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Reply to Isgur's Comments on Valence QCD

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Abstract

We reply to Nathan Isgur's critique that is directed at some of the conclusions drawn from the lattice simulation of valence QCD, regarding the valence quark model and effective chiral theories.

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1 Introduction

With the goal of understanding the complexity of QCD and the role of symmetry in dynamics, we studied a field theory called Valence QCD (VQCD) [1] in which the Z graphs are forbidden so that the Fock space is limited to the valence quarks. We calculated nucleon form factors, matrix elements, and hadron masses both with this theory and with quenched QCD on a set of lattices with the same gauge background. Comparing the results of the lattice calculations in these two theories, we drew conclusions regarding the $SU(6)$ valence quark model and chiral symmetry. While recognizing the goal of VQCD, Nathan Isgur disagrees on some of the conclusions we have drawn [2].

The foremost objection raised in [2] is to our suggestion that the major part of the hyperfine splittings in baryons is due to Goldstone boson exchange and not one-gluon-exchange (OGE) interactions. The logic of Isgur's objection is that VQCD yields a spectroscopy vastly different from quenched QCD and therefore the structure of the hadrons (to which hyperfine splittings in a quark model are intimately tied) is also suspect so no definite conclusions are possible. To put this into perspective it should be emphasized at the outset that spectroscopy is only one aspect of hadron physics examined in [1]. We have studied the axial and scalar couplings of nucleon in terms of F_A/D_A and F_S/D_S, the neutron to proton magnetic moment ratio μ_n/μ_p, and various form factors. None of these results reveal any pathologies of hadron structure and turn out to be close to the $SU(6)$ relations, as expected. In fact this is what motivated the study of valence degrees of freedom via VQCD.

In Sec. 2 we address specific issues related to spectroscopy in VQCD. Isgur also presented more general agruments against the idea of boson exchange as a contributor to hyperfine effects. A cornerstone of his discussion is the unifying aspect of OGE in a quark model picture. We believe that it is also natural and economical to identify chiral symmetry as the common origin for much of the physics being discussed here. Therefore in Sec. 3 we take the opportunity to sketch out an an effective theory that may serve as a framework to interpret the numerical results of VQCD.

2 Hadron Spectrum

2.1 Meson excitation — $a_1 - \rho$ mass difference

Isgur argues that even with the 'constituent quark' mass shift incorporated into VQCD which lifts the baryon masses by $\sim 3m_{const}$ and the mesons by $\sim 2m_{const}$, it does not restore the $a_1 - \rho$ mass splitting. This is a good point. However, the author's objection that the a_1 does not have an orbital excitation energy relative to the ρ is based on the non-relativistic picture that the axial vector meson has a p-wave excitation as compared to the s-wave description of the vector meson. This is not necessarily true for the relativistic system of light quarks. For example, in a chirally-symmetric world, there are degenerate states due to parity doubling. The pion would be degenerate with the scalar and a_1 would be degenerate with ρ. This is indeed expected at high temperature where the chiral symmetry breaking order parameter, $\langle\bar{\Psi}\Psi\rangle$, goes to zero.

For heavy quarks, we think VQCD should be able to describe the vector – axial-vector meson difference based on the non-relativistic picture. As seen from Figs. 25 and 28 in Ref. [1], from $m_q a = 0.25$ on, the axial-vector meson starts to lie higher than the vector meson. In the charmonium region ($\kappa = 0.1191$), we find the mass difference between them to be 502 ± 80 MeV. Indeed, this is close to the experimental difference of 413 MeV between χ_{c1} and J/Ψ.

In the light quark region the near degeneracy of a_1 and ρ is interpreted as due to the fact that axial symmetry breaking scale, as measured by the condensates $\langle\bar{u}u\rangle$ and $\langle\bar{v}v\rangle$, is small in VQCD as compared to $\langle\bar{\Psi}\Psi\rangle$ in QCD [1]. As a result, there are near parity doublers in the meson spectrum. Note that it is consistent with the observation that dynamical mass generation, another manifestion of spontaneously broken chiral symmetry, is also very small in VQCD.

