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THE SPIN OF THE PROTON *

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I argue that our response to the spin crisis should not be to abandon the naïve quark model baby, but rather to allow it to mature. In particular, I advocate dressing the baby in $q\bar{q}$ pairs, first showing that this can be done without compromising the naïve quark model's success with either spectroscopy or the OZI rule. Finally, I show that despite their near invisibility elsewhere, pairs do play an important role in the proton's spin structure by creating an antipolarized $q\bar{q}$ sea. In the context of an explicit calculation I demonstrate that it is plausible that the entire "spin crisis" arises from this effect.

1 Introductory Remarks

Sometimes, it seems to me, we get a little carried away with the beauty and precision of deep inelastic scattering as a probe of the quark-gluon structure of the proton. It is undeniably satisfying to be able to make rigorous statements about strongly interacting systems: it is such a rare experience. However, deep inelastic scattering can, *via* perturbative QCD, only tell us "what is there". It is thus only a prelude to understanding *via* Strong QCD [1] "why it is there".

There are those who would argue that it is sufficient to compare the results of experiments against lattice simulations of QCD. I certainly agree that this is our principal route toward precision tests of Strong QCD. However, I would argue that such comparisons cannot act as a substitute for *understanding* QCD, and that it would be a great failure of physics if we were to forfeit thinking about QCD to our computers.

QCD is undeniably very complex, so if we are to understand it, we will clearly have to find some way to simplify it. This is, of course, not new to QCD. Superconductivity is also intractable as a theory of 10^{23} crytalline nuclei and 10^{23} electrons, but it can be understood once the right low energy effective degrees of freedom - - - Cooper pairs and phonons - - - are identified. An example closer to home is nuclear physics: nuclear structure may be understood very well in terms of the nucleonic low energy degrees of freedom interacting *via* few-body effective potentials. Similarly, I believe that the central issue in Strong QCD is to identify the correct low energy degrees of freedom at the "quark model" scale $\mu_{QM} \sim 1$ GeV.

In this talk I will argue that there are many good reasons to believe that the constituent quark model *is* a good starting point in this quest. When extended *via* the flux tube model, wherein the gluonic degrees of freedom are subsumed into flux tubes, the constituent quark model can be mapped onto QCD in the large N_c limit [2]. While adequate for many purposes, this quark model must be further extended by "unquenching" (a $1/N_c$ effect) if it is to provide a satisfactory qualitative picture of Strong QCD. That is, I suggest that dynamical $q\bar{q}$ pairs are the key missing ingredient of the constituent quark model. In particular, I will show in an explicit model how one can "unquench" the quark model without spoiling its spectroscopic successes or ruining the OZI rule [3]. At the same time, we will see that while each light quark flavor may make a relatively small negative contribution to the net proton spin of order $1/N_c$, N_f such contributions can account for the observed "spin crisis".

In particular, in the following, I will present arguments in support of a very conservative solution to the spin crisis: that the valence quarks have their expected "not-so-naïve"

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polarization $\Sigma_v \simeq 0.75$, but that sea quarks are polarized with $\Sigma_{sea} \simeq -0.15$ per flavor of light quark (see Figure 1). To make this case I will describe an explicit model in which the degrees of freedom of the flux tube model are expanded by allowing for the creation of $q\bar{q}$ loops ("unquenching"). Then I will describe the results of this model for Δs and for the strange electric and magnetic form factors of the proton, and speculate on an extension to $u\bar{u}$ and $d\bar{d}$ pairs which would resolve the "spin crisis".

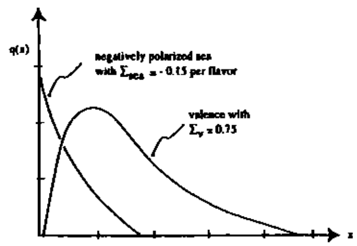

FIG. 1. A schematic of the polarization of the valence and the sea quarks.

2 Dressing the Baby

While it provides a good description of low-energy strong interaction phenomena, the constituent quark model appears to be inconsistent with many fundamental characteristics of QCD. Foremost among these inconsistencies is a "degree of freedom problem": the quark model declares that the low energy spectrum of QCD is built from the degrees of freedom of spin-1/2 fermions confined to a $q\bar{q}$ or qqq system. Thus, for mesons the quark model predicts - - - and we seem to observe - - - a "quarkonium" spectrum. In the baryons it predicts - - - and we seem to observe - - - the spectrum of the two relative coordinates of three spin-1/2 degrees of freedom.

These quark model degrees of freedom are to be contrasted with the most naïve interpretation of QCD which would lead us to expect a low energy spectrum exhibiting 36 quark and antiquark degrees of freedom (3 flavors $\times$ 2 spins $\times$ 3 colors for particle and antiparticle), *and* 16 gluon degrees of freedom (2 spins $\times$ 8 colors). Less naïve pictures exist, but none evade the first major "degree of freedom problem": the gluonic degrees of freedom appear to be missing from the low energy spectrum. This issue, being one of the most critical in Strong QCD, is being addressed by many theoretical and experimental programs.

The second major "degree of freedom problem" has to do with $q\bar{q}$ pair creation. *A priori*, one would expect pair creation to be so probable that a valence quark model would fail dramatically, while empirically pair creation is suppressed: the observed hadronic spectrum is dominated by narrow resonances, while the naïve picture would predict resonances with widths Γ comparable to their masses m.

The dressing of the constituent quark model with glue in the context of the flux tube model [4] has been widely discussed [5]. Here I will focus be on the effects of the $q\bar{q}$ sea.

