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Abstract

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by

Michael S. Chanowitz

ABSTRACT

The properties of QCD which imply the existence of gluonic states are reviewed. The problem of discovering the spectrum of gluonic states is discussed in general and illustrated with examples from current data. Higher statistics fixed target experiments, such as could be performed at LAMPF II, are essential for further progress.

I. INTRODUCTION

I am happy for the chance to give this talk because I believe that identification of the spectrum of gluonic states will require a new generation of higher statistics fixed target experiments such as could be performed at LAMPF II. My talk is in two parts. In the first I want to show you the beautiful equations that define quantum chromodynamics (QCD) and why they lead us to expect the existence of glueballs, hadrons which to first approximation are composed entirely of gluons. For the remainder I will discuss the problem of experimentally verifying this prediction.

The experimental reality is more complex and obscure than the beautiful equations. There are two reasons for this. First, theorists have not yet been able to make reliable quantitative predictions of glueball masses and decays, so the experimenters do not know exactly what to look for. Second, the relevant mass region, say from ~ $1\frac{1}{2}$ to ~ $2\frac{1}{2}$ GeV, is filled with an enormously complex spectrum of "ordinary" $\bar{q}q$ mesons,[1,2] not to mention the extraordinary variants which may also exist (e.g., cavity or string excitations[1]). To identify the glueballs we will have to understand the ordinary $\bar{q}q$ spectrum far better than we do now.

The plan of the talk is as follows. After the beautiful equations of Section II I will present in Section III the very conservative criteria which I believe can reliably be used to identify at least some of the glueball states. I will illustrate the discussion with a brief review of the two leading glueball candidates, stressing the need for additional fixed target data. In Section IV I will discuss the mixed states of quarks and glue which we in Berkeley call meiktons[3] (pronounced "make-ton", from the classical Greek for a mixed thing) and which our English cousins at Rutherford Lab call hermaphrodites.[4] (See also ref. 5 for other approaches to $\bar{q}qg$ states.) After a brief review of the conceptual issues I will discuss a meikton candidate state which, again, requires higher statistics fixed target studies to understand. Finally in Section V I will try to give a scale for the kind of statistical increase that is needed. As is fully appreciated by the participants in this workshop, this level of statistics requires not only higher flux beams but also new developments in detectors and in both on-line and off-line analysis.

II. LOCAL SYMMETRY: GLUEBALLS AS A FUNDAMENTAL CONSEQUENCE OF QCD

QCD is an example of what we call, following Pauli, a "gauge" theory. The simplest gauge theory is QED, quantum electrodynamics. The hallmark of any gauge theory is a local symmetry (called gauge invariance in the jargon). In QED this symmetry is just multiplication by an imaginary phase. The symmetry is local in the sense that we require invariance while allowing the phase to be an arbitrary (though smooth) function of space and time. QED is an "Abelian" gauge theory because multiplication by phases is commutative, $e^{i\alpha}e^{i\beta} = e^{i\beta}e^{i\alpha}$.

Local symmetry is a very strong demand to make of a theory. It means that we can change the phases of the fields in different parts of this room without changing the observable physics. It is a much more stringent and remarkable requirement than "global" symmetry, invariance under multiplication by a phase that is the same for all space and time. It is not surprising that locally symmetric theories have very special properties.

In QED the ingredients are one "matter" field (the electron) with charge -1

$$\psi(x) = \psi(\bar{x},t) \qquad Q = -1 \quad , \tag{1}$$

and the gauge field (the photon) which is electrically neutral

$$A_\mu(x) \qquad \mu = 0,\ 1,\ 2,\ 3 \qquad Q = 0 \quad . \tag{2}$$

Under a gauge transformation we multiply the matter field by an imaginary phase

$$\psi(x) \rightarrow e^{iQ\Lambda(x)}\ \psi(x) \tag{3}$$

where $\Lambda(x)$ is an arbitrary function of $x = \bar{x},t$. Since Q appears in the exponent, electrically neutral matter fields are not transformed. For small Λ we can expand to first order,

