instituto de Ciencias Nucleares, Universidad Nacional Autónoma de México, A.P. 70-543, 04510 México, D.F., México

E-mail: bijker@nucleares.unam.mx

Elena Santopinto

I.N.F.N., Sezione di Genova, via Dodecaneso 33, I-16145, Italy

E-mail: santopinto@ge.infn.it

Abstract. We present a review of some recent results obtained in the unquenched quark model for baryons. The effects of sea quarks are taken into account in an explicit form via a QCD-inspired creation mechanism of the quark-antiquark pairs. In this approach, the contribution of the quark-antiquark pairs can be studied for any inital baryon and for any flavor of the $q\bar{q}$ pairs ($u\bar{u}$, $d\bar{d}$ and $s\bar{s}$). It is shown that, whereas the inclusion of $q\bar{q}$ pairs (or meson loops) does not affect the baryon magnetic moments, it immediately leads to an excess of $\bar{d}$ over $\bar{u}$ in the proton and introduces a sizeable contribution of orbital angular momentum to the spin of the proton. The contribution of $s\bar{s}$ pairs to the magnetic moment and the radius of the proton is found to be small, in agreement with the latest experimental results and recent lattice calculations.

1. Introduction

In the constituent quark model (CQM) hadrons are described as a system of constituent (or valence) quarks and antiquarks, qqq for baryons and $q\bar{q}$ for mesons. Despite the success of the quark model there is strong evidence for the existence of exotic degrees of freedom (other than valence quarks) in hadrons. Common features of constituent quark models for baryons are the effective degrees of freedom of three constituent quarks (qqq configurations), the $SU(6)$ spin-flavor symmetry and a long-range confining potential. In general, CQMs reproduce the mass spectrum of baryon resonances reasonably well, but at the same time, they show very similar deviations for other observables, such as photocouplings, helicity amplitudes and strong decays. Since the photocouplings depend mostly on the spin-flavor structure, all models that have the same $SU(6)$ structure in common, show the same behavior, *e.g.* the photocouplings for the $\Delta(1232)$ are underpredicted by a large amount, even though their ratio is reproduced correctly. In general, the helicity amplitudes (or transition form factors) show deviations from CQM calculations at low values of Q^2 [1, 2, 3]. The problem of missing strength at low Q^2 indicates that some fundamental mechanism is lacking in the dynamical description of hadronic structure. This mechanism can be identified with the production of quark-antiquark pairs, which becomes more important in the outer region of the nucleon.

Additional evidence for such higher Fock components in the baryon wave function ($qqq - q\bar{q}$ configurations) comes from CQM studies of strong decay widths of baryon resonances that are on average underpredicted by CQMs [4, 5, 6], the spin-orbit splitting of $\Lambda(1405)$ and $\Lambda(1520)$, and the large η decay widths of $N(1535)$, $\Lambda(1670)$ and $\Sigma(1750)$. More direct evidence for the importance of quark-antiquark components in the proton comes from measurements of the $\bar{d}/\bar{u}$ asymmetry in the nucleon sea [7, 8], the proton spin crisis [9, 10] and parity-violating electron scattering experiments which report a nonvanishing strange quark contribution, albeit small, to the charge and magnetization distributions [11, 12].

The aim of this contribution is to present a review of some results recently obtained for the unquenched quark model, in which the effects of quark-antiquark pair creation ($u\bar{u}$, $d\bar{d}$ and $s\bar{s}$) are taken into account in an explicit form via a $^3P_0$ coupling mechanism [13, 14, 15]. In order to test the consistency of the formalism we first calculate the baryon magnetic moments [13] which constitute one of the early successes of the CQM. Next we discuss the spin and flavor content of the proton. Whereas the flavor asymmetry and the orbital angular momentum are dominated by pion loops, the contribution of the sea quark spins to the proton spin arises almost entirely from (excited) vector meson loops [16]. Finally, we study the strange magnetic moment and the strange radius of the proton [17].

2. Unquenched quark model

The role of higher Fock components in baryon wave functions has been studied by many authors in the context of meson cloud models, such as the cloudy bag model, meson convolution models and chiral models [7, 18]. In these models, the flavor asymmetry of the proton can be understood in terms of couplings to the pion cloud. There have also been several attempts to study the importance of higher Fock components in the context of the constituent quark model. In this respect we mention the work by Riska and coworkers who introduce a small number of selected higher Fock components which are then fitted to reproduce the experimental data [19]. However, these studies lack an explicit model or mechanism for the mixing between the valence and sea quarks. The Rome group studied the pion and nucleon electromagnetic form factors in a Bethe-Salpeter approach, mainly thanks to the dressing of photon vertex by means of a vector-meson dominance parametrization [20]. Koniuk and Guiasu used a convolution model with CQM wave functions and an elementary emission model for the coupling to the pion cloud to calculate the magnetic moments and the helicity amplitudes from the nucleon to the Δ resonance [21]. It was found that the nucleon magnetic moments were unchanged after renomalization of the parameters, but that the missing strength in the helicity amplitudes of the Δ could not be explained with pions only.