2.1 Unquenching the Quark Model: Overview

There are three puzzles associated with the nature and importance of $q\bar{q}$ pairs in low energy hadron structure:

1) the origin of the apparent valence structure of hadrons (since even as $N_c \to \infty$, Z-graphs would produce pairs unless the quarks were heavy),

2) the apparent absence of unitarity corrections to naïve quark model spectroscopy, despite one's expectation of mass shifts $\Delta m \sim \Gamma$ (where Γ is a typical hadronic width), and

3) the systematic suppression of OZI-violating amplitudes A_{OZI} relative to one's expectation (from unitarity) that $A_{OZI} \sim \Gamma$.

In this section I will describe the solutions I see to these puzzles.

2.2 The Origin of the Valence Approximation

A weak form of the valence approximation seems to emerge from the large N_c limit [2] in the sense that diagrams in which only valence quark lines propagate through hadronic two-point functions dominate as $N_c \to \infty$. However, this dominance does not seem to correspond to the usual valence approximation since the Z-graph pieces of such diagrams will produce a $q\bar{q}$ sea.

Consider, however, the Dirac equation for a single light quark interacting with a static color source (or a single light quark confined in a bag). This equation represents the sum of a set of Feynman graphs which also include Z-graphs, but the effects of those graphs is captured in the lower components of the single-particle Dirac spinor. *I.e.*, such Z-graphs correspond to relativistic corrections to the quark model. That such corrections are important in the quark model has been known for a long time [6]. For us the important point is that while they have quantitative effects on quark model predictions like g_A, they do not qualitatively change the single-particle nature of the spectrum of the quark of our example, nor would they qualitatively change the spectrum of $q\bar{q}$ or qqq systems. Note that this interpretation is consistent with the fact that Z-graph-induced $q\bar{q}$ pairs do *not* correspond to the usual partonic definition of the $q\bar{q}$ sea since Z-graphs vanish in the infinite momentum frame. Thus the $q\bar{q}$ sea of the parton model is also associated with the $q\bar{q}$ loops of unquenched QCD.

2.3 The $\Delta m << \Gamma$ Problem

Consider two resonances which are separated by a mass gap δm in the narrow resonance approximation. In general we would expect that departures from the narrow resonance approximation, which produce resonance widths Γ, ought also to produce mass shifts Δm of order Γ. Yet even though a typical hadronic mass spectrum is characterized by mass gaps

δm of order 500 MeV, and typical hadronic widths are of order 250 MeV, this does not seem to happen.

A simple resolution of this puzzle has been proposed [7]. As discussed in Section 4.1, in the flux tube model [4], the quark potential model arises from an adiabatic approximation to the gluonic degrees of freedom embodied in the flux tube. At short distances where perturbation theory applies, the effect of N_f types of light $q\bar{q}$ pairs is (in lowest order) to shift the coefficient of the Coulombic potential from $\alpha_s^{(0)}(Q^2) = \frac{12\pi}{33 ln(Q^2/\Lambda_0^2)}$ to $\alpha_s^{(N_f)}(Q^2) = \frac{12\pi}{(33-2N_f) ln(Q^2/\Lambda_{N_f}^2)}$. The net effect of such pairs is thus to produce a *new* effective short distance $Q\bar{Q}$ potential. Similarly, when pairs bubble up in the flux tube (*i.e.*, when the flux tube breaks to create a $Q\bar{q}$ plus $q\bar{Q}$ system and then "heals" back to $Q\bar{Q}$), their net effect is to cause a shift $\Delta E_{N_f}(r)$ in the ground state gluonic energy which in turn produces a new long-range effective $Q\bar{Q}$ potential [8].

It has indeed been shown [7] that the net long-distance effect of the bubbles is to create a new string tension b_{N_f} (*i.e.*, that the potential remains linear). Since this string tension is to be associated with the observed string tension, after renormalization *pair creation has no effect on the long-distance structure of the quark model in the adiabatic approximation.* Thus the net effect of mass shifts from pair creation is much smaller than one would naïvely expect from the typical width Γ: such shifts can only arise from nonadiabatic effects. For heavy quarkonium, these shifts can in turn be associated with states which are strongly coupled to nearby thresholds.

It should be emphasized that it was necessary to sum over very large towers of $Q\bar{q}$ plus $q\bar{Q}$ intermediate states to see that the spectrum was only weakly perturbed (after unquenching and renormalization). In particular, no simple truncation of the set of meson loop graphs can reproduce such results.

2.4 The Survival of the OZI Rule

There is another puzzle of hadronic dynamics which is reminiscent of this one: the success of the OZI rule [3]. A generic OZI-violating amplitude A_{OZI} can also be shown to vanish like $1/N_c$. However, there are several unsatisfactory features of this "solution" to the OZI mixing problem [9]. Consider ω-ϕ mixing as an example. This mixing receives contributions from both "hairpin graphs" and from the virtual hadronic loop process $\omega \to K\bar{K} \to \phi$, both steps of which are OZI-allowed, and each of which scales with N_c like $\Gamma^{1/2} \sim N_c^{-1/2}$ (see Figure 2). The large N_c result that this OZI-violating amplitude behaves like N_c^{-1} is thus not peculiar to large N_c: it just arises from "unitarity" in the sense that the real and imaginary parts of a generic hadronic loop diagram will have the same dependence on N_c. The usual interpretation of the OZI rule in this case - - - that "double hairpin graphs" are dramatically suppressed - - - is untenable in the light of these OZI-allowed loop diagrams. They expose the deficiency of the large N_c argument since $A_{OZI} \sim \Gamma$ is *not* a good representation of the OZI rule. (Continuing to use ω-ϕ mixing as an example, we note that $m_\omega - m_\phi$ is numerically comparable to a typical hadronic width, so the large N_c result would predict an ω-ϕ mixing angle of order unity in contrast to the observed pattern of very weak mixing which implies that $A_{OZI} << \Gamma << m$.)