$$\psi(x) \rightarrow \psi(x) + iQ\Lambda(x)\psi(x) \quad . \tag{4}$$

Now the Lagrangian which defines QED contains a term, related to the electron kinetic energy, which is

$$\bar{\psi}\gamma_\mu \frac{\partial}{\partial x_\mu} \psi . \qquad (5)$$

For the physics to be locally symmetric the entire Lagrangian must be invariant under (3) [or to first order under (4)]. If Λ were just a constant (5) would be invariant but because of the derivative $\partial/\partial x_\mu$ ($x_0 = t$, $x_{(1,2,3)} = \bar{x}$) (5) is not invariant when $\Lambda = \Lambda(x)$. Pauli realized that local symmetry would be restored by replacing the ordinary derivative in (5) with his "gauge covariant" derivative

$$\mathcal{D}^\mu = \frac{\partial}{\partial x_\mu} - ieQA^\mu \quad , \qquad (6)$$

<u>and</u> requiring the photon field to transform under the gauge transformation like

$$A^\mu(x) \rightarrow A^\mu(x) + \frac{1}{e}\frac{\partial}{\partial x_\mu}\Lambda(x) \quad . \qquad (7)$$

Now instead of just (5) the Lagrangian contains the term

$$\bar{\psi}\gamma_\mu \mathcal{D}^\mu \psi \qquad (8)$$

which <u>is</u> invariant: the unwanted term that appeared when we transformed (5) is just cancelled by the transformation of A^μ in (7). For the rest of the Lagrangian we follow Maxwell, defining the field strength tensor

$$F^{\mu\nu} = \frac{\partial}{\partial x_\mu} A^\nu - \frac{\partial}{\partial x_\nu} A^\mu \qquad (9)$$

which is invariant under (7)

$$F^{\mu\nu} \rightarrow F^{\mu\nu} . \qquad (10)$$

The complete locally invariant Lagrangian of QED is then

$$\mathcal{L} = \bar{\psi}(i\gamma_\mu \mathcal{D}^\mu - m)\psi - \frac{1}{4} F_{\mu\nu}F^{\mu\nu} \qquad (11)$$

where the electron-photon interaction is hidden in the covariant derivative.

QCD is just like QED but with one crucial difference: the local symmetry of QCD is that of a non-Abelian group. That is, the transformations which are the counterparts of (3) do not commute with one another. The unidimensional charge Q of QED is replaced by a multi-dimensional, non-commuting collection of charges, Q_α. The symmetry of QED is U(1), the symmetry of the unit circle. The symmetry of QCD is the group SU(3) with $3^2-1 = 8$ charges or generators Q_α, $\alpha = 1, \ldots, 8$, analogous to the $2^2-1 = 3$ generators T_i of the SU(2) of isospin. In the typically tasteless jargon of particle physics we use the term "color" for the degrees of freedom of the QCD SU(3), analogous to the isospin degrees of freedom of SU(2) or to the unidimensional charge of the QED U(1). But unlike the SU(2) of isospin (in truly tasteless fashion particle physicists refer to isospin as a "flavor" symmetry) which is a global symmetry, the color SU(3) of QCD is a local symmetry. In nuclear physics we have approximate symmetry under global isospin rotations, which might, for instance, interchange all protons and neutrons. The gauge theory analogue would be much stronger: it would require exact symmetry under locally space-time dependent isospin rotations, which might, for instance, interchange protons and neutrons in one corner of the room while doing some quite different isospin rotation in another corner.

The central point is this: in order to implement Pauli's trick for the non-Abelian case there must be a gauge boson corresponding to each charge operator. In QED we have one photon corresponding to the single charge operator Q of U(1). In QCD we have eight gluons for the eight color charges Q_α of SU(3). Like the charges Q_α, the gluons also transform under the group, therefore, unlike the photon which is electrically neutral, they are not color neutral. Therein lies the tale! -- asymptotic freedom, confinement, and the existence of glueballs.