The impact of $q\bar{q}$ pairs in hadron spectroscopy was originally studied by Törnqvist and Zenczykowski in a quark model extended by the $^3P_0$ model [22]. Even though their model only includes a sum over ground state baryons and ground state mesons, the basic idea of the importance to carry out a sum over a complete set of intermediate states was proposed in there. Subsequently, the effects of hadron loops in mesons was studied by Geiger and Isgur in a flux-tube breaking model in which the $q\bar{q}$ pairs are created in the $^3P_0$ state with the quantum numbers of the vacuum [23, 24, 25]. In this approach, the quark potential model arises from an adiabatic approximation to the gluonic degrees of freedom embodied in the flux-tube [26]. It was shown that cancellations between apparently uncorrelated sets of intermediate states occur in such a way that the modification in the linear potential can be reabsorbed, after renormalization, in the new strength of the linear potential [24]. In addition, the quark-antiquark pairs do not destroy the good CQM results for the mesons [24] and preserve the OZI hierarchy [25] provided that the sum be carried out over a large tower of intermediate states. A first application of this procedure to baryons was presented in [27] in which the importance of $s\bar{s}$ loops in the proton were studied by taking into account the contribution of the six diagrams of Fig. 1 in combination with harmonic oscillator wave functions for the baryons and mesons and a $^3P_0$ pair creation mechanism. This approach has the advantage that the effects of quark-antiquark pairs are introduced explicitly via a QCD-inspired pair-creation mechanism, which opens the

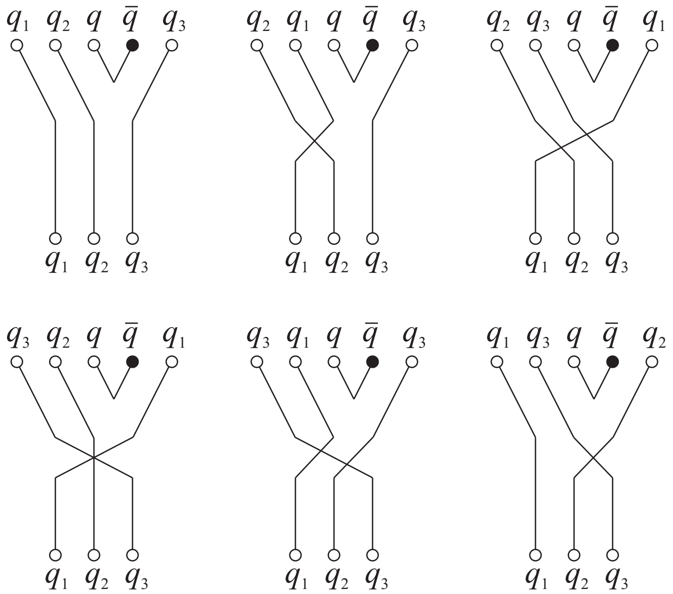

Figure 1. Quark line diagrams for $A \rightarrow BC$ with $q\bar{q} = s\bar{s}$ and $q_1q_2q_3 = uud$.

possibility to study the importance of $q\bar{q}$ pairs in baryons and mesons in a systematic and unified way.

The present approach is motivated by these earlier studies on extensions of the quark model to include the effects of $q\bar{q}$ pairs [22, 27]. Our approach is based on a CQM to which the quark-antiquark pairs with vacuum quantum numbers are added as a perturbation employing a $^3P_0$ model for the $q\bar{q}$ pair creation [13, 14, 27]. The pair-creation mechanism is inserted at the quark level and the one-loop diagrams are calculated by summing over a complete set of intermediate baryon-meson states. Under these assumptions, the baryon wave function consists of a zeroth order three-quark configuration $\mid A\rangle$ plus a sum over all possible higher Fock components due to the creation of $^3P_0$ quark-antiquark pairs

$$\mid \psi_A\rangle \;=\; \mathcal{N}\left[\mid A\rangle + \sum_{BCl}\int d\vec{k}\,\mid BC\vec{k}\,lJ\rangle\frac{\langle BC\vec{k}\,lJ\mid T^{\dagger}\mid A\rangle}{M_A - E_B - E_C}\right] \;. \tag{1}$$

Here A denotes the initial baryon, B and C represent the intermediate baryon and meson, and M_A, E_B and E_C are their respective energies, $\vec{k}$ and l the relative radial momentum and orbital angular momentum of B and C, and J is the total angular momentum $\vec{J} = \vec{J}_B + \vec{J}_C + \vec{l}$. The operator $T^{\dagger}$ creates a quark-antiquark pair in the $^3P_0$ state with the quantum numbers of the vacuum: $L = S = 1$ and $J = 0$ [13, 14, 28]

$$\begin{aligned} T^{\dagger} \;&=\; -3\sum_{ij}\int d\vec{p}_i\, d\vec{p}_j\, \delta(\vec{p}_i + \vec{p}_j)\, C_{ij}\, F_{ij}\, \Gamma(\vec{p}_i - \vec{p}_j) \\ &\qquad [\chi_{ij} \times \mathcal{Y}_1(\vec{p}_i - \vec{p}_j)]^{(0)}\, b_i^{\dagger}(\vec{p}_i)\, d_j^{\dagger}(\vec{p}_j)\;. \end{aligned} \tag{2}$$