Unquenching the quark model thus endangers the naïve quark model's agreement with the OZI rule. It has been shown [10] how this disaster is naturally averted in the flux tube

model through a "miraculous" set of cancellations between mesonic loop diagrams consisting of apparently unrelated sets of mesons (*e.g.*, the $K\bar{K}$, $K\bar{K}^* + K^*\bar{K}$, and $K^*\bar{K}^*$ loops tend to strongly cancel against loops containing a K or K^* plus one of the four strange mesons of the $L = 1$ meson nonets).

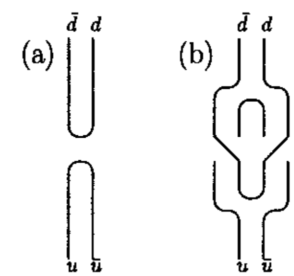

Figure 2: (a) OZI-violation in a meson propagator by "pure annihilation", corresponding to a disconnected double-hairpin diagram. (b) A different time ordering of the same Feynman graph gives an OZI-violating loop diagram via two OZI-allowed amplitudes.

Of course the "miracle" occurs for a good reason. In the flux tube model, where pair creation occurs in the $^3P_0$ state, the overlapping double hairpin graphs which correspond to OZI-violating loop diagrams (see Fig. 2), cannot contribute in a closure-plus-spectator approximation since the 0^{++} quantum numbers of the produced (or annihilated) pair do not match those of the initial and final state for any established nonet. In fact [10] this approximation gives zero OZI violation in all but the (still obscure) 0^{++} nonet. In addition, corrections to the closure-plus-spectator approximation are small, so that the observed hierarchy $A_{OZI} << \Gamma$ is reproduced.

I emphasize once again that such cancellations require the summation of a very large set of meson loop diagrams with cancellations between what are apparently unrelated sets of intermediate states, and that I know of no low energy hadronic effective theory that could reproduce such behaviour.

2.5 Some Comments

We believe the preceding discussion strongly suggests that models [11] which have not addressed the effects of unquenching on spectroscopy and the OZI rule should be viewed very skeptically as models of the effects of the $q\bar{q}$ sea on hadron structure: large towers of mesonic loops are required to understand how quarkonium spectroscopy and the OZI rule survive once strong pair creation is turned on. In particular, while pion and kaon loops (which tend to break the closure approximation due to their exceptional masses) have a special role to play, they cannot be expected to provide a reliable guide to the physics of $q\bar{q}$ pairs.

3 A Pair Creation Model for the Strangeness of the Proton

The following discussion of the strangeness content of the proton will be based on the quark-level process shown in Fig. 3(b). The main new feature of the calculation on which this discussion is based [12] is a sum over a *complete set* of strange intermediate states, rather than just a few low-lying states. As explained above, this is *necessary* for consistency with the OZI rule and the success of quark model spectroscopy.

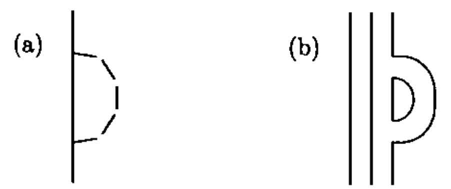

Figure 3: A meson loop correction to a baryon propagator, drawn at (a) the hadronic level, and (b) the quark level.

The lower vertex in Fig. 3(b) arises when $q\bar{q}$ pair creation perturbs the initial nucleon state vector so that, to leading order in pair creation,

$$|p\rangle \rightarrow |p\rangle + \sum_{Y^*K^*\ell S} \int q^2 dq \; |Y^*K^*q\ell S\rangle \frac{\langle Y^*K^*q\ell S\,|h_{q\bar{q}}|\,p\rangle}{M_p - E_{Y^*} - E_{K^*}} \,, \tag{1}$$

where $h_{q\bar{q}}$ is a quark pair creation operator, Y^* (K^*) is the intermediate baryon (meson), q and ℓ are the relative radial momentum and orbital angular momentum of Y^* and K^*, and S is the sum of their spins. Of particular interest is $s\bar{s}$ pair creation by the pair creation operator $h_{s\bar{s}}$, which will generate non-zero expectation values for strangeness observables:

$$\langle O_s\rangle = \sum_{\substack{Y^*K^*\ell S\\ Y^{*\prime}K^{*\prime}\ell' S'}} \int q^2 dq\, q'^2 dq' \frac{\left\langle p\,|h_{s\bar{s}}|\,Y^{*\prime}K^{*\prime}q'\ell' S'\right\rangle}{M_p - E_{Y^{*\prime}} - E_{K^{*\prime}}}$$

$$\times\left\langle Y^{*\prime}K^{*\prime}q'\ell' S'\,|O_s|\,Y^*K^*q\ell S\right\rangle \frac{\langle Y^*K^*q\ell S\,|h_{s\bar{s}}|\,p\rangle}{M_p - E_{Y^*} - E_{K^*}}. \tag{2}$$

The derivation of this simple equation, including the demonstration that it is gauge invariant, is straightforward [12]. We will be considering the cases $O_s = \Delta s$, R_s^2, and μ_s, where Δs the matrix element $\langle p,+|\bar{s}\gamma_3\gamma_5 s|p,+\rangle$ which measures the contribution of strange pairs to the net spin of the proton, and where R_s^2 and and μ_s are the strangeness radius and magnetic moments to be defined more precisely below. The value of Δs can also be associated (via

small scale-dependent QCD radiative corrections) with the contribution of strange quarks to the deep inelastic spin-dependent structure functions.