To exhibit the similarities and differences I will write down the QCD counterparts of the QED equations (1) through (11).

The matter field (the quark) is in the $\underline{3}$ representation of SU(3).

$$\psi_a(x) \in \underline{3} \qquad a = 1, 2, 3 \tag{1'}$$

("a" is the "color" index) while the gauge fields (gluons) are in the $\underline{8}$

$$A^\mu_\alpha \in \underline{8} \qquad \alpha = 1, 2, \ldots, 8 \ . \tag{2'}$$

(I will always use Latin letters for the $\underline{3}$, a,b = 1, 2, 3 and Greek for the $\underline{8}$, $\alpha,\beta,\gamma = 1, \ldots, 8$.) Then under a local SU(3) rotation specified by $\Lambda_\alpha(x)$ the quarks transform as

$$\psi_a(x) \to e^{i(Q_\alpha)_{ab}\Lambda_\alpha(x)} \psi_b(x) \tag{3'}$$

where repeated indices are summed and $(Q_\alpha)_{ab}$ is the 3×3 matrix representation of the Q_α's in the $\underline{3}$ representation. Thus the quark fields rotate in color space with axes and amounts specified by $\Lambda_\alpha(x)$. For small Λ_α we have to first order

$$\psi_a(x) \rightarrow \psi_a(x) + i(Q_\alpha)_{ab}\Lambda_\alpha(x)\psi_b(x) \quad . \tag{4'}$$

The statements that $\psi_a \varepsilon \underline{3}$ and $A_\alpha \varepsilon \underline{8}$ are analogous to the charge assignments of the electron ($Q = -1$) and the photon ($Q = 0$), since they tell us how the fields change under gauge transformation. As in QED the term in the Lagrangian

$$\bar{\psi}_a \gamma^\mu \frac{\partial}{\partial x^\mu} \psi_a \tag{5'}$$

is not invariant under (3') because of the action of the derivative on $\Lambda_\alpha(x)$.

Again we define a covariant derivative

$$\mathcal{D}^\mu_{ab} = \frac{\partial}{\partial x_\mu} \delta_{ab} - ig(Q_\alpha)_{ab}A^\mu_a \tag{6'}$$

where δ_{ab} is the Kronecker delta and g, the analogue of e in (6), is the strong interaction coupling constant. But compared to (7) the gauge transformation of the gluon field has an extra term (which I show to first order in small Λ_α)

$$A^\mu_\alpha \rightarrow A^\mu_\alpha + \frac{1}{g}\frac{\partial}{\partial x_\mu} \Lambda_\alpha + i(Q_\gamma)_{\alpha\beta}\Lambda_\gamma A^\mu_\beta \quad . \tag{7'}$$

Now we have successfully duplicated Pauli's trick since we can replace (5') by

$$\bar{\psi}_a \gamma_\mu \mathcal{D}^\mu_{ab} \psi_b \tag{8'}$$

which is gauge invariant. As before the second term in the transformation of the gluon field, (7'), cancels the noninvariance of (5').

The new feature, the third term in (7') arises because the gluon carries color so that it too rotates in color space just as the quark does in (4'). Because of this extra term in (7') the gluon field strength tensor contains an extra term not found in (9) which is bilinear in the gluon field,

$$F^{\mu\nu}_\alpha = \frac{\partial}{\partial x_\mu} A^\nu_\alpha - \frac{\partial}{\partial x_\nu} A^\mu_\alpha + ig(Q_\alpha)_{\beta\gamma}A^\mu_\beta A^\nu_\gamma \tag{9'}$$

and which is required so that $F^{\mu\nu}_\alpha$ rotates correctly (covariantly) under the transformation (7'):

$$F^{\mu\nu}_{\alpha} \to F^{\mu\nu}_{\alpha} + i(Q_{\gamma})_{\alpha\beta}\Lambda_{\gamma}F^{\mu\nu}_{\beta} \quad . \tag{10'}$$