Here, $b_i^{\dagger}(\vec{p}_i)$ and $d_j^{\dagger}(\vec{p}_j)$ are the creation operators for a quark and antiquark with momenta $\vec{p}_i$ and $\vec{p}_j$, respectively. The quark pair is characterized by a color singlet wave function C_{ij}, a flavor singlet wave function F_{ij} and a spin triplet wave function χ_{ij} with spin $S = 1$. The solid harmonic $\mathcal{Y}_1(\vec{p}_i - \vec{p}_j)$ indicates that the quark and antiquark are in a relative P wave. The $SU(3)$ flavor symmetry of the valence quark configuration $\mid A\rangle$ is broken by the quark-antiquark pairs via the energy denominator, but the $SU(2)$ isospin symmetry is still preserved. In the special case of the closure limit in which the energy denominator of Eq. (1) is a constant, the flavor symmetry of the valence quark configuration is recovered.

Since the operator $T^\dagger$ creates a pair of constituent quarks, a Gaussian quark-antiquark creation vertex function was introduced by which the pair is created as a finite object with an effective size, rather than as a pointlike object. In momentum space it is given by

$$\Gamma(\vec{p}_i - \vec{p}_j) \;=\; \gamma_0 \, \mathrm{e}^{-r_q^2(\vec{p}_i - \vec{p}_j)^2/6} \; . \tag{3}$$

The width has been determined from meson decays to be approximately 0.25−0.35 fm [25, 27, 29]. Here we take the average value, $r_q = 0.30$ fm. Finally, the dimensionless constant γ_0 is the intrinsic pair creation strength which can be determined from the strong decays of baryons [4].

The strong coupling vertex

$$\langle BC\vec{k}\, lJ \mid T^\dagger \mid A\rangle \; , \tag{4}$$

was derived in explicit form in the harmonic oscillator basis [28]. In the present calculations, we use harmonic oscillator wave functions in which there is a single oscillator parameter for the baryons and another one for the mesons which, following [27], are taken to be $\beta_{\rm baryon} = 0.32$ GeV [30] and $\beta_{\rm meson} = 0.40$ GeV [23], respectively.

In order to calculate the effects of quark-antiquark pairs on an observable, one has to evaluate the contribution of all possible intermediate states. By using a combination of group theoretical and computational techniques, the sum over intermediate states is carried out up to saturation and not only for the first few shells as in previous studies [22, 27]. 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 CQMs in hadron spectroscopy. In addition, the contributions of quark-antiquark pairs can be evaluated for any initial baryon (ground state or resonance) and for any flavor of the $q\bar{q}$ pair (not only $s\bar{s}$ as in [27], but also $u\bar{u}$ and $d\bar{d}$), and for any model of baryons and mesons, as long as their wave functions are expressed in the basis of harmonic oscillator wave functions [13, 14].

In the calculations presented in this contribution, we use harmonic oscillator wave functions up to five oscillator shells for the intermediate baryons and mesons. All parameters were taken from the literature without attempting to optimize their values in order to improve the agreement with experimental data [13, 14].

2.1. Closure limit

Before discussing an application of the unquenched model to baryon magnetic moments and spins, we study the so-called closure limit in which the intermediate states appearing in Eq. (1) are degenerate in energy and hence the energy denominator becomes a constant independent of the quantum numbers of the intermediate states. In the closure limit, the evaluation of the contribution of the quark-antiquark pairs (or the higher Fock components) simplifies considerably, since the sum over intermediate states can be solved by closure and the contribution of the quark-antiquark pairs to the matrix element reduces to

$$\mathcal{O}_{\rm sea} \;\propto\; \langle A \mid T\, \hat{\mathcal{O}}\, T^\dagger \mid A\rangle \; . \tag{5}$$

Since the $^3P_0$ pair-creation operator of Eq. (2) is a flavor singlet and the energy denominator in Eq. (1) is reduced to a constant in the closure limit, the higher Fock components of the baryon wave function have the same $SU(3)$ flavor symmetry as the valence quark configuration $\mid A\rangle$.

At a qualitative level, the closure limit helps to explain the phenomenological success of the CQM because the $SU(3)$ flavor symmetry of the baryon wave function is preserved. As an example, the strange content of the proton vanishes in the closure limit due to many cancelling contributions in the sum over intermediate states in Eq. (1). Away from the closure limit, the strangeness content of the proton is expected to be small, in agreement with the experimental data from parity-violating electron scattering (for some recent data see [11, 12]). Even though in this case the cancellations are no longer exact, many intermediate states contribute with opposite signs, and the net result is nonzero, but small. This means that even if the flavor symmetry of the CQM is broken by the higher Fock components, the net results are still to a large extent determined by the flavor symmetry of the valence quark configuration. Similar arguments were applied to the preservation of the OZI hierarchy in the context of the flux-tube breaking model [25]. Therefore, the closure limit not only provides simple expressions for the relative flavor content of physical observables, but also gives further insight into the origin of cancellations between the contributions from different intermediate states.