To calculate the $p \rightarrow Y^*K^*$ vertices in Eq. (1), the flux-tube-breaking model was used. This model, which reduces to the well-known $^3P_0$ decay model in a well-defined limit, had its origin in applications to decays of mesons [13,14] and baryons [16]. The model assumes that a meson or baryon decays when a chromoelectric flux tube breaks, creating a constituent quark and antiquark on the newly exposed flux tube ends. The pair creation operator is taken to have $^3P_0$ quantum numbers:

$$h_{q\bar{q}}(t,\mathbf{x}) = \gamma_0 \left(\frac{3}{8\pi r_q^2}\right)^{3/2} \int d^3z \; \exp\left(-\frac{3z^2}{8r_q^2}\right) \; q^\dagger(t,\mathbf{x}+\frac{\mathbf{z}}{2})\alpha\cdot\nabla q(t,\mathbf{x}-\frac{\mathbf{z}}{2}) \,. \tag{3}$$

The dimensionless constant γ_0 is the intrinsic pair creation strength, a parameter which we fit to the $\Delta \rightarrow N\pi$ width. The operator (3) creates *constituent* quarks, hence the pair creation point is smeared out by a gaussian factor whose width, r_q, is another parameter of the model. The parameter r_q is constrained by meson decay data to be approximately 0.25 fm [7,10].

Once an $s\bar{s}$ pair is created, the decay proceeds by quark rearrangement, as shown in Fig. 4. The $p \rightarrow Y^*K^*$ decay amplitude of the first of Figs. 4 may be written as

$$\langle Y^*K^*|h_{s\bar{s}}|p\rangle = \gamma_0\, \vec{\Sigma}\cdot\vec{I}, \tag{4}$$

where $\vec{\Sigma}$ is a spin overlap which can be expressed in terms of the baryon and meson spin wavefunctions as

$$\vec{\Sigma} \equiv \sum_{s_1\cdots s_5} \chi^{*Y^*}_{s_1 s_2 s_4}\, \chi^{*K^*}_{s_3 s_5}\, \chi^{p}_{s_1 s_2 s_3}\, \vec{\chi}_{s_4 s_5} \,, \tag{5}$$

with

$$\vec{\chi}_{s_4 s_5} \equiv \begin{pmatrix} 2\delta_{s_4\uparrow}\delta_{s_5\uparrow} \\ -\delta_{s_4\downarrow}\delta_{s_5\uparrow} - \delta_{s_4\uparrow}\delta_{s_5\downarrow} \\ -2\delta_{s_4\downarrow}\delta_{s_5\downarrow} \end{pmatrix} \,, \tag{6}$$

and $\vec{I}$ a spatial overlap:

$$\vec{I} = 2\gamma_0\left(\frac{3}{4\pi b}\right)^{3/2} \int d^3k\, d^3p\, d^3s \; \exp\left(\frac{-s^2}{2b}\right)$$

$$\times\Phi^*_{Y^*}\left[\mathbf{k},\, -\sqrt{\frac{3}{2}}\left(\mathbf{p}-\frac{\mathbf{s}}{2}-\frac{m_s}{m_{uus}}\mathbf{q}\right)\right]\; \Phi^*_{K^*}\left[\mathbf{p}+\frac{\mathbf{s}}{2}-\frac{m_s}{m_{us}}q\right]$$

$$\times\mathbf{p}\exp\left(-\frac{2}{3}r_q^2p^2\right)\; \Phi_p\left[\mathbf{k},\, -\sqrt{\frac{3}{2}}\left(\mathbf{p}+\frac{\mathbf{s}}{6}-\mathbf{q}\right)\right]. \tag{7}$$

Here the Φ's are momentum space wavefunctions, $\mathbf{q}$ is the momentum of Y^*, and the m_i's are quark masses (m_{uus} is short for $2m_u + m_s$, *etc.*). The factor $\exp(-s^2/2b)$ is the overlap of the initial and final-state flux tube wavefunctions; its size is controlled by the physical string tension b.

For the remaining quark line diagrams in Fig. 4, the decay amplitude still has the form (4), but the spin indices in Eq. (5) become permuted. The spatial overlap in (7) remains the same thanks to the assumed symmetry of the proton's spatial wavefunction.

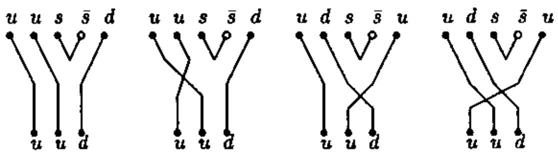

Figure 4: Quark line diagrams for $p \to \Sigma^* K^*$ and $p \to \Lambda^* K^*$.

Faced with the large number of states that contribute to the sum in Eq. (2), it was necessary to use simple harmonic oscillator (SHO) wavefunctions for the baryons and mesons in (7). The oscillator parameters β (defined by $\Phi(\mathbf{k}) \sim e^{-k^2/2\beta^2}$), were taken to be $\beta_{meson} =$ 0.4 GeV for mesons [14] and $\beta_{baryon} = 0.32$ GeV for baryons [15]. As discussed below, the results are quite insensitive to changes in the β's (mainly because Eq. (2) is independent of the choice of wavefunctions in the closure limit - - - any complete set gives the same result - - - and the full calculation with energy denominators does not deviate much from this limit.)