Finally we can write the full locally symmetric Lagrangian of QCD (for one quark flavor)

$$\mathcal{L} = \bar{\psi}_a(i\gamma_\mu \mathcal{D}^\mu_{ab} - m\,\delta_{ab})\,\psi_b - \frac{1}{4}F^{\mu\nu}_{\alpha}F_{\alpha\mu\nu} \quad . \tag{11'}$$

Apart from the color subscripts which decorate (11') the QCD Lagrangian looks just like the QED Lagrangian (11). The crucial difference is in the last term of (11'), where it is hidden by the compact notation. Because $F^{\mu\nu}_{\alpha}$ contains terms linear and bilinear in the gluon field, the F^2 term in (11') contains three and four point gluon interaction vertices which have no counterpart in QED.

Since the basic simple ideas may have gotten lost in the unfamiliar mathematical expressions, I will summarize the argument in words. Local symmetry is implemented by Pauli's minimal substitution trick which requires the gauge bosons (photon/gluons) to interact with all quanta that carry the appropriate (electric/color) charge -- see (8) and (8'). Photons are electrically neutral and are not self-coupled but local non-Abelian symmetry requires that gluons carry nonvanishing color charge. Therefore gluons interact with themselves, as shown by the three and four point interaction vertices in (11').

These gluon-gluon interactions are the cause of the remarkable dynamical properties of QCD which distinguish it from QED. The first of these is asymptotic freedom, the "anti-screening" of the QCD vacuum which makes color charges appear weaker at short distances. The flip-side of asymptotic freedom is confinement, which is to say that quarks and gluons are confined to net-color-neutral hadrons by potentials which rise linearly with increasing separation. (Confinement is a proven property of space-time lattice models and is widely believed to be true in the continuum limit.) Finally the third remarkable property of QCD is the existence of purely gluonic states, glueballs, which have no counterparts in QED.

The expected existence of glueballs follows from color confinement and the fact that gluons carry color charge. According to confinement, only color neutral states, that is, singlets of the color SU(3), are directly observable in the laboratory. Thus a meson made of a quark-antiquark pair is the color singlet combination of $\bar{q}q$ pairs,

$$|\,\text{meson}\,\rangle = \frac{1}{\sqrt{3}}\sum_{a=1}^{3}|\bar{q}_a q_a\rangle \quad . \tag{12}$$

Similarly two gluons cannot separate by an arbitrarily large distance because of the confining potential but they can form a color singlet combination

$$|\,\text{glueball}\,\rangle = \frac{1}{\sqrt{8}}\sum_{\alpha=1}^{8}|g_\alpha g_\alpha\rangle \quad . \tag{13}$$

Equation (13) suggests that glueballs are made of "valence" gluons as mesons and baryons are known to be made of "valence" quarks. This is in fact a controversial point: valence glue is inescapable in the bag model but is not evident in the coarse-grained limit of the lattice calculations. I will say more of this in the discussion of meiktons in Section IV.

In looking for the glueball spectrum we are going to the heart of the remarkable properties of QCD.

III. GLUEBALLS

While the mathematical formulations of QED and QCD in Section II are equally beautiful and elegant, experimental verification of QCD is by no means as stunningly clear as for QED. To date the principal evidence for QCD is based on asymptotically-free perturbation theory and is plagued by ambiguities due to the still sizeable value of $\alpha_s = g^2/4\pi \cong 1/3 - 1/10$ at the distance scales probed. Such ambiguities do not afflict QED because of the comfortably small value of $\alpha = e^2/4\pi = 1/137$. For large distances, $\ell \gtrsim 10^{-13}$ cm. = 1 fm., theorists are even less able to extract the consequences of the equations of Section II. In particular there are currently no reliable quantitative predictions for the masses, much less the widths and decay modes, of the expected glueballs. Lattice calculations may ultimately succeed but for now the artifacts and uncertainties due to the small sizes of the lattices prevent them from being quantitatively reliable.