In the chiral theory, Weinberg's second sum rule gives the relation $m_{a_1} = \sqrt{2} m_\rho$ and the improved sum rule, taking into account of the experimental a_1 and ρ decay constants, gives $m_{a_1} = 1.77 m_\rho$ [3]. This relation is based on chiral symmetry, current algebra, vector meson dominance, and the KSFR relation. These are based on the premise of spontaneous symmetry breaking (SSB). Otherwise, one would expect parity doubling for a_1 and ρ. Thus, to explain the spectrum, we argue that it is sufficient to implement SSB chiral symmetry, not necessarily the p-wave orbital excitation as in the non-relativistic theory. In other words, by restoring the spontaneously broken $SU(3)_L \times SU(3)_R \times U_A(1)$ symmetry to VQCD which has only $U_q(6) \times U_{\bar{q}}(6)$, it is possible to restore the physical mass difference between a_1 and ρ to be consistent with Weinberg's sum rule.

2.2 Hyperfine splittings

As for hyperfine splittings, we have argued that the one-gluon-exchange is not the major source since OGE is still contained in VQCD. Being magnetic in origin, the color-spin interaction is related to the hopping of the quarks in the gauge background in the spatial direction [4]. VQCD does not change this from QCD; the $\vec{\sigma} \cdot \vec{B}$ term is present in the Pauli spinor representation of the VQCD action. Thus, we are forced to draw the conclusion that one-gluon-exchange type of color-spin interaction, i.e. $\lambda_i^c \cdot \lambda_j^c \vec{\sigma}_i \cdot \vec{\sigma}_j$, cannot be responsible for the majority part of the hyperfine splittings between N and Δ and between ρ and π. While we suggested that the Goldstone boson exchange is consistent with the Z-graphs and maybe responsible for the missing hyperfine interaction in the baryons (Fig. 1), it is correctly pointed out by Isgur that there is no such $q\bar{q}$ exchange between the quark and anti-quark in the meson.

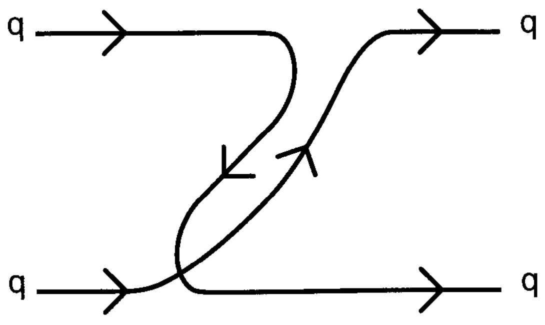

Figure 1: Z-graph between two quarks in a baryon.

One therefore has to consider the possibility that the hyperfine splitting mechanism in the light quark sector is different in mesons from that in the baryons. The numerical results of QCD and VQCD do not, by themselves, reveal the interaction mechanism. A mapping to some model is necessary to make an interpretation. We consider the $SU(3)$ Nambu-Jona-Lasinio NJL model as an example. Starting with a color current-current coupling [5]

$$-9/8G(\bar{\psi}t^a\gamma_\mu\psi)^2, \tag{1}$$

it is convenient to consider Fierz transform to include the exchange terms. The Lagrangian for the color-singlet $q\bar{q}$ meson then takes the following $SU(3)_L \otimes SU(3)_R$ symmetric form with dimension-6 operators for the interaction

$$\begin{aligned} \mathcal{L}_{NJL} &= \bar{\psi}(i\,\partial\!\!\!/ - m_0)\psi + G\sum_i[(\bar{\psi}\frac{\lambda_i}{2}\psi)^2 + (\bar{\psi}\frac{\lambda_i}{2}i\gamma_5\psi)^2] \\ &- G/2\sum_i[(\bar{\psi}\frac{\lambda_i}{2}\gamma_\mu\psi)^2 + (\bar{\psi}\frac{\lambda_i}{2}\gamma_\mu\gamma_5\psi)^2]. \end{aligned} \tag{2}$$

The scalar four-fermion interaction can generate a dynamical quark mass

$$m_d = G\langle\bar{\psi}\psi\rangle. \tag{3}$$

in the mean-field approximation. This is illustrated in Fig. 2. While all the meson masses are lifted up by the dynamical quark masses, the attractive pseudo-scalar four-fermion interaction brings the pion mass back to zero making it a Goldstone boson. The repulsive vector and axial-vector four-fermi interaction makes the ρ, at $\sim$ 770 MeV, slightly higher than twice m_d = 360 MeV. Similarly, the a_1 mass is calculated at $m_{a_1} \simeq 1.2$ GeV, which is not far from the Weinberg's sum rule relation

$m_{a_1} = \sqrt{2} m_\rho$. We see that with one parameter, G, the meson masses can be reasonably described in the NJL model without the $q\bar{q}$ type of meson exchange as in Fig. 1. In addition, current algebra relations such as the Gell-Mann-Oakes-Renner relation

$$m_\pi^2 f_\pi^2 = -\frac{m_u^0 + m_d^0}{2} \langle \bar{u}u + \bar{d}d \rangle, \tag{4}$$

are satisfied. The crucial ingredient here is spontaneous chiral symmetry breaking which is characterized by non-vanishing f_π and quark condensate $\langle \bar{\Psi}\Psi \rangle$, and the existence of Goldstone bosons.

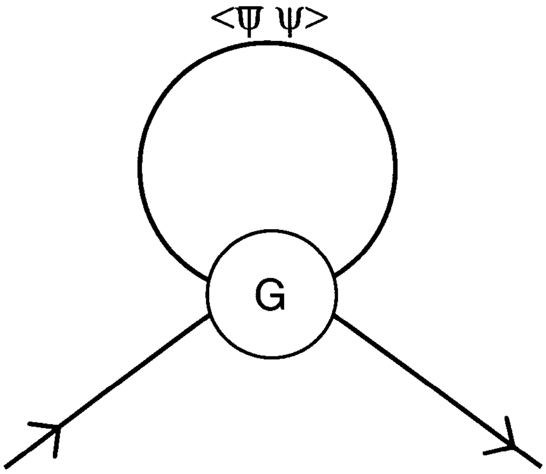

Figure 2: The dynamical mass is generated through the four-fermi interaction with a mean-field approximation.

We should point out that although the color current–current coupling in Eq. (1) is reminiscent of the one-gluon-exchange interaction with the q^2 in the gluon propagator replaced by a cut-off Λ^2 which reflects the short-range nature of the interaction, it is the covariant form for relativistic quarks not the one-gluon exchange potential in the non-relativistic reduction. It is the latter which has been considered as the standard form for hyperfine and fine splittings in the valence quark model.

As illustrated through the NJL model, it is possible to have different mechanisms for hyperfine splitting in the baryons and mesons. In the baryons, the hyperfine splitting *can* be largely due to the meson exchanges between the quarks in the t-channel (Fig. 1); whereas in the mesons, it is the s-channel short-range four-fermion coupling (Fig. 3) that give rise to the hyperfine splittings. Although they appear to be different mechanisms, both of them are based on spontaneously broken chiral symmetry.

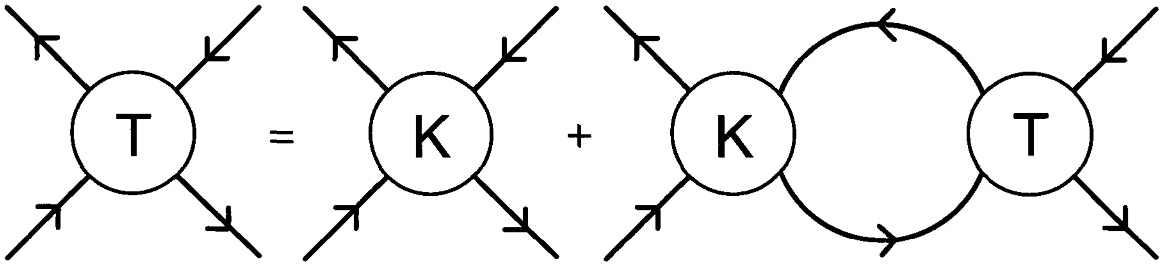

Figure 3: Bethe-Salpeter equation for the meson T-matrix.