Even with SHO wavefunctions, the sum over intermediate states would be very difficult were it not for an important selection rule: inspection of the quark line diagrams in Fig. 4 shows that the relative coordinate of the non-strange quarks in baryon Y^* is always in its ground state. Only the relative coordinate between the strange and non-strange quarks (*i.e.*, the λ_{Y^*}-oscillator) can become excited. This drastically reduces the number of states that must be summed over. Unfortunately, this simplification does not apply for $u\bar{u}$ or $d\bar{d}$ pair creation.

It is useful to refer to the closure-spectator limit of Eq. (2). This is the limit in which the energy denominators do not depend strongly on the quantum numbers of Y^* and K^*, so that the sums over intermediate states collapse to 1, giving

$$\langle O_s \rangle \propto \langle p \, | h_{s\bar{s}} O_s h_{s\bar{s}} | \, p \rangle \propto \langle 0 \, | h_{s\bar{s}} O_s h_{s\bar{s}} | \, 0 \rangle \ , \tag{8}$$

where the second step follows since $h_{s\bar{s}}$ does not couple to the motion of the valence spectator quarks. We see that the expectation value of O_s is taken between the $^3P_0$ states created by $h_{s\bar{s}}$. From the J^{PC} of the $^3P_0$ pair it then follows that $\Delta s = R_s^2 = \mu_s = 0$ in the closure-spectator limit (a result which would not be seen if only the lowest term, or lowest few terms, were included in the closure sum).

In the next Section, I will discuss the results for the expectation values defined by Eq. (2) for the quantities Δs, R_s^2, and μ_s. We will see that delicate cancellations lead to small values for these observables even though the probability of $s\bar{s}$ pairs in the proton is substantial.

3.1 Strange Spin Content

Δs, the fraction of the proton's spin carried by strange quarks, is given by twice the expectation value of the s and $\bar{s}$ spins :

$$\Delta s = 2 \left\langle S_z^{(s)} + S_z^{(\bar{s})} \right\rangle . \tag{9}$$

Let us first examine the contribution to Δs from just the lowest-lying intermediate state, ΛK. The P-wave ΛK state with $J = J_z = \frac{1}{2}$ is

$$\left|(\Lambda K)_{P\frac{1}{2}}\right\rangle = \sqrt{\frac{2}{3}} \left|(\Lambda_\downarrow K)_{m=1}\right\rangle - \sqrt{\frac{1}{3}} \left|(\Lambda_\uparrow K)_{m=0}\right\rangle . \tag{10}$$

The $\bar{s}$ quark in the kaon is unpolarized, while the s quark in the Λ carries all of the Λ's spin; because of the larger coefficient multiplying the first term in (10), the ΛK intermediate state alone gives a negative contribution to Δs.

When we add in the $(\Lambda K^*)_{P\frac{1}{2}}$ and $(\Lambda K^*)_{P\frac{3}{2}}$ states (note that the subscripts denote the quantities ℓS defined previously), we have

$$\Delta s \propto \begin{pmatrix} 1 & -\sqrt{\frac{1}{3}} & \sqrt{\frac{8}{3}} \end{pmatrix} \frac{1}{18} \begin{bmatrix} -3 & \sqrt{3} & -\sqrt{24} \\ & -1 & \sqrt{8} \\ & & 10 \end{bmatrix} \begin{pmatrix} 1 \\ -\sqrt{\frac{1}{3}} \\ \sqrt{\frac{8}{3}} \end{pmatrix} \tag{11}$$

in the closure limit. Here the matrix is just $2(S_z^{(s)} + S_z^{(\bar{s})})$ (which is of course symmetric), and the vectors give the relative coupling strengths of the proton to $\{ (\Lambda K)_{P\frac{1}{2}}, (\Lambda K^*)_{P\frac{1}{2}}, (\Lambda K^*)_{P\frac{3}{2}} \}$. There are a couple of things to note here:

(1) The matrix multiplication in (11) evaluates to zero; there is no net contribution to Δs from the ΛK and ΛK^* states in the closure limit. There are in fact many such "sub-cancellations" in the closure sum for Δs: for each fixed set of spatial quantum numbers in the intermediate state, the sum over quark spins alone gives zero (because $\langle S_z^{(s)} \rangle = \langle S_z^{(\bar{s})} \rangle = 0$ in the $^3P_0$ state). That is, each $SU(6)$ multiplet inserted into Eq. (2) separately sums to zero. Moreover, the Δs operator does not cause transitions between $I = 0$ and $I = 1$ strange baryons so that the Λ and Σ sectors are decoupled, hence they individually sum to zero.

(2) Only the diagonal term in Eq. (11) corresponding to $p \to (\Lambda K^*)_{P\frac{3}{2}} \stackrel{\Delta s}{\to} (\Lambda K^*)_{P\frac{3}{2}} \to p$ gives a positive contribution to Δs. (Here $\stackrel{\Delta s}{\to}$ denotes the action of the Δs operator.) All of the other terms give negative contributions. In the full calculation with energy denominators, the negative terms are enhanced because they contain kaon (rather than K^*) masses. The full calculation gives $\Delta s = -0.065$ from ΛK and ΛK^* states. The largest individual contribution is -0.086, from the off-diagonal term $p \to (\Lambda K)_{P\frac{1}{2}} \stackrel{\Delta s}{\to} (\Lambda K^*)_{P\frac{3}{2}} \to p$.