In this context it seems to me that there are only two properties which we can reliably use to identify glueball states. I call them the M.O. or modus operandi of the glueball. The first is a tautology: glueballs do not fit in the $\bar{q}q$ multiplets of the quark model. The second requires some quantum mechanics: glueballs are copiously produced in hard gluon channels. This second part of the M.O. is most evident if glueballs are made of valence gluons.

A good example of a hard gluon channel is radiative decay of $J/\psi(3095)$, $J/\psi \to \gamma + X$. In perturbation theory this decay is dominated by $J/\psi \to \gamma + gg$ where the two gluons are in a net color singlet. This is therefore a beautiful channel to look for glueballs with positive C-parity.

In the literature other properties have been proposed as a basis for identifying glueballs. One is that glueballs, being flavor singlets, should have flavor symmetric decays.[6] Another is that glueball widths should be the geometric mean of OIZ allowed and forbidden decays.[7] I do not believe either proposition is reliable. The first overlooks large symmetry breaking effects of both dynamical and kinematical origin. The second actually assumes much more than the OIZ rule. My criticisms of these propositions are presented elsewhere so I will not discuss them any further here.[8]

The good news is that the conservative M.O. is very powerful. It applies to glueballs even when they are mixed with $\bar{q}q$ states. The bad news is that it is not easy to apply, because knowing that a particle is not a $\bar{q}q$ meson requires a thorough understanding of the $\bar{q}q$ spectrum. But this is good news for proponents of high intensity facilities like LAMPF II which are essential to master the $\bar{q}q$ spectrum in the very complex region above ~1½ GeV.

I will illustrate these points with a brief discussion of two blue-ribbon glueball candidates that have been found in radiative J/ψ decay. The first, $\iota(1440)$, is seen at a large rate,

$$B(J/\psi \to \gamma\iota) \cdot B(\iota \to \bar{K}K\pi) \cong 4 \cdot 10^{-3}$$

large both as a fraction ($\gtrsim$ 5%) of all radiative decays and as larger than the previously most prominent state in the channel, the $\eta'(958)$. I argued that this state was not the $J^P = 1^+$ E(1420) but another state, with $J^P = 0^-$, discovered at CERN in $\bar{p}p$ annihilation at rest in 1965.[9] This appears to be confirmed by the subsequent spin-parity analysis of the Crystal Ball group which yielded $J^P = 0^-$ for the iota.[10] The analysis of reference (9) illustrates how complicated this mass region can be and how crucial good quality, high statistics data is to explicate the spectrum.

The second blue-ribbon candidate is $\theta(1700)$. It was seen first in radiative J/ψ decay in $\theta \to \eta\eta$,[11] subsequently in $\bar{K}K$,[12] and perhaps also in $\rho\rho$.[13] J^P is not conclusively measured though there is a reported[11] preference for 2^+. If we add the rates for these three modes we find

$$B(\psi \to \gamma\theta) \gtrsim 5 \cdot 10^{-3}$$

even larger than $\iota(1440)$ and therefore the most prominent state to date in the channel. Another striking piece of evidence comes from comparing $\gamma\gamma \to \bar{K}K$[14] with $J/\psi \to \gamma\bar{K}K$,[12] from which we learn that

$$\frac{\sigma(\gamma\gamma \to \theta)}{\sigma(\gamma\gamma \to f')} << \frac{\Gamma(J/\psi \to \gamma\theta)}{\Gamma(J/\psi \to \gamma f')} .$$

This is just what we would expect from a pristine, unmixed glueball which would couple much more strongly to two gluons (in $J/\psi \to \gamma_{gg}$) than to two photons.