In addition, the closure limit imposes very stringent conditions on the numerical calculations, since it involves the sum over all possible intermediate states. Therefore, the closure limit provides a highly nontrivial test of the computer codes which involves both the spin-flavor sector, the permutation symmetry, the construction of a complete set of intermediate states in spin-flavor space for each radial excitation and the implementation of the sum over all of these states.

In the following, we study the effects of quark-antiquark pairs on the magnetic moments, radii, and the flavor and spin content of baryons in the general case, *i.e.* beyond the closure limit.

2.2. Magnetic moments

The unquenching of the quark model has to be carried out in such a way as to preserve the phenomenological successes of the constituent quark model. In applications to mesons, it was shown that the inclusion of quark-antiquark pairs does not destroy the good CQM results [24] and preserves the OZI hierarchy [25]. In a similar fashion, in this Section we will show that the CQM results for the magnetic moments of the octet baryons also hold in the unquenched CQM [13].

It is well known that the CQM gives a good description of the baryon magnetic moments, even in its simplest form in which the baryons are treated in terms of three constituent quarks in a relative S-wave. The quark magnetic moments are determined by fitting the magnetic moments of the proton, neutron and Λ hyperon to give $\mu_u = 1.852$, $\mu_d = -0.972$ and $\mu_s = -0.613\ \mu_N$ [31].

In the unquenched CQM the baryon magnetic moments also receive contributions from the quark spins of the pairs and the orbital motion of the quarks

$$\vec{\mu} \;=\; \sum_q \mu_q \left[2\vec{s}(q) + \vec{l}(q) - 2\vec{s}(\bar{q}) - \vec{l}(\bar{q}) \right] \;, \tag{6}$$

where $\mu_q = e_q\hbar/2m_qc$ is the quark magnetic moment. In Fig. 2 we show a comparison between the experimental values of the magnetic moments of the octet baryons (circles) and the theoretical values obtained in the CQM (squares) and in the unquenched quark model (triangles). The results for the unquenched quark model were obtained in a calculation involving a sum over intermediate states up to five oscillator shells for both baryons and mesons. The results obtained in the unquenched quark model are practically identical to the ones in the CQM, which shows that the addition of the quark-antiquark pairs preserves the good CQM results for the baryon magnetic moments. The effect of the $q\bar{q}$ pairs could be absorbed into renormalized values of the quark magnetic moments to $\mu_u = 2.066$, $\mu_d = -1.110$ and $\mu_s = -0.633\ \mu_N$, A similar feature was found in the context of the flux-tube breaking model for mesons in which it was shown that the inclusion of quark-antiquark pairs preserved the linear behavior of the confining potential as well as the OZI hierarchy [25]. The change in the linear potential caused

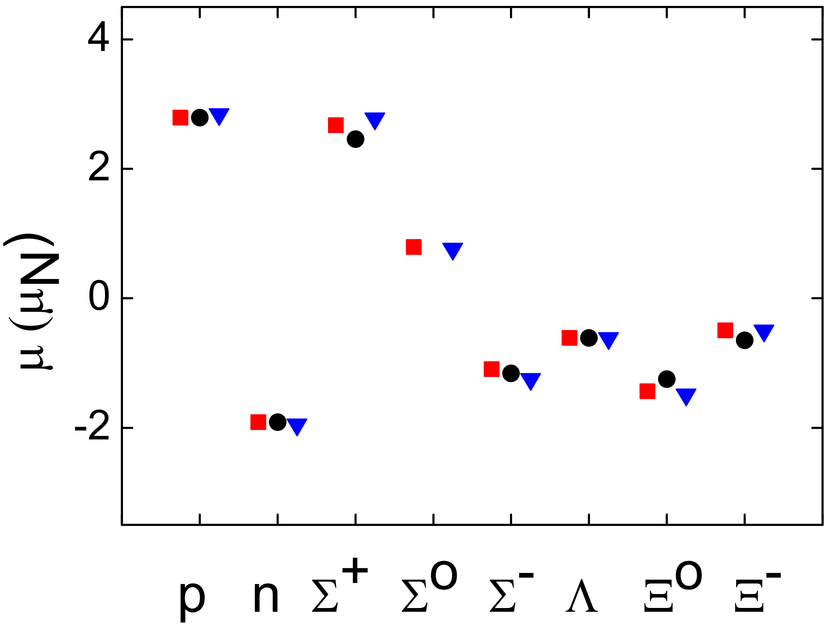

Figure 2. Magnetic moments of octet baryons: experimental values from PDG [31] (circles), CQM (squares) and unquenched quark model (triangles).

by the bubbling of the pairs in the string could be absorbed into a renormalized strength of the linear potential. The largest difference is observed for the charged Σ hyperons, but the relation between the magnetic moments of Σ hyperons $\mu(\Sigma^0) = [\mu(\Sigma^+) + \mu(\Sigma^-)]/2$ [32] is preserved in the unquenched calculation as a consequence of isospin symmetry.