For intermediate states containing Σ and Σ^* baryons, one finds

$$2(S_z^{(s)} + S_z^{(\bar{s})}) = \frac{1}{54}\begin{bmatrix} 3 & -12\sqrt{2} & 3\sqrt{3} & -6\sqrt{6} & 0 & 0 \\ & 15 & 0 & 0 & 6\sqrt{3} & -3\sqrt{15} \\ & & -7 & -10\sqrt{2} & -4\sqrt{2} & -4\sqrt{10} \\ & & & 10 & 4 & 4\sqrt{5} \\ & & & & -2 & -2\sqrt{5} \\ & & & & & 17 \end{bmatrix} \tag{12}$$

in the basis $\{ (\Sigma K)_{P\frac{1}{2}}, (\Sigma^* K)_{P\frac{3}{2}}, (\Sigma K^*)_{P\frac{1}{2}}, (\Sigma K^*)_{P\frac{3}{2}}, (\Sigma^* K^*)_{P\frac{1}{2}}, (\Sigma^* K^*)_{P\frac{3}{2}} \}$. The corresponding relative couplings to the proton are $[-\frac{1}{3}, -\sqrt{\frac{8}{9}}, \sqrt{\frac{25}{27}}, \sqrt{\frac{8}{27}}, \sqrt{\frac{8}{27}}, \sqrt{\frac{40}{27}}]$.

Again, the net Δs from these states is zero in the closure limit, but this time the insertion of energy denominators does not spoil the cancellation very much: the full calculation gives $\Delta s = -0.003$ in this sector.

P-wave hyperons and kaons contribute another -0.04 to Δs, and the net contribution from all higher states is -0.025. *Thus, the result of the calculation [12] is* $\Delta s = -0.13$, in quite good agreement with the most recent extractions from experiment [17,18], which give [19] $\Delta s = -0.10 \pm 0.03$. It should be emphasized that all parameters of this calculation were fixed by spectra and decay data. Moreover, the result is quite stable to parameter changes, varying by at most ± 0.025 when r_q, b, $\beta_{\rm baryon}$ and $\beta_{\rm meson}$ are individually varied by 30%.

For comparison with other calculations, note that the ΛK intermediate state alone contributes -0.030 to Δs, and the contribution from the ΛK, ΣK, and $\Sigma^* K$ states together is (coincidentally) also -0.030.

It is interesting to observe that Δs is driven mainly by meson, rather than baryon mass splittings: if one sets $m_\Lambda = m_\Sigma = m_{\Sigma^*}$, then Δs decreases by only about 30%, whereas it drops by about 80% if one sets $m_K = m_{K^*}$. Finally, an analogous calculation gives for the charm-quark contribution to the proton spin $\Delta c \approx -0.01$.

3.2 Strangeness Radius

Figure 5 defines our spatial variables for the quarks in an intermediate state. The (squared) distances of the s and $\bar{s}$ quarks from the baryon-meson center of mass are

$$r_s^2 = (\mathbf{r}_4 - \mathbf{R}_{\rm cm})^2 = \left[-\sqrt{6}\left(\frac{m_u}{m_{uus}}\right)\lambda_{Y^*} + \epsilon_{K^*}\mathbf{r} \right]^2 \tag{13}$$

$$r_{\bar{s}}^2 = (\mathbf{r}_5 - \mathbf{R}_{\rm cm})^2 = \left[-\left(\frac{m_u}{m_{us}}\right)\mathbf{r}_{K^*} - \epsilon_{Y^*}\mathbf{r} \right]^2 , \tag{14}$$

where $\epsilon_{K^*} \equiv M_{K^*}/M_{Y^*K^*}$ and $\epsilon_{Y^*} \equiv M_{Y^*}/M_{Y^*K^*}$, while by definition

$$R_s^2 \equiv r_s^2 - r_{\bar{s}}^2 \tag{15}$$

is the strangeness radius.

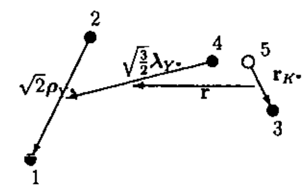

Figure 5: Quark coordinates in an intermediate state.

The calculation of R_s^2 is more difficult than the calculation of Δs, for several reasons. First, the operators appearing in R_s^2 cause orbital and radial transitions among the intermediate states. Thus SHO transitions satisfying $\Delta n = 0, \pm 1$ and/or $\Delta \ell = 0, \pm 1$ are allowed, so there are many more terms to calculate (n and ℓ are orbital and radial SHO quantum numbers). Moreover, the sub-cancellations discussed above no longer occur, so that R_s^2 converges more slowly than Δs: more states must be included in Eq.(2) to obtain good accuracy. In addition, the basic matrix elements are more complicated: in a basis of states with good magnetic quantum numbers (m, m'), for example,

$$\begin{aligned} \langle n'\ell'm'|r_{K^*z}|n\ell m\rangle = \\ \delta_{\ell'\,\ell-1}\delta_{m'\,m}\beta_{K^*}A_{\ell,m}\left(\sqrt{n+\ell+1/2}\,\delta_{n'\,n} - \sqrt{n+1}\,\delta_{n'\,n+1}\right) \\ +\delta_{\ell'\,\ell+1}\delta_{m'\,m}\beta_{K^*}A_{\ell+1,m}\left(\sqrt{n+\ell+3/2}\,\delta_{n'\,n} - \sqrt{n}\,\delta_{n'\,n-1}\right) \end{aligned} \tag{16}$$

Table 1: Proton strangeness radius from hadronic loops (in fm^2). The rows give the running totals as progressively more excited intermediate states are added into the calculation. The final column thus shows the total from all intermediate states.