Because of their prominence in $J/\psi \to \gamma X$ both ι and θ merit consideration as possible glueballs. We want to consider whether they have natural assignments in the $\bar{q}q$ spectrum. In the case of $\iota(1440)$ I have considered this question in considerable detail.[8,9,1] Iota could be the missing ninth member of the π' nonet, the radial excitation of the pion nonet, for which there are now eight candidates: $\pi'(\sim1270)$, $K'(\sim1450)$ and an isoscalar $\zeta(1275)$. Iota would then be the isoscalar partner of the ζ. I have argued on various grounds that this is unlikely,[8,9] but what is needed is not more arguments but one good experiment to explore the mass region above 1500 MeV where the ninth member of the nonet (assuming it is not iota) is likely to appear. I call this ninth pseudoscalar ζ' and would expect it to be produced in $\pi^- p \to \bar{K}K\pi n$ and perhaps also in $\pi p \to \eta\pi\pi n$ or $\pi p \to \eta'\pi\pi n$. No experiment reported to date would have been able to observe the ζ'. In fact the $\zeta(1275)$ has only been seen so far by one experiment[15] because only this experiment studied the channel $\pi^- p \to \eta\pi\pi n$ with high enough statistics to perform a partial wave analysis and discover the 70 MeV $\zeta(1275)$ beneath the narrower $J^P = 1^+$ D(1280). Previous experiments which only looked at the $\eta\pi\pi$ mass histogram undoubtedly confused these two states, so the properties ascribed to D(1280) in the Particle Data booklet cannot be taken at face value without looking back critically at the experimental sources.

The question of the assignment of $\theta(1700)$ is less well formulated since we still lack a definite spin-parity assignment. The situation will be clearer if θ is a tensor rather than a scalar, since the $\bar{q}q$ scalars are particularly poorly understood.[1,2] If θ is a tensor then it seems rather too near f and f' to be their radial excitation and too strongly produced in $\psi \to \gamma X$. To strengthen the glueball interpretation we would, in analogy to ι and ζ', like to have experimental knowledge of the masses of the radially excited tensor nonet. For both ι and θ the glueball interpretation hangs on further studies of the meson spectrum which require a new generation of fixed target experiments.

IV. MEIKTONS

The gluonic degree of freedom might also be observed by finding the mixed $\bar{q}qg$ states[3,4,5] which I will call meiktons. I will briefly describe the bag model prediction for the ground state meikton nonets. If these states were observed it would confirm the existence of valence gluons in the particular form required by the bag model.

I mentioned in Section II that the idea of valence gluons is controversial. In fact we do not understand why there are even valence quarks! -- though the regularities of the meson and baryon spectra leave no doubt about the usefulness of the concept of valence quarks. The question is why mesons have many of the properties of $\bar{q}q$ states rather than say $\bar{q}\bar{q}qq$, $\bar{q}\bar{q}\bar{q}qqq$ ... as one might expect of very strongly interacting quark quanta. I want to suggest an answer based on two facts we have learned in recent years.[1] First, deep inelastic scattering experiments have taught us that asymptotic freedom extends out to larger distances than we had previously thought, to about one fermi rather than to a fraction of a fermi. Second, lattice studies show that the transition from strong to (asymptotically free) weak coupling occurs very abruptly as a function of distance and that the change occurs at about one fermi! Since hadron radii are about one fermi, this all suggests that perturbation theory may be a reasonable qualitative or even semi-quantitative guide to the physics of hadron interiors. Hence valence quarks and gluons may exist because of the surprising relevance of perturbation theory. In cavity perturbation theory, as done in the bag model, additional convergence is gained because the vertices are not point-like but are proportional to small overlap integrals of cavity eigenfunctions.