The results for the magnetic moments can be understood qualitatively in the closure limit in which the relative contribution of the quark spins from the quark-antiquark pairs is the same as that from the valence quarks [13]. Moreover, since in the closure limit the contribution of the orbital angular momentum is small in comparison to that of the quark spins, the results for the baryon magnetic moments are almost indistinguishable from those of the CQM. Away from the closure limit, even though the relations between the different contributions no longer hold exactly, they are still valid approximately. In addition, there is now a contribution from the orbital part (at the level of ~ 5 %) which is mainly due to the baryon-pion channel.

In summary, the inclusion of the effects of quark-antiquark pairs preserves, after renormalization, the good results of the CQM for the magnetic moments of the octet baryons.

2.3. Flavor content

The flavor asymmetry of the proton $\mathcal{A}(p)$ is related to the Gottfried integral S_G for the difference of the proton and neutron electromagnetic structure functions as

$$S_G \;=\; \int_0^1 \frac{F_2^p(x) - F_2^n(x)}{x} dx \;=\; \frac{1}{3} - \frac{2}{3}\int_0^1 \left[\bar{d}_p(x) - \bar{u}_p(x)\right] dx \;=\; \frac{1}{3}[1 - 2\mathcal{A}(p)] \; . \tag{7}$$

Under the assumption of a flavor symmetric (or rather flavor independent) sea one obtains the Gottfried sum rule $S_G = 1/3$ [8, 33], whereas any deviation from this value is an indication of the $\bar{d}/\bar{u}$ asymmetry of the nucleon sea, thus providing evidence of the existence of higher

Table 1. Contributions to the flavor asymmetry of the proton [14].

	Unquenched QM	
	0-4 $\hbar\omega$	0 $\hbar\omega$
$N\pi$	0.195	0.177
$\Delta\pi$	–0.016	–0.010
$N\pi\eta_8\eta_1$	–0.028	–0.018
$N\rho$	0.050	0.012
$\Delta\rho$	–0.017	–0.003
$N\rho\omega_8\omega_1$	–0.033	–0.010
Total	0.151	0.147

Fock components (such as $qqq - q\bar{q}$ configurations) in the proton wave function. The first clear evidence of a violation of the Gottfried sum rule came from the New Muon Collaboration (NMC) [37] which was later confirmed by Drell-Yan experiments [34, 35] and a measurement of semi-inclusive deep-inelastic scattering [36]. All experiments show evidence that there are more $\bar{d}$ quarks in the proton than there are $\bar{u}$ quarks [8]. The final NMC value is 0.2281 ± 0.0065 at $Q^2 = 4$ $(\mathrm{GeV/c})^2$ for the Gottfried integral over the range $0.004 \leq x \leq 0.8$ [37], which implies a flavor asymmetric sea. The observed flavor asymmetry is far too large to be accounted for by processes that can be described by QCD in perturbative regime and therefore has to attributed to non-perturbative QCD mechanisms. It was shown in the framework of the meson-cloud model, that the coupling of the nucleon to the pion cloud provides a mechanism that is able to produce a flavor asymmetry due to the dominance of $n\pi^+$ among the virtual configurations [38].

In the unquenched quark model, the flavor asymmetry of the proton can be calculated directly from the difference of the number of $\bar{d}$ and $\bar{u}$ sea quarks in the proton, even in the absence of explicit information on the (anti)quark distribution functions. Table 1 shows that the flavor asymmetry for the proton in the UCQM is 0.151 [14] which corresponds to a value of the Gottfried integral of 0.232, remarkably close to the experimental value. The main contribution to the flavor asymmetry of the proton is due to the pion loops, especially the $n\pi^+$ intermediate state, thus confirming in an explicit calculation the explanation given in Ref. [38] in the context of the meson-cloud model. In addition, we find that there are important contributions from the $\Delta\pi$ channel and, especially, from the off-diagonal terms $p\pi^0$-$p\eta_8$ and $p\pi^0$-$p\eta_1$ which together are of the order of 15-20 % of that of the $N\pi$ channel, but with the opposite sign (see Table 1). The contribution of the intermediate vector mesons is very small due to a cancelation between the $n\rho^+$ and the $\Delta\rho$ channels and the cross terms $p\rho^0$-$p\omega_8$ and $p\rho^0$-$p\omega_1$. Kaon loops do not contribute to the proton flavor asymmetry. Table 1 shows that the full four-shell calculation is dominated by the contribution of the ground state intermediate baryons and mesons (0 $\hbar\omega$). Both columns show the same qualitative behavior: dominance of the pion loops with a small negative correction of the order of 10-15 % due to the off-diagonal terms involving π and η pseudoscalar mesons and an almost vanishing contribution from the vector mesons.

Since the unquenched quark model is valid not only for the proton, but for all baryons (ground state or resonance), it is straightforward to calculate the flavor asymmetries of the other octet baryons. For the Σ^+ hyperon and the Ξ^0 cascade particle we find $\mathcal{A}(\Sigma^+) = 0.126$ and $\mathcal{A}(\Xi^0) = -0.001$ [14], respectively. The flavor asymmetries of the remaining octet baryons can be obtained by using the isospin symmetry of the unquenched quark model [14]. For example, the excess of $\bar{d}$ over $\bar{u}$ in the proton is related to the excess of $\bar{u}$ over $\bar{d}$ in the neutron, $\mathcal{A}(p) = -\mathcal{A}(n)$. Similar relations hold for the other octet baryons: $\mathcal{A}(\Sigma^+) = -\mathcal{A}(\Sigma^-)$, $\mathcal{A}(\Xi^0) = -\mathcal{A}(\Xi^-)$ and

Table 2. Relative flavor asymmetries of octet baryons

Model	$\mathcal{A}(\Sigma^+)/\mathcal{A}(p)$	$\mathcal{A}(\Xi^0)/\mathcal{A}(p)$	Ref.
Unquenched CQM	0.833	–0.005	[14]
Octet couplings	0.353	–0.647	[39]
Chiral QM	2	1	[40]
Balance Model	3.083	2.075	[41]

$\mathcal{A}(\Lambda) = \mathcal{A}(\Sigma^0) = 0$. Just as for the proton, the flavor asymmetry of the other octet baryons is expected to be dominated by pion loops, whereas the other contributions are suppressed by the energy denominator in Eq. (1). For the Σ hyperon this is indeed the case, but for the cascade particles the pion loops are suppressed by the value of the $SU(3)$ flavor coupling which is a factor of 5 smaller than that for the proton. Hence for the Ξ hyperons there is no dominant contribution. Since for the Ξ hyperon all contributions are roughly of the same order and small, and moreover some with a positive and others with a negative sign, the value of the flavor asymmetry of the cascade particles is calculated to be small [14].

In Table 2, we show a comparison of some predictions for the flavor asymmetry of the Σ^+ and Ξ^0 hyperons relative to that of the proton. In the unquenched quark model, the flavor asymmetry of the proton is predicted to be of the same order as that of the Σ^+ hyperon and much larger than that of the cascade particle

$$\mathcal{A}(p) \sim \mathcal{A}(\Sigma^+) \gg |\mathcal{A}(\Xi^0)| \ . \tag{8}$$

This behavior is very different from that obtained in the chiral quark model $\mathcal{A}(\Sigma^+) = 2\mathcal{A}(p) = 2\mathcal{A}(\Xi^0)$ [40], the balance model $\mathcal{A}(\Sigma^+) > \mathcal{A}(\Xi^0) > \mathcal{A}(p)$ [41], and the octet model $\mathcal{A}(p) > |\mathcal{A}(\Xi^0)| > \mathcal{A}(\Sigma^+)$ [39]. The values for the chiral quark model and the balance model were taken from [42].

In order to distinguish between the predictions of the different models and to obtain a better understanding of the non-perturbative structure of QCD, new experiments are needed to measure the flavor asymmetry of hyperons. In particular, the flavor asymmetry of charged Σ hyperons can obtained from Drell-Yan experiments using charged hyperon beams on the proton [39] or by means of backward $K^\pm$ electroproduction [43].

2.4. Spin content

The contribution of the quark spins to the spin of the proton can be obtained from the proton spin structure function g_1^p in combination with the neutron and hyperon semileptonic decays [10]. The observation by the European Muon Collaboration that the total quark spin constitutes only a small fraction of the spin of the nucleon [9] sparked an enormous interest in the spin structure of the proton [10]. Recent experiments show that approximately one third of the proton spin is carried by quarks [44, 45], and that the gluon contribution is rather small (either positive or negative) and compatible with zero [46]. This rules out the possibility that most of the missing spin be carried by the gluon and indicates that the origen of the missing spin of the proton has to be attributed to other mechanisms.

In the unquenched quark model, the effect of hadron loops on the fraction of the proton spin carried by the quark (antiquark) spins and orbital angular momentum can be studied in an explicit way [13]. As in other effective models [10], gluonic effects associated with the axial anomaly are not included, and therefore the contribution from the gluons is missing from the outset. The total spin of the proton can then be written as the sum of the contributions from

Table 3. Contribution of quark spins $\Delta\Sigma$ and orbital angular momentum ΔL to the spin of the proton and the Λ hyperon

		CQM	Unquenched QM		
			Valence	Sea	Total
p	$\Delta\Sigma$	1	0.378	0.298	0.676
	$2\Delta L$	0	0.000	0.324	0.324
	$2\Delta J$	1	0.378	0.622	1.000
Λ	$\Delta\Sigma$	1	0.422	0.429	0.851
	$2\Delta L$	0	0.000	0.149	0.149
	$2\Delta J$	1	0.422	0.578	1.000

the quark (and antiquark) spins and orbital angular momentum

$$1 \;=\; 2\Delta J \;=\; \Delta\Sigma + 2\Delta L \; . \tag{9}$$

Table 3 shows that the inclusion of the quark-antiquark pairs has a dramatic effect on the spin content of the proton. Whereas in the CQM the proton spin is carried entirely by the (valence) quarks, in the unquenched calculation the contributions of the valence quark spins, the sea quark spins and the orbital angular momentum to the proton spin are comparable in size and equal to approximately 38, 30 and 32 %, respectively. The importance of orbital angular momentum to the proton spin was discussed many years ago by Sehgal [47] and Ratcliffe [48] in the context of the quark-parton model and, more recently, by Myhrer and Thomas in framework of the bag model [49].

The large contribution of orbital angular momentum is a consequence of the fact that the effects of pion loops for the proton flavor asymmetry and the contribution of orbital angular momentum to the proton spin are identical $\mathcal{A}(p) = \Delta L$ due to the spin and isospin properties [50]. Since in the unquenched calculations both the flavor asymmetry and the orbital angular momentum are dominated by pion loops, this relation is to a good approximation still valid in the UCQM, $\mathcal{A}(p) = 0.151$ and $\Delta L = 0.162$, respectively.

The situation for the quark spins is completely different. In the unquenched calculations, the contributions of valence and sea quarks are given by $\Delta\Sigma_{\rm val} = 0.378$ and $\Delta\Sigma_{\rm sea} = 0.298$, respectively. While the orbital angular momentum arises almost entirely from the $N\pi$ channel, the sea quark spins are dominated by the intermediate vector mesons with a relatively small contribution from the pseudoscalar mesons [17]. Since the convergence of the sum over the contribution of intermediate vector mesons is much slower, the sum was carried out over five complete oscillator shells for both the intermediate baryons and mesons [13].

The experimental data on the spin structure of the proton have raised many questions about the contributions of valence and sea quarks, gluons and orbital angular momentum to the proton spin. In this respect it is of interest to investigate the spin structure of other octet baryons, in particular the Λ hyperon. In most studies, additional assumptions had to be made about the sea quarks in order to get an estimate of its spin content. For example, the assumption that both valence and sea quarks are related by $SU(3)$ flavor symmetry, allows to express the spin content of the Λ hyperon in terms of that of the proton [51, 52, 53] and gives rise to equal contributions of the quark spins $(\Delta\Sigma)_\Lambda = (\Delta\Sigma)_p$. In the unquenched quark model there is no need to make additional assumptions about the nature of the sea. Table 3 shows that the contribution of quark spins for the Λ is larger than that for the proton, $(\Delta\Sigma)_\Lambda > (\Delta\Sigma)_p$, which is a result of $SU(3)$ flavor breaking by the sea quarks.

2.5. *Strangeness*

Even though the nucleon carries no net strangeness, it may have a nonvanishing distribution of strangeness. In 1987 Kaplan and Manohar [54] observed that neutral current experiments could provide information on the strange matrix elements of the nucleon. The parity-violating elastic scattering (PVES) of electrons on nucleon, as suggested by Beck and McKeown in 1988 [55], has demonstrated to provide a powerful tool to determine the strangeness contribution to by combining the weak form factors $G^{Z,p}$ of the proton with the electromagnetic form factors of the nucleon $G^{\gamma,p}$ and $G^{\gamma,n}$ [56, 57, 58, 59, 60, 61]. The strangeness contribution may be extracted by performing a flavor decomposition [55, 62]

$$\begin{aligned}
G^{\gamma,p}_{E,M} &= \frac{2}{3}G^u_{E,M} - \frac{1}{3}\left(G^d_{E,M} + G^s_{E,M}\right) , \\
G^{\gamma,n}_{E,M} &= \frac{2}{3}G^d_{E,M} - \frac{1}{3}\left(G^u_{E,M} + G^s_{E,M}\right) , \\
G^{Z,p}_{E,M} &= (1 - \frac{8}{3}\sin^2\theta_W)G^u_{E,M} + (-1 + \frac{4}{3}\sin^2\theta_W)\left(G^d_{E,M} + G^s_{E,M}\right) .
\end{aligned} \tag{10}$$

Even though the first measurements indicated large and positive values of G^s_M, *e.g.* the SAMPLE collaboration in 1999 found $G^s_M = 0.61 \pm 0.17 \pm 0.21 \pm 0.19$ at $Q^2 = 0.1$ (GeV/c)^2 [63], the more recent values obtained by the HAPPEX, A4, G0 collaborations for G^s_M and G^s_E are much smaller. The most recent measurements show that the contribution of strange quarks to the electric and magnetic form factors is compatible with zero within the experimental errors [58, 59, 60]. Another possibility to determine the strangeness in the proton was suggested by Pate in 2004 [64] by combining experimental data on neutrino scattering with the electromagnetic form factors of the nucleon.

Theoretically, the strangeness of the proton has been addressed in terms of two static observables, the strange magnetic moment and the strange radius. They can be extracted from the behavior of the strange form factors of the nucleon near the origin $Q^2 = 0$ as

$$\begin{aligned}
\mu_s &= e_s\, G^s_M(0) , \\
R^2_s &= -6e_s\, \frac{dG^s_E}{dQ^2}\bigg|_{Q^2=0} .
\end{aligned} \tag{11}$$

We note, that this definition differs from the one in [62] by the electric charge of the strange quark. The values of μ_s and R^2_s can be extracted from the experimental data as follows. In Ref. [56], after measuring G^s_M at $Q^2 = 0.1$ (GeV/c)^2, the strange magnetic moment μ_s is extrapolated by considering the momentum dependence of $G^s_M(Q^2)$ taken from Ref. [65]. The resulting value is $\mu_s = -0.003 \pm 0.097 \pm 0.103 \pm 0.023\ \mu_N$. In Ref. [66] the strange magnetic moment of the proton was obtained from a fit to the complete set of parity violating scattering experimental data to give $\mu_s = -0.04 \pm 0.18 \pm 0.02\ \mu_N$.

Since R^2_s is proportional to the slope of G^s_E in the origen, the evaluation of the strange radius R^2_s requires a measurement of the strange electric form factor at a small value of Q^2 in combination with $G^s_E(0) = 0$ (the nucleon carries no net strangeness). The values of the strange radius obtained in two global analysis of the available experimental data are $R^2_s = 0.002 \pm 0.014$ fm^2 [66] and $R^2_s = -0.006 \pm 0.013$ fm^2 [67].

The first theoretical calculation of the strange form factors of the nucleon, performed by Jaffe in 1989 [68], reported quite large results for G^s_E and G^s_M, thus triggering the interest for this kind of observables. Subsequent theoretical calculations of μ_s and R^2_s, obtained through lattice QCD calculations, hadronic models and effective hadronic theory, vary widely both in absolute value and in the predicted sign $\mu_s(p)$ and $R^2_s(p)$ [62, 69]. The results from recent lattice QCD calculations show small values for both $\mu_s(p)$ and $R^2_s(p)$ [70, 71].

Finally, we discuss the calculation of the strange magnetic moment and radius of the proton in the unquenched CQM. The results are obtained in a calculation involving a sum over intermediate states up to four oscillator shells for both baryons and mesons. In the UCQM formalism, the strange magnetic moment of the proton is given by the expectation value of the operator

$$\vec{\mu}_s = \sum_i \mu_{i,s} \left[2\vec{s}(q_i) + \vec{\ell}(q_i) - 2\vec{s}(\bar{q}_i) - \vec{\ell}(\bar{q}_i) \right] \ . \tag{12}$$

Here $\mu_{i,s}$ is the magnetic moment of the quark i times a projector. In the UCQM the strange magnetic moment of the proton arises from the sea quarks. There is a contribution from the quark spins of the $s\bar{s}$ pair $-0.0004\ \mu_N$, as well as from its orbital motion $-0.0002\ \mu_N$. Both contributions are small and give a total strange magnetic moment $\mu_s = -0.0006\ \mu_N$ [17].

Similarly, the strange radius of the proton is calculated as the expectation value of the operator

$$R_s^2 = \sum_{i=1}^{5} e_{i,s} \left(\vec{r}_i - \vec{R}_{CM} \right)^2 \ , \tag{13}$$

where $e_{i,s}$ is the electric charge of the quark i times a projector on strangeness, and $\vec{r}_i$ and $\vec{R}_{CM}$ are the coordinates of the quark i and the center of mass, respectively. The strange radius of the proton is calculated to be -0.004 fm^2 [17].

In conclusion, the effects of the higher Fock components on the strange magnetic moment and the strange radius of the proton are found to be negligible. Our results are compatible with the latest experimental data and recent lattice calculations.

3. Summary and conclusion

In this contribution, we presented a short review of some general features of an unquenched quark model for baryons in which the effects of sea quarks are taken into account in an explicit form via a $^3P_0$ creation mechanism of the quark-antiquark pairs ($u\bar{u}$, $d\bar{d}$ and $s\bar{s}$). This provides the possibility to address many open problems in baryon structure and spectroscopy. It was shown that, whereas the inclusion of $q\bar{q}$ pairs (or meson loops) does not affect the baryon magnetic moments, it immediately leads to an excess of $\bar{d}$ over $\bar{u}$ in the proton and introduces a sizeable contribution of orbital angular momentum to the spin of the proton. In addition, the contribution of $s\bar{s}$ pairs to the magnetic moment and the radius of the proton was found to be small, in agreement with the latest experimental results and recent lattice calculations.

Even though different models of hadron structure may show similar results for the properties of the proton, often their predictions for the other octet baryons exhibit large variations. Therefore, in order to be able to distinguish between the predictions of different models of hadron structure and to obtain a better understanding of the non-perturbative structure of QCD, new experiments are needed to measure the properties of other octet baryons, such as the Σ and Λ hyperons. In particular, the flavor asymmetry of charged Σ hyperons can obtained from Drell-Yan experiments using charged hyperon beams on the proton [39] or by means of backward $K^{\pm}$ electroproduction [43].

The results for the magnetic moments, the spin and flavor content of octet baryons are very promising and encouraging. The inclusion of the effects of quark-antiquark pairs in a general and consistent way, as suggested here, may provide a major improvement to the constituent quark model which increases considerably its range of applicability.

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