	S-waves	plus P-waves	plus D-waves and S-wave radial excitations	all states
r_s^2	.097	.198	.210	.173
$r_{\bar{s}}^2$	.094	.139	.185	.210
R_s^2	.003	.059	.025	$-.04$

for matrix elements of the meson internal coordinate and

$$\langle q'\ell'm'|r_z|q\ell m\rangle = i\delta_{m'\,m}\left\{\delta_{\ell'\,\ell-1}A_{\ell,m}\left[-\frac{d}{dq}+\frac{\ell-1}{q}\right]\right.$$
$$\left.-\ \delta_{\ell'\,\ell+1}A_{\ell+1,m}\left[\frac{d}{dq}+\frac{\ell+2}{q}\right]\right\}\frac{\delta(q-q')}{q^2} \tag{17}$$

for matrix elements of the $Y^* - K^*$ relative coordinate, where $A_{\ell m} = \sqrt{\frac{(\ell+m)(\ell-m)}{(2\ell+1)(2\ell-1)}}$. These matrix elements must be coupled together to give $\langle R_s^2\rangle$ between states of definite ℓ and S with total angular momentum $\frac{1}{2}$, leading to formulas which become quite lengthy, especially for excited intermediate states. There is fortunately a stringent check of the results: when one equates all of the energy denominators in Eq. (2), the closure-spectator result, $R_s^2 = 0$, must be obtained.

The results for R_s^2 are shown in Table 1. With the standard parameter set, $R_s^2 = -0.04\text{fm}^2$. For reasonable parameter variations, R_s^2 ranges between -0.02 and -0.06fm^2. Table 1 shows that the lowest-lying $SU(6)$ multiplets of intermediate states (*i.e.*, the S-wave hyperons and kaons) account for about half of r_s^2 and $r_{\bar{s}}^2$. Most of the remaining contributions come from P-wave hyperons and kaons. However, R_s^2 involves a large cancellation between r_s^2 and $r_{\bar{s}}^2$, and its value doesn't settle down until we add in quite highly excited intermediate states. For this reason, the precise numerical value (and perhaps even the sign) of R_s^2 cannot be considered definitive: the conclusion is rather that R_s^2 is small, about an order of magnitude smaller than r_s^2 and $r_{\bar{s}}^2$. This result is not too surprising: R_s^2 is exactly zero in the closure limit, and previous hadronic loop studies [7,10] led one to expect that the full calculation with energy denominators would preserve the qualitative features of this limit.

Note that the ΛK intermediate state alone gives $R_s^2 \approx -0.01\text{fm}^2$ (the sign is as expected from the usual folklore) while the ΛK, ΣK, and $\Sigma^* K$ states together give -0.017fm^2. Nevertheless, although the sum over all states gives the same sign and order of magnitude as these truncations, Table 1 shows that this is just a coincidence.

3.3 Strange Magnetic Moment

The strange and antistrange quarks carry magnetic moments $-\frac{1}{3}\mu^{(s,\bar{s})}$ where

$$\mu^{(s)} = \frac{1}{2m_s}\langle 2S_z^{(s)} + L_z^{(s)}\rangle \tag{18}$$

$$\mu^{(\bar{s})} = -\ \frac{1}{2m_s}\langle 2S_z^{(\bar{s})} + L_z^{(\bar{s})}\rangle \tag{19}$$

and we denote the net strange magnetic moment by μ_s:

$$\mu_s \equiv \mu^{(s)} + \mu^{(\bar{s})}. \tag{20}$$

The spin expectation values are already in hand from the Δs calculation. Referring again to Fig. 5, we see that the s and $\bar{s}$ orbital angular momenta are given by

$$L_s = (\mathbf{r}_4 - \mathbf{R}_{cm}) \times \mathbf{p}_4$$
$$= \left[-\sqrt{6}\left(\frac{m_u}{m_{uus}}\right)\lambda_{Y^*} + \epsilon_{K^*}\mathbf{r}\right] \times \left[-\frac{2}{\sqrt{6}}\mathbf{p}_{\lambda_{Y^*}} + \left(\frac{m_s}{m_{uus}}\right)\mathbf{q}\right] \tag{21}$$

$$L_{\bar{s}} = (\mathbf{r}_5 - \mathbf{R}_{cm}) \times \mathbf{p}_5$$
$$= \left[-\left(\frac{m_u}{m_{us}}\right)\mathbf{r}_{K^*} - \epsilon_{Y^*}\mathbf{r}\right] \times \left[-\mathbf{p}_{K^*} - \left(\frac{m_s}{m_{us}}\right)\mathbf{q}\right]\ . \tag{22}$$

Computing the expectation values of these operators presents no new difficulties beyond those encountered in the R_s^2 calculation. In fact, there are no radial transitions in this case, so there are fewer states to sum over and the sum converges more quickly.