In the bag model the lowest energy quark mode has $J^P = \frac{1}{2}^+$ and energy $E = 2.04/R$ where R is the cavity radius. The lowest energy gluon mode is the transverse electric (TE) mode with, surprisingly, axial vector quantum numbers $J^P = 1^+$ and energy $E = 2.74/R$. The ground state meiktons are constructed from a $\bar{q}q$ pair, either the spin singlet with $J^{PC} = 0^{-+}$ or the triplet with $J^{PC} = 1^{--}$, combined with the TE gluon with $J^{PC} = 1^{+-}$. The result is four nonets having $J^{PC} = 1^{--}$, $(0,1,2)^{-+}$. Three groups[3,4] have now computed the masses of these nonets through $O(\alpha_s)$ in cavity perturbation theory and are in agreement except for differences in the treatment of quark and gluon self energies. The results from reference (3) are shown in Table 1 for three values of the ratio of gluon mode self energies $C_{TE}/C_{TM} = 2, 1, \frac{1}{2}$. This ratio is fixed at $\frac{1}{2}$ or 2 if the spin of $\theta(1700)$ is $J = 0$ or $J = 2$ respectively.

For the preliminary indication that θ is a tensor, the masses range from 1.2 to 2.1 GeV. The 1^{--} nonet complicates the already complicated situation expected in the nonrelativistic quark model which may have <u>two</u> $\bar{q}q$ nonets in this region: the radial excitation, $L = 0$, $N = 2$, and the d-wave orbital excitation, $L = 2$, $N = 1$. The 0^{-+} nonet falls in the range of the radially excited π' $\bar{q}q$ nonet with $L = 0$, $N = 2$. The 2^{-+} nonet is near the region of the d-wave spin singlet $\bar{q}q$ nonet, L = 2, N = 1. But the 1^{-+} nonet is a quark model exotic; that is, $J^{PC} = 1^{-+}$ does not appear in the spectrum of the nonrelativistic $\bar{q}q$ model (although 1^{-+} states do appear as cavity excitations of $\bar{q}q$ states in the bag model[1,3]). It is therefore particularly interesting to look for the states of the 1^{-+} nonet.

These $\bar{q}qg$ states are likely to decay by formation of a $\bar{q}q$ pair from the gluon, $g \to \bar{q}q$, followed by disassociation of the resultant $\bar{q}\bar{q}qq$ state into two $\bar{q}q$ mesons. Because of parity the TE gluon does not couple to an s-wave pair $\bar{q}_s q_s$ (we use j - j coupling in the bag) but to $\bar{q}_p q_s$ or $\bar{q}_s q_p$. The result then is either two L = 0 mesons in a relative p-wave or an L = 0 and an L = 1 meson in a relative s-wave,

$$\bar{q}_s q_s g(TE) \rightarrow \begin{cases} (\bar{q}q)_s + (\bar{q}q)_s & L = 1 \\ (\bar{q}q)_s + (\bar{q}q)_p & L = 0 \end{cases}$$

Examples of these two kinds of decays for the isovector member of the exotic 1^{-+} nonet are

$$"\rho"(1^{-+}) \rightarrow \begin{cases} \pi\eta & L = 1 \\ \pi D(1280) & L = 0 \end{cases}$$

The $\pi\eta$ channel is particularly nice because it is an experimentally clean two body channel and because $\pi\eta$ in a p-wave uniquely signals the 1^{-+} quantum numbers.

Here even more than for the glueballs we depend upon the results of high quality, fixed target experiments. For instance, too many states of a given quantum number could indicate the existence of meikton nonets. There is an intriguing example of this already in the experimental literature. The ACCMOR collaboration at CERN accumulated 600,000 events in the reaction $\pi^- p \rightarrow \pi^+\pi^-\pi^- p$ from which many interesting results were obtained[16] -- on the previously controversial A_1 meson, on the radially excited π' and K', and on a 2^{-+} isovector the $A_3(1700)$.

They confirmed the existence of $A_3(1700)$, primarily in the $f\pi$ s-wave, though they saw it less clearly also in the $f\pi$ d-wave, $\rho\pi$ p-wave, and $\varepsilon\pi$ d-wave. And they saw a second bump in the 2^{-+} channel at 1850 MeV, which I will call the A_3', only 150 MeV above the A_3! This second bump appeared in $(f\pi)_d$, $(\rho\pi)_p$, and $(\varepsilon\pi)_d$ but not in the $f\pi$ s-wave. From Table I with $C_{TE}/C_{TM} = \frac{1}{2}$ as for $J(\theta) = 2$, the $"\rho"(2^{-+})$ meikton is expected in just this region, at 1790 MeV. Now if the 2^{-+} meikton and the 2^{-+} d-wave $\bar{q}q$ isovectors had nearly equal masses between 1750 and 1800 MeV they would mix strongly. The mixing might naturally be dominated by the s-wave $f\pi$ intermediate state in which case the levels would "repel" and one of the eigenstates would tend to decouple from the $f\pi$ s-wave, leaving a picture like what is perhaps observed.

I say "perhaps" because the mass and even the existence of the A_3' are by no means clear. The experimenters have found a second interpretation of their data in which the bump at 1850 results from the interference of $A_3(1700)$ with a second state at ~ 2100 MeV. My private suspicion, which I have not yet been able to confirm with the principals, is that they were moved to find this second solution by the perception that nobody would love a second $I,J^{PC} = 1,2^{-+}$ state just 150 MeV above the A_3. Indeed such a state could not be explained in the $\bar{q}q$ model as an excitation of the A_3. Even the 400 MeV splitting corresponding to an $A_3'(2100)$ seems rather small for a radial excitation.

Another experiment with even more statistics is probably needed to decide the existence and mass of the A_3'. However the initial results from ACCMOR are a good case study in how careful study of the meson spectrum may turn up the new physics we are seeking.

V. THE FRONTIER IN STATISTICS

Given the present inability to calculate reliably the long distance properties of QCD, such as hadron masses and decays, it is clear that the discovery and identification of hadrons with gluon constituents requires a much more thorough experimental exploration of the meson spectrum than exists to date. In fact this will be so even if our theoretical understanding improves, because even if we did have reliable predictions of glueball and meikton masses and decays, positive identification of the new gluonic states would still in most cases require disentangling them from nearby $\bar{q}q$ mesons with the same quantum numbers. This in turn means that we still would need to know in detail the composition of the $\bar{q}q$ nonets.

The point is that high statistics is as much a frontier as high energy. In particle physics we go to higher energy to be able to resolve structure and dynamics at smaller distances. But in order to resolve the structure and dynamics of the meson spectrum, which is fundamental to QCD, we need not higher energy but higher statistics. In the past five years each increase in statistics has brought important new results. There is no reason to think that this progression has reached an end. Many of the most fundamental questions, such as the existence and nature of gluonic states, remain to be answered.

The particular statistical level required for the next step will vary from case to case. For example, the questions raised in Section III about iota and the π' nonet could probably be settled by experiments in $\pi^- p \to \bar{K}K\pi n$, $\eta\pi\pi n$ with good acceptance out to $\lesssim 1.7$ GeV and with statistical power comparable to or perhaps even less than that of the ACCMOR experiment mentioned in Section IV. But higher statistical levels will probably be needed to map the spectrum at higher masses, say from 1.5 to 2.5 GeV. For instance, to settle the question of the possible meikton candidate A_3' (1850/2100) raised by the ACCMOR collaboration, considerably more than their 600,000 event sample would be needed.

Here we move into the realm where a high intensity source such as LAMPF II is essential. In addition we begin to move beyond the state of the art in detectors and data analysis, matters being discussed in this workshop. To record the required data we will need faster detectors and/or new on-line triggers. To process an order of magnitude or more data than was accumulated by ACCMOR, progress will also be needed in off-line computing power. These are formidable technical problems but in the assumed ~ 10 year time frame we can be optimistic they can be solved. The result will be a quantum jump in our ability to confront and master some of the most fundamental problems of QCD.

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