The author displayed the spectrum ranging from heavy–heavy mesons $(b\bar{b}, c\bar{c})$ to light–light mesons ($s\bar{s}$ and isovector light quarkonia) in Fig. 4 of his paper [2] which suggests a smooth trend as a function of the quark mass and argues for a universal OGE hyperfine interaction with a strength proportional to $1/m_Q^2$. We have pointed out in our VQCD paper [1] from the outset that we believe the heavy–heavy mesons are well described by a non-relativistic potential model including the OGE; this is supported by the lattice calculations [6, 7, 8, 9]. It is the validity of OGE in the light–light mesons sector that we question. What have been neglected in Fig. 4 of Ref. [2] are the 1^{++} and 0^{++} mesons. Had these been put in, one would have seen that $a_0(1430)$ lies higher than $a_1(1260)$ and $a_2(1320)$. This ordering between 1^{++} and 0^{++} mesons is reversed from that in the charmonium family where $\chi_{c1}(3510)$ lies higher than $\chi_{c0}(3415)$. There is an indication from the lattice calculation that this cross-over occurs at about the strange mass region [10]. As far as we know, this pattern of order reversal in the fine splitting as the quark mass becomes light cannot be accommodated in the OGE picture.

Also shown in Fig. 5(a) of Ref. [2] are the hyperfine splittings of the ground state heavy-light mesons. We concur that the splittings of $B^*(5325) - B(5279)$ and $D^*(2010) - D(1869)$ are quite consistent with the matrix elements of the hyperfine interaction $\vec{\sigma}_Q \cdot \vec{B}/2m_Q$ and that it clearly demonstrates the $1/m_Q$ behavior of the heavy quark. We never questioned the relativistic corrections of the heavy quarks. It is with light quarks that we think OGE has problems. For example, consider the similar splittings for the heavy–light mesons with different light quarks. The mass difference between $D^*(2010)$ and $D(1869)$ is $140.64 \pm 0.10\,\mathrm{MeV}$. This is practically the same as that between $D_s^*(2110)$ and $D_s(1969)$ which is $143.9 \pm 0.4\,\mathrm{MeV}$. There is no indication of the $1/m_q$ dependence on the light quark mass as required by the OGE potential. Similarly we find that $m_{B^*} - m_B = 45.78 \pm 0.35\,\mathrm{MeV}$ is identical to $m_{B_s^*} - m_{B_s} = 47.0 \pm 2.6\,\mathrm{MeV}$. Again, there is no $1/m_q$ dependence.

3 Effective Theory for Both Mesons and Baryons

Besides commenting on the spectroscopy specific to VQCD, Isgur also questioned the meson exchange picture on more general grounds. Since this issue has been raised, we take the opportunity to extend our discussion although it is outside the scope of VQCD.

Perhaps the most serious challenge to the meson exchange picture in the baryons is the possibility of meson exchanges between the quark and anti-quark in the iso-singlet meson. It is pointed out by Isgur that the annihilation diagram depicted in Fig. 6 in Ref. [2] in terms of the quark lines is OZI suppressed in QCD. We should add that it is $O(1/N_c^2)$ suppressed as compared to one-pion-exchange between the quark pairs in the baryon (Fig. 1) in the large N_c analysis. On the other hand, interpreting this as a Goldstone boson exchange between the quark and anti-quark in the iso-singlet mesons, such as a kaon exchange, leads to large $\omega - \phi$ mixing. How does one reconcile the apparent contradiction? The short answer is that there is no such process in the effective theory of mesons. It is inconsistent, within the renormalization group approach to effective theories, to consider this QCD annihilation process as a meson exchange between the quark and anti-quark in the meson. To see this, we shall use the NJL model as an illustration.

3.1 Bosonization

We shall follow the example given by U. Vogl and W. Weise [5] for a simple $U(1)_V \otimes U(1)_A$ symmetric Lagrangian

$$\mathcal{L} = \bar{\psi}(i \not{\partial} - m_0)\psi + G[(\bar{\psi}\psi)^2 + (\bar{\psi} i \gamma_5 \psi)^2]. \tag{5}$$

To bosonize this theory, one needs to integrate out the fermions. One can follow the Hubbard-Stratonovich transformation [11] by introducing Gaussian auxiliary boson fields σ and π with the Lagrangian $-\mu^2/2(\sigma^2+\pi^2)$ and the partition function becomes

$$\mathcal{Z} = \mathcal{N} \int \mathcal{D}\sigma \mathcal{D}\pi \mathcal{D}\bar{\psi} \mathcal{D}\psi \, e^{i \int d^4x \bar{\psi}[i \not{\partial} - m_0 - \mu\sqrt{2G}(\sigma + i\gamma_5 \pi)]\psi - \mu^2/2(\sigma^2+\pi^2)}, \tag{6}$$

after a linear shift of the fields σ and π. Note here, the σ and π are the auxiliary fields with no kinetic terms.

At this stage, one can integrate the fermion field with the quadratic action to obtain the fermion determinant. This gives an effective action with the $\mathrm{tr}\ln M$ Lagrangian, where M is the inverse quark propagator between the square brackets in Eq. (6). Expanding the $\mathrm{tr}\ln M$ to the second order in the derivative ∂_μ for the low energy long wavelength approximation, the effective Lagrangian becomes

$$\mathcal{L}_{eff}(\sigma,\pi) = \frac{1}{2}[(\partial_\mu\sigma)^2 + (\partial_\mu\pi)^2] - \frac{1}{2}m_\pi^2\pi^2 - \frac{1}{2}m_\sigma^2\sigma^2$$

$$-\frac{2m^2}{f_\pi}\sigma(\sigma^2+\pi^2)-\frac{m^2}{2f_\pi}\pi(\sigma^2+\pi^2)^2, \qquad (7)$$

where $m = m_0 + \mu\sqrt{2G}\langle\sigma\rangle = m_0 - 2G\langle\bar{\Psi}\Psi\rangle$. Besides giving π and σ masses as the physical mesons, it also gives the explicit meson-meson couplings.

Thus, to construct an effective theory below the meson confinement scale, which corresponds to the chiral symmetry breaking scale $\Lambda_\chi = 4\pi f_\pi \simeq 1\,\mathrm{GeV}$ as we shall see later, one can take the following equivalent approaches: In the first one, one can introduce higher dimensional operators like $(\bar{\psi}\psi)^2, (\bar{\psi}i\gamma_5\psi)^2, (\bar{\psi}\gamma_\mu\psi)^2, (\bar{\psi}\gamma_\mu\gamma_5\psi)^2$ to the usual QCD Lagrangian and tune the couplings to match to QCD above Λ_χ. Many improved lattice actions are constructed this way in order to do numerical simulation at a lower lattice cut-off or larger lattice spacing in order to save computer time [12]. In the second approach, one can introduce auxiliary fields π, σ, ρ, a_1, etc. to replace the four-fermion operators with couplings to fermion bilinears and multi-auxiliary-field couplings as in Eq. (6). This form has been considered in lattice QCD simulations [13, 14] to control the singular nature of the massless Dirac operator. The third approach is to bosonize the theory by integrating out the fermion fields and performing derivative expansion of the $\mathrm{tr}\ln M$ action from the fermion loop as in Eq. (7). An extensive and successful model of this kind has been developed [15] where ρ is predicted to be close to the experimental value and a_1 mass is related to the ρ via the modified Weinberg sum rule [3]. VMD and the KSFR relation are satisfied. In addition, the pion form factor, $\pi\pi$ scattering, and a host of meson decays are all in good agreement with the experiments.

We see that in none of the above three equivalent approaches is there a coupling between the quark and physical mesons. Thus, there is no OPE between the quark-anti-quark pair in the meson. Since one is below the meson confinement scale Λ_χ, the meson fields are the relevant degrees of freedom. Once one integrates out the fermion fields in the meson in favor of the physical meson fields, it would be inconsistent to construct a meson model with couplings between quarks and physical mesons. Of course, this does not preclude short-range couplings between $u\bar{u}$, $d\bar{d}$ and $s\bar{s}$ in the s-channel to resolve the $U_A(1)$ anomaly and give η' a large mass via the contact term of the topological susceptibility [17].

Then how does one justify the σ - quark model that one proposes as an effective theory for the baryons? To realize this one has to make a distinction between the meson and the baryon.

3.2 Chiral effective theory for baryons

In view of the observation that mesons have form factors in the monopole form and baryons have form factors in the dipole form, the πNN form factor is much softer than the $\rho\pi\pi$ form factor, we suggest that the confinement scale of quarks in the baryon l_B is larger than l_M – the confinement scale between the quark and anti-quark in the meson; that is,

$$l_B > l_M. \tag{8}$$

This is consistent with the large N_c approach where the mesons are treated as point-like fields and the baryons emerge as solitons with a size of order unity in N_c. Taking the l_M from the $\rho\pi\pi$ form factor gives $l_M \sim 0.2\,\mathrm{fm}$. This is very close to the chiral symmetry breaking scale set by $\Lambda_\chi = 4\pi f_\pi$. We consider them to be the same, i.e. below Λ_χ, operators of mesons fields become relevant operators. As for the baryon confinement scale, we take it to be the size charactering the meson-baryon-baryon form factors. Defining the meson-baryon-baryon form factors from taking out the respective meson poles in the nucleon pseudoscalar, vector, and axial form factors (see Fig. 17 in Ref. [1]), we obtain $l_B \sim 0.6 - 0.7\,\mathrm{fm}$. This satisfies the inequality in Eq. (8). Thus, in between these two scales l_M and l_B, one could have coexistence of mesons and quarks in a baryon.

We give an outline to show how to construct a chiral effective theory for baryons. In the intermediate length scale between l_M and l_B, one needs to separate the fermion field into a long-range one and a short-range one

$$\psi = \psi_L + \psi_S, \tag{9}$$

where ψ_L/ψ_S represent the infrared/ultraviolet part of the quark field with momentum components below/above $1/l_M$ or Λ_χ. We add to the ordinary QCD Lagrangian irrelevant higher dimension operators with coupling between bilinear quark fields and auxiliary fields as given in Ref. [15]. However, we interpret these quark fields as the short-range ones, i.e. ψ_S and $\bar{\psi}_S$. Following the procedure in Ref. [15], one can integrate out the ψ_S and $\bar{\psi}_S$ fields and perform the derivative expansion to bosonize the short-range part of the quark fields. This leads to the Lagrangian with the following generic form:

$$\mathcal{L}_{\chi QCD} = \mathcal{L}_{QCD'}(\bar{\psi}_L, \psi_L, A_\mu^L) + \mathcal{L}_M(\pi, \sigma, \rho, a_1, G, ...) + \mathcal{L}_{\sigma q}(\bar{\psi}_L, \psi_L, \pi, \sigma, \rho, a_1, G, ...). \tag{10}$$

$\mathcal{L}_{QCD'}$ includes the original form of QCD but in terms of the quark fields $\bar{\psi}_L, \psi_L$, and the long-range gauge field A_μ^L with renormalized couplings; it also includes higher-order covariant derivatives [18]. $\mathcal{L}_M$ is the meson effective Lagrangian, e.g. the one derived by Li [15] which should include the glueball field G. Finally, $\mathcal{L}_{\sigma q}$ gives the coupling between the $\bar{\psi}_L, \psi_L$, and mesons. As we see, in this intermediate scale, the quarks, gluons, and mesons coexist and meson fields *do* couple to the quark fields, but it is ψ_L that the mesons couple to, *not* ψ_S. Going further down below the baryon confinement scale $1/l_B$, one can integrate out $\bar{\psi}_L, \psi_L$ and A_μ^L, resulting in an effective Lagrangian $\mathcal{L}(\bar{\Psi}_B, \Psi_B, \pi, \sigma, \rho, a_1, G, ...)$ in terms of the baryon and meson fields [19]. This would correspond to an effective theory in the chiral perturbation theory.

Fig. 4 is a schematic illustration of effective theories partitioned by the two scales of l_M and l_B. We should point out that although we adopt two scales here, they are

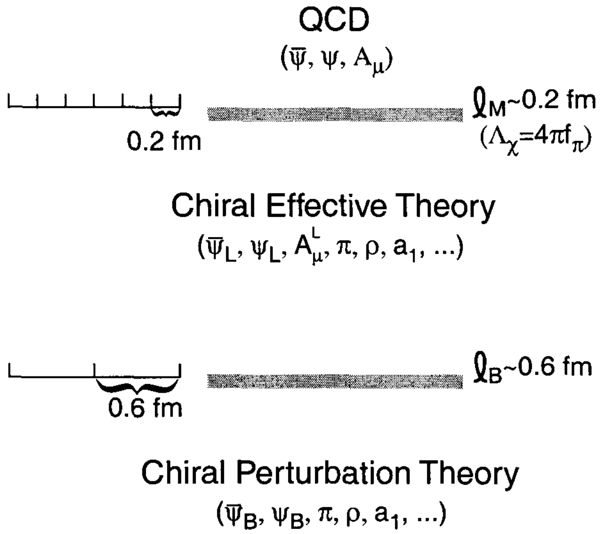
QCD
$(\bar{\psi}, \psi, A_\mu)$
0.2 fm
ℓ_M~0.2 fm
$(\Lambda_\chi=4\pi f_\pi)$
Chiral Effective Theory
$(\bar{\psi}_L, \psi_L, A^L_\mu, \pi, \rho, a_1, ...)$
0.6 fm
ℓ_B~0.6 fm
Chiral Perturbation Theory
$(\bar{\psi}_B, \psi_B, \pi, \rho, a_1, ...)$

Figure 4: A schematic illustration of the the two-scale delineation of the effective theories. The shaded bars mark the positions of the cutoff scales l_M and l_B separating different effective theories.

distinct from those of Manohar and Georgi [20]. In the latter, the σ – quark model does not make a distinction between the baryons and mesons. As such, there is an ambiguity of double counting of mesons and $q\bar{q}$ states. By making the quark-quark confinement length scale l_B larger than the quark–anti-quark confinement length scale l_M, one does not have this ambiguity. The outline we give here is a systematic way of constructing the effective theory at appropriate scales following Wilson's renormalization group approach [21, 22].

We see from Fig. 5 that the $\mathcal{L}_{\sigma q}$ part of the effective chiral theory in Eq. (10) is capable of depicting meson dominance (Fig. 5(a)), the quark Z-graphs and cloud degree of freedom via the meson exchange current (Fig. 5(b)), and the sea quarks in the disconnected insertion via the meson loop (Fig. 5(c)) in a baryon. These correspond to the dynamical quark degrees of freedom in QCD as we alluded to in the study of baryon form factors in the path-integral formulation [1]. On the other hand, when one considers the chiral perturbation theory at energy lower than $1/l_B \sim 300\,\mathrm{MeV}$, the dressing of baryons with meson clouds (Fig. 6) no longer distinguishes the cloud-quarks from the sea-quarks.

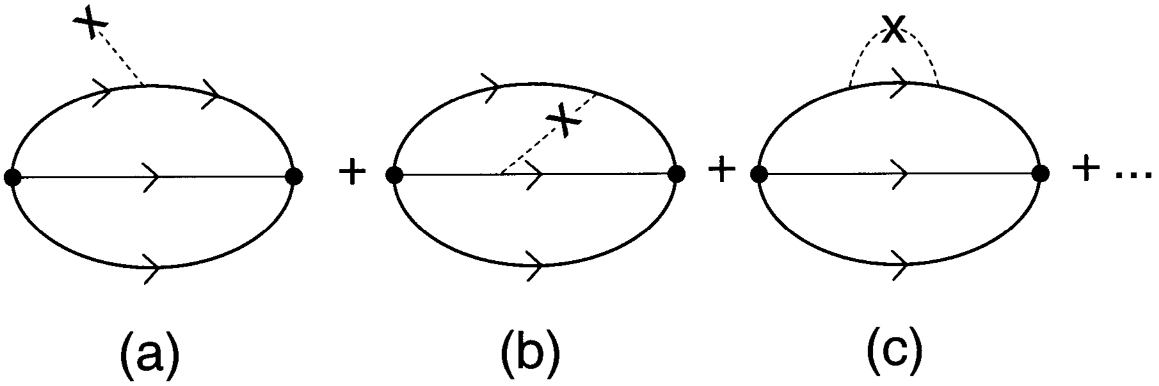

Figure 5: The σ-quark model description of (a) meson dominance, (b) cloud quarks via meson exchange current, and (c) sea quarks via the meson loop.

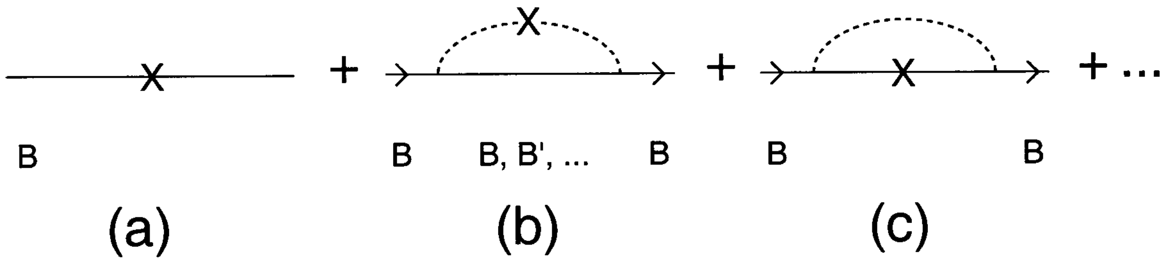

Figure 6: (a) Direct baryon contribution and (b) & (c) meson loop contribution in the chiral perturbation theory.

One important aspect of constructing effective theories based on the renormalization group is that chiral symmetry and other symmetries of the theory should be preserved as one changes the cut-off so as to ensure universality.

As we see from the above construction of effective chiral theories, there is no large OZI-violating meson exchange between the quark and anti-quark in an iso-singlet meson. The problem that Isgur perceives for the meson exchange in the iso-singlet meson is simply not there.

4 Conclusions

As stressed at the beginning, hadron spectroscopy is only one of the many facets of hadron physics. At low energies, there is a lot of evidence that chiral symmetry is playing a crucial role, for example, in the $\pi\pi$ scattering, the Goldberger-Treiman relation, the Gell-Mann-Oakes-Renner relation, the Kroll-Ruderman relation, the KSRF relation, and Weinberg sum rules.

As far as light hadrons are concerned, it is natural to expect chiral symmetry to play a role in spectroscopy also. For many years, various chiral models have been successful in describing the pattern of masses in the meson sector in addition to scattering and decays. Now it appears that the chiral quark picture can give a reasonable explanation of the baryon spectroscopy as well as structure.

Finally, we echo Isgur's comment 'while qQCD describes both the $\rho - \pi$ and $\Delta - N$ splittings, they are both poorly described in vQCD. It would be natural and economical to identify a common origin for these problems.' It is proposed that chiral symmetry is this common origin, albeit it may have different dynamical realization in mesons and baryons. We suggest it is chiral symmetry that is the essential physics multilated in VQCD and that this is manifested by the suppression of dynamical mass generation, approximate parity doublets, the incorrect $U(6)$ symmetry and the disappearance of hyperfine splittings. We expect that effective chiral theories or models that incorporate the spontaneously broken $SU(3)_L \times SU(3)_R \times U_A(1)$ symmetry will have the relevant dynamical degrees of freedom necessary to delineate the structure and spectroscopy of both mesons and baryons of light quarks at a scale below $\sim 1\,\text{GeV}$.

5 Acknowledgment

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