The results obtained with the standard parameter set are

$$\begin{aligned} \langle 2S_z^{(s)}\rangle &= -0.058 & \langle 2S_z^{(\bar{s})}\rangle &= -0.074 \\ \langle L_z^{(s)}\rangle &= 0.043 & \langle L_z^{(\bar{s})}\rangle &= 0.038 \\ \mu^{(s)} &= -0.025\mu_N & \mu^{(\bar{s})} &= 0.060\mu_N \\ & & \mu_s = 0.035\mu_N & \end{aligned} \tag{23}$$

The result is a positive (albeit small) value for μ_s, in disagreement with most other models. Where does the positive sign originate? First note that the signs of $\langle S_z^{(s)}\rangle$, $\langle L_z^{(s)}\rangle$, and $\langle L_z^{(\bar{s})}\rangle$ are correctly given by just the lowest lying intermediate state, ΛK of Eq. (10). (Note that the L_z's have similar magnitudes so that orbital angular momentum contributes very little to μ_s in any case.) On the other hand, the ΛK state has $\langle S_z^{(\bar{s})}\rangle = 0$, whereas $\langle S_z^{(\bar{s})}\rangle$ is quite large and negative. (The main contribution comes from the off-diagonal process $p \to (\Lambda K)_{P\frac{1}{2}} \xrightarrow{S_z^{(\bar{s})}} (\Lambda K^*)_{P\frac{1}{2}} \to p$, although there is also a significant contribution from $p \to (\Lambda(1405)K)_{S\frac{1}{2}} \xrightarrow{S_z^{(\bar{s})}} (\Lambda(1405)K^*)_{S\frac{1}{2}} \to p$.) These important terms, which drive μ_s positive, are omitted in calculations which include only kaon loops. (The ΛK intermediate state alone contributes $-0.080\mu_N$ to μ_s, and the contribution from ΛK, ΣK, and $\Sigma^* K$ together is $-0.074\mu_N$.)

3.4 Comments on the Effects of $s\bar{s}$ Loops

The calculations just described here represent parameter-free calculations of the effects of the $s\bar{s}$ sea generated by strong Y^*K^* loops on the low energy, nonperturbative structure of the nucleons. They are to my knowledge the first such results within a framework which has been demonstrated to be consistent with the many empirical constraints which should be applied to such calculations, namely consistency with the success of the quark potential model's spectroscopy and especially with the validity of the OZI rule.

The results indicate that observable effects from the strange sea generated by such loops arise from delicate cancellations between large contributions involving a suprisingly massive tower of virtual meson-baryon intermediate states. If correct, these conclusions rule out the utility of a search for a simple but predictive low energy hadronic truncation of QCD. While complete (in the sense of summing over all OZI-allowed Y^*K^* loops) and gauge invariant, the calculation has ignored pure OZI-forbidden effects as well as certain loop diagrams directly generated by the probing current (independently gauge invariant contact terms). As a consequence, they cannot strictly speaking be taken as predictions for Δs, R_s^2, or μ_s. Rather, the calculation shows that a complete set of strong Y^*K^* loops, computed in a model consistent with the OZI rule, gives very small observable $s\bar{s}$ effects. While such OZI-allowed processes *might* dominate, one cannot rule out the possibility (as was also the case with $\omega - \phi$ and other meson mixing [10]) that direct OZI violation (and in this case contact graphs as well) could make additional contributions of a comparable magnitude.

The small residual effect calculated for the loop contributions to Δs seems consistent with the most recent analyses of polarized deep inelastic scattering data. The calculations also give small residual strange quark contributions to the charge and magnetization distributions inside the nucleons. If these contributions are dominant, it will be a challenge to devise experiments that are capable of seeing them. Indeed, they are sufficiently small that their observation would appear to require the development of special apparatus dedicated to this task. Given the fundamental nature of the puzzling absence of other signals for the strong $q\bar{q}$ sea in low energy phenomena, this effort seems very worthwhile.

It would be desirable to devise tests of the mechanisms underlying the delicate cancellations which conspire to hide the effects of the sea in this picture. It also seems very worthwhile to extend such a calculation to $u\bar{u}$ and $d\bar{d}$ loops. Such an extension could reveal the origin of the observed violations [20] of the Gottfried Sum Rule [21] and also complete our understanding of the origin of the spin crisis. Since the effects of the $s\bar{s}$ loops are generally at high mass, it seems likely that the Pauli principle will have only a minor effect on $u\bar{u}$ and $d\bar{d}$ loops. Thus, from this calculation, I consider it likely that these lighter quarks will carry an even larger negative polarization than strange quarks, making it plausible that the "missing spin" of the proton is in orbital angular momentum.

4 Summary

In these lectures I have advocated treating the phenomenology of QCD in two steps. In the zeroth order, Strong QCD is approximated by a relativistic constituent quark model with flux tube gluodynamics. As a second step, $q\bar{q}$ sea and other $1/N_c$ effects are added as perturbations.

We have seen here how the quark model might be "unquenched" in a way that preserves its spectroscopic successes and respects the OZI rule. All of the results presented are qualitative, but the model appears to be a viable candidate to explain the underlying physics.

If the picture I have advocated is correct, there are some immediate consequences:

1. Low energy hadronizations of QCD are in trouble, since sums over large towers of states were required to preserve the spectrum and the OZI rule,

2. $\Delta s \simeq -0.13$, suggesting $\Delta u \sim \Delta d \sim -0.2$, implies that when combined with $\Sigma_{valence} \sim 0.75$, the missing spin of the proton would reside in orbital angular momentum, and

3. μ_s and $r^2{}_s$ are small.

This discussion has also made it clear that "singlet physics" has revealed a serious flaw in the constituent quark model which had been hidden by minus signs:

$$\Delta u - \Delta d = (\Delta u_v + \Delta u_{sea}) - (\Delta d_v + \Delta d_{sea}) \simeq \Delta u_v - \Delta d_v, \tag{24}$$

so that non-singlet quantities like the Bjorken Sum Rule [22] are misleading about the success of the quenched quark model. However, it now seems plausible that the unquenched quark model can successfully describe those properties strongly affected by the sea, so I will close by declaring:

THE NAIVE QUARK MODEL IS DEAD
LONG LIVE THE QUARK MODEL.

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