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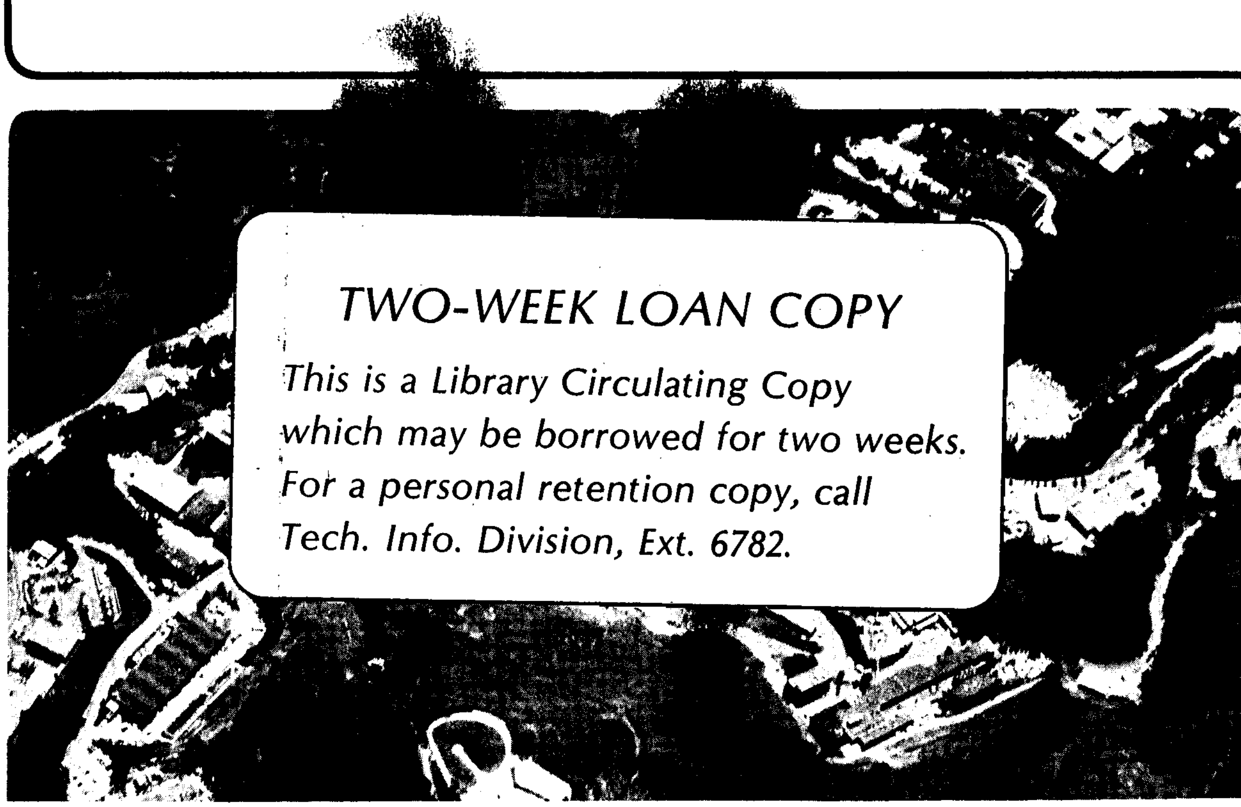

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J/ψ PHYSICS AT BEPC

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Talk presented at the First Workshop
on Colliding Beam Physics in China

Abstract

J/ψ physics is discussed which will be of interest at $T > 1988$, the period of operation of the Beijing Electron Positron Collider. Emphasis is placed on the gluonic states which are best studied in radiative J/ψ decay. The difficulties of these studies are discussed and the need for very high statistics is stressed. In particular it is essential to partial-wave-analyze the hadronic final states produced in $J/\psi \to \gamma X$. An estimate using fixed target data suggests that $0(10^8)$ J/ψ decays are needed to do an unambiguous partial wave analysis for hadron masses up to about 2 GeV. This requirement is an excellent match to the BEPC design parameters, which imply production of $0(10^8)$ J/ψ's per year. With a J/ψ production rate an order of magnitude greater than other electron-positron storage rings, BEPC will be a unique world facility for these studies.

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I. Introduction: the need for high statistics J/ψ studies

The discovery of J/ψ(3095) in 1974 was a momentous event for high energy physics. The dramatic two body $\bar{c}c$ spectrum — the ground states J/ψ and η_c, the radial excitations ψ′ and η'_c, and the orbitally excited χ states — convinced the last skeptics of the reality of quarks, despite the failure to observe them directly. Furthermore the successful description of the spectrum by potential models which are Coulombic at short distances and growing at long distances was an important verification of asymptotic freedom and quark confinement as expected in Quantum Chromodynamics.

Today the study of the decays of J/ψ remains at the center of the effort to understand Quantum Chromodynamics. In this talk I will try to show why this is still likely to be true toward the end of the decade when BEPC will begin operation. The reason very simply is that the questions which remain to study are of fundamental importance and they are exceedingly difficult. The answers will require very high statistics, perhaps many tens of millions of J/ψ decays. If BEPC approaches its design specifications for operation at the ψ, it will be the premier world machine for these studies. It will be a unique source of important physics results at the end of the decade and beyond.

The ψ decays of greatest current interest are the radiative decays, $\psi \to \gamma$ + hadrons, predicted to be about 8% of all ψ decays.[1] The experimental rate is consistent with this prediction.[2] The prediction is based on the lowest order Feynman diagram, $\psi \to \gamma + 2$ gluons, in which the two gluons interact to form the hadrons. Because the two gluons are in a color singlet, this is a beautiful channel to search for the glueball states which are expected as a unique consequence of the non-Abelian dynamics of QCD.[3] The masses of the lightest glueballs are thought to be between 1 and 2-1/2 GeV, which is well-matched to the masses that can be produced in radiative ψ decay. The enormous cross section for ψ production and its complete dominance over the continuum background make it the ideal channel in which to study the glueball spectrum.

Let me pause a moment to elaborate on the statement that glueballs "are expected as a unique consequence of the non-Abelian dynamics of QCD." QCD is a gauge theory like the more familiar QED or Quantum Electrodynamics. Both are based on symmetry, the difference being that the symmetry operations of QED commute with one another [the Abelian group U(1)] while those of QCD do not [the non-Abelian group SU(3)]. There is a gauge boson for each symmetry axis of the gauge group, hence QED has the one photon while QCD has $3^2 - 1 =$ 8 gluons. This is the crucial difference between QED and QCD and is responsible for the unusual properties of QCD: asymptotic freedom and confinement of color charge. Photons are electric charge neutral, and they do not interact to form "photonium" or "lightballs". But the eight gluons are color charged so they interact directly with each other and they must be confined. Therefore two gluons in a net color neutral configuration, a "color singlet", are expected to bind to form a glueball. The existence of glueballs is among the most striking properties expected of QCD. We will not be sure that we have really understood the strong interactions until we have verified this prediction. There are hints in existing data of new particles which might be glueballs, but the situation is unclear and controversial, for reasons I will discuss in detail below.

Although radiative ψ decay is the ideal place to look for glueballs, the search is very difficult, for two reasons. First, there is still no quantitatively reliable theoretical calculation of the glueball spectrum, so the experimenters don't know exactly where to look or what to look for. This may well have changed by the end of the decade, as progress is made in lattice techniques for calculating the spectrum. Second, the spectrum in the 1 to 2-1/2 GeV mass region is very complex. It contains over 26 $\bar{q}q$ nonets, of which only a few are completely known and many are not known at all. Inevitably many of these particles will overlap in mass with one another and with glueball states. They will also mix, to varying degrees, with glueballs of the same quantum numbers. The situation is further complicated by the possibility of other kinds of interesting new physics, such as $\bar{q}qg$ states (variously called meiktons, hermaphrodites, or hybrids) and "cryptoexotic" $\bar{q}\bar{q}qq$ states.

I will discuss below the problem of how we can hope to identify at least some of the glueball states despite these difficulties. Here I only wish to emphasize one point, which is the crucial importance of partial wave analysis. "Bump hunting" with mass histograms will not be sufficient. Only a few of the most obvious states are visible in mass histograms and even these cannot be understood without spin-parity determination. Many states will not even be visible above the background unless the data are partial wave analyzed.

We can get an idea of how much statistics is required from the experience of the Mark II, Crystal Ball, and Mark III detectors, and also by considering examples from meson partial wave analysis in fixed target experiments. The $\iota(1440)$ $\bar{K}K\pi$ resonance found with the Mark II[4] in $\psi \rightarrow \gamma\bar{K}K\pi$ could not be definitively distinguished from the $J^P = 1^+$ E(1420) until the Crystal Ball[5] observed 2,000,000 ψ decays and established $J^P = 0^-$. More recently we have learned from the Mark III study[6] of 2,500,000 ψ decays that the situation may be more complicated still: the structure of the $\bar{K}K\pi$ Dalitz plot as a function of the $\bar{K}K\pi$ mass and observation of a $\rho\gamma$ signal raise the possibility that the iota region might contain more than one state. To put this in perspective, it is important to realize that the iota is as prominent a state as we can hope to encounter in radiative ψ decay, since it appears in a relatively background free setting at a rate that is between 5 and 10% of all radiative ψ decays. Yet 2.5 million ψ decays have not been sufficient to give us a clear understanding of its structure.

Another instructive example is the very interesting $\xi(2220)$ discovered with the Mark III[6] in $\psi \rightarrow \gamma\bar{K}K$. Because of the low level of the $\bar{K}K$ background, the ξ could be seen in the mass histogram with a signal of only ~ 30 K^+K^- events. Its possible narrowness has led to speculation that it could be a Higgs boson, in which case it must have spin zero. Other hypotheses are that it is a $J = 0$ or $J = 2$ meikton or glueball or a high spin $\bar{q}q$ resonance. To measure its spin will require from 10,000,000 to 20,000,000 ψ decays.

A second interesting example from the Mark III[7] is a possible resonant structure in $\rho\rho$ seen in the $J^P = 0^-$ channel from 1.6 to 1.9 GeV. This is of great interest because in the nonrelativistic $\bar{q}q$ spectrum, excited 0^- isoscalar states can only be radial excitations. More pseudoscalars than required by the $\bar{q}q$ spectrum would be a sure sign of new physics. Much higher statistics are needed to verify and study this possible $\rho\rho$ resonance.

Fixed target experiments have identified almost all the particles in the four $L = 1$ qq nonets and at least some members of 22 additional excited nonets. With higher statistics, experiments have been able to extend partial wave analyses to ever larger masses. A notable example is the ACCMOR experiment[8] at CERN, which detected 600,000 events in the diffractive reaction $\pi^-p \rightarrow \pi^+\pi^-\pi^-p$. With these 600,000 $\pi^+\pi^-\pi^-$ with three pion mass between 1 and 3 GeV, they were able to resolve the puzzle of the $J^{PC} = 1^{++} A_1$ meson and to do an unambiguous analysis of the $J^{PC} = 2^{-+}$ channel to $M \approx 1700$ MeV, confirming the existence of the $A_3(1700)$, presumeably the $L = 2\ \bar{q}q$ isovector. However a definitive analysis was not possible for higher masses. In particular, a possible isovector structure at 1850 MeV was seen but not established, which might be related to the $L = 0, J^{PC} = 2^{-+}\ \bar{q}qg$ meikton. Even with 600,000 $\pi^+\pi^-\pi^-$ events, it was not possible to push the analysis unambiguously to 2 GeV.

From the results obtained already by the Mark III there is clearly strong motivation to increase the data sample to 10 to 20 million ψ's. The example of ACCMOR suggests that we will eventually need an even larger data sample. For instance, if we wanted to detect 100,000 events in a prominent channel (e.g., $\psi \to \gamma\eta\pi\pi$) which might correspond to ~ 10% of all radiative ψ decays, then even assuming a generous detection efficiency of 10% we would need ~ 100,000,000 ψ decays. We would then have ~ 1/6 as many events as ACCMOR accumulated in $\pi^+\pi^-\pi^-$ in a comparable mass region to radiative ψ decay, though this might be compensated by more favorable signal characteristics of the ψ data.[9] It is quite possible that an effort of this magnitude would be necessary to extract the full richness of the radiative ψ decay channel.

Depending on the success of the mini-beta system to be installed at SPEAR this summer, the Mark III might be able to detect 10-20 million more ψ's in the coming years. But there is no possibility in sight for producing 10^8 ψ's at SPEAR. If BEPC approaches its design specifications for luminosity and energy spread of the beam at $E_{beam} = M_\psi/2 = 1.55$ GeV, then it will surpass SPEAR performance (without mini-beta) by about an order of magnitude. If SPEAR gains the factor of ~ 3 in luminosity which the mini-beta system is hoped to provide, then we can also expect a comparable improvement from the addition of mini-beta to BEPC. In either case it is necessary to consider the problems of detecting and analyzing such an enormous number of events. If the detector can handle 10^8 ψ's per year, as I understand that it can, then the limiting factor would be the analysis, since the presently planned off-line computing system could not process more than ~ 5,000,000 ψ's per year (see presentation of T. Shalk). This problem would be partially alleviated if a hard photon trigger could be devised to select the radiative ψ decays, or if radiative ψ decays could be selected by an initial off-line screening procedure performed before the events were fully analyzed.

Since the BEPC luminosity is optimized at the maximum energy, $E_{beam} = 2.8$ GeV, it is crucial for ψ physics that the luminosity vary as close as possible to the $L \propto E^2$ limit for energies down to $E_{beam} = 1.5$ GeV. There are wigglers in the BEPC design for this purpose. If this goal is met the peak design luminosity of $1.7 \cdot 10^{31}$ cm.$^{-2}$ sec.$^{-1}$ at 2.8 GeV would imply a peak luminosity of $5 \cdot 10^{30}$ cm.$^{-2}$ sec.$^{-1}$ at $\psi(3095)$, about an order of magnitude above the peak luminosities achieved in ψ running at SPEAR. (This is partly because at SPEAR the luminosity is[10] approximately proportional to E^4 rather than E^2.) When SPEAR runs in top-off mode the average luminosity, taking account of injection time, beam loss, and other interruptions, has typically been $L_{average} \approx 1/2\ L_{peak}$, but out of top-off mode it is $L_{average} \approx 1/4\ L_{peak}$.[11] BEPC will inject at $E_{beam} = 1 - 1.4$ GeV, so it will not operate in top-off mode.

Since the ψ is a narrow resonance, $\Gamma_\psi = 63$ keV $<< \Delta E_{beam}$, the ψ yield is inversely proportional to ΔE_{beam}, the energy spread of the beam. The natural RMS value for the BEPC design is $\Delta E_{beam}/E_{beam} = 7.4 \cdot 10^{-4}$ at $E_{beam} = 2.8$ GeV. If the expected energy dependence, $\Delta E/E \propto E$, is realized, then at ψ we will have $\Delta E_{beam}/E_{beam} = 4.1 \cdot 10^{-4}$, nearly a factor 3 smaller than what has been achieved at SPEAR. Because of the design of the BEPC vacuum system, the actual $\Delta E/E$ may be close to the natural RMS limit.

Putting all this together and assuming $L_{average} \approx 1/4\, L_{peak}$, we find that BEPC could have a ψ event rate an order of magnitude greater than SPEAR. The figure of merit is $L_{average}/\Delta E_{beam}$ which is $\sim 2\frac{1}{2} \cdot 10^{29}$ cm.$^{-2}$ sec.$^{-1}$ MeV^{-1} for SPEAR compared to $\sim 2 \cdot 10^{30}$ cm.$^{-2}$ sec.$^{-1}$ MeV^{-1} for BEPC (with the above assumptions). To calculate the event rate, we begin with the theoretical cross section

$$\sigma_\psi = \frac{12\pi}{M_\psi^2} B(\psi \to e^+e^-) \cong 1.13 \cdot 10^{-28} \text{cm.}^2 \tag{1.1}$$

which must be smeared by the beam spread to find the observed cross section,

$$\bar{\sigma}_\psi = \sigma_\psi \cdot \frac{\Gamma_\psi(TOT)}{2\sqrt{2}\,\Delta E_{beam}} \cong 4 \cdot 10^{-30} \text{cm.}^2 \tag{1.2}$$

assuming ΔE_{beam} has the natural RMS value for the BEPC design. Finally for a run of $T = 10^7$ sec. = 1/3 year and assuming $L_{average} = 1/4\, L_{peak}$ we get the event yield

$$N_\psi = \bar{\sigma}_\psi\, T\, \mathcal{L}_{average} \cong 5 \cdot 10^7 \tag{1.3}$$

or about 4,000,000 radiative ψ decays. If this is achieved, we must also consider the computing facilities needed to carry out multi-amplitude partial wave analysis with data samples like 10^5 $\eta\pi\pi$.

The plan for the rest of the talk is as follows. Section II is a brief look at the beautiful equations that define non-Abelian gauge theories and QCD in particular. The purpose is to show why non-Abelian gauge invariance probably implies the existence of glueballs. This section is meant as a "cultural interlude" and can be omitted by the reader who wants to proceed directly to more "practical" matters.

The next sections are concerned with the new spectroscopic studies that could be done with twenty to a hundred million ψ decays and, in addition to some still open questions about the charmonium spectrum.

In Section III I discuss glueballs, beginning with a theoretical review of what we do and do not know about them (there is much more of the latter than the former). Despite the bad news about how much we do not know, there is still some good news: that the lightest glueballs are likely to be in the mass range that can be studied at BEPC and that, despite our theoretical ignorance, it should be possible to identify a few examples of glueballs with a lot of experimental hard work, much of which might only be possible at BEPC.

Sections IV and V concern two other examples of new spectroscopy, the $\bar{q}qg$ meikton/hermaphrodites and the $\bar{q}\bar{q}qq$ cryptoexotics. Again I will give brief reviews and then discuss the phenomenological problems of finding and identifying these new objects. As with the glueballs, the discussion will be illustrated with examples from existing data.

In Section VI I will present my understanding of the present status of charmonium studies and the value of pursuing additional measurements. Although the nonrelativistic potential model is a good approximate description of the $\bar{c}c$ spectrum, there are large relativistic and rescattering corrections which are not under precise theoretical control. Therefore it seems appropriate to concentrate on the remaining qualitative open questions, such as the mass of the $^1P_1$ state and the structure of the reasonance region between 4 and 4-1/2 GeV.

Section VII is a brief review of the work that has been done on the possibility that $\xi(2220)$ is a Higgs boson.

Section VIII is a summary and statement of conclusions.

II. Local Symmetry: glueballs as a fundamental consequence of QCD

This section can be omitted by those interested in the strictly phenomenological issues. It is meant to provide a quick look at the mathematical structure of a non-Abelian gauge theory and to show in particular why gauge invariance in QCD implies that gluons can bind to one another to form glueballs. My purpose is to emphasize that the effort to discover the glueball states is directly related to the most elegant aspects of the mathematical structure 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

$$\psi(x) = \psi(\vec{x},t) \qquad Q = -1 \tag{2.1}$$

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

$$A_\mu(x) \qquad \mu = 0, 1, 2, 3 \qquad Q = 0 \tag{2.2}$$

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

$$\psi(x) \longrightarrow e^{iQ\Lambda(x)}\,\psi(x) \tag{2.3}$$

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

$$\Psi(x) \to \Psi(x) + i Q \Lambda(x) \Psi(x) \tag{2.4}$$

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

$$\bar{\Psi}(x) \gamma_\mu \frac{\partial}{\partial x_\mu} \Psi(x) \tag{2.5}$$

For the physics to be locally symmetric the entire Lagrangian must be invariant under (2.3) [or to first order under (2.4)]. If Λ were just a constant (2.5) would be invariant but because of the derivative $\partial/\partial x_\mu$ (2.5) is not invariant when $\Lambda = \Lambda(x)$. Pauli realized that local symmetry would be restored by replacing the ordinary derivative in (2.5) with his "gauge covariant" derivative

$$\mathcal{D}^\mu = \partial^\mu - ieQA^\mu \tag{2.6}$$

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

$$A^\mu(x) \longrightarrow A^\mu(x) + \frac{1}{e} \frac{\partial}{\partial x_\mu} \Lambda(x) \tag{2.7}$$

Now instead of just (2.5) the Lagrangian contains the term

$$\bar{\Psi}(x) \gamma_\mu \mathcal{D}^\mu \Psi(x) \tag{2.8}$$

which <u>is</u> invariant: the unwanted term that appeared when we transformed (2.5) is just cancelled by the transformation of A^μ in (2.7). For the rest of the Lagrangian we follow Maxwell, defining the field strength tensor

$$F^{\mu\nu} = \frac{\partial}{\partial x_\mu} A^\nu - \frac{\partial}{\partial x_\nu} A^\mu \tag{2.9}$$

which is invariant under (2.7)

$$F^{\mu\nu} \longrightarrow F^{\mu\nu} \tag{2.10}$$

The complete locally invariant Lagrangian of QED is then

$$\mathcal{L}(x) = \overline{\Psi}(i\gamma_\mu D^\mu - m)\Psi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} \tag{2.11}$$

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 (2.3) do not commute with one another. The unidimensional charge 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, $\alpha = 1, ..., 8$, analogous to the $2^2 - 1 = 3$ generators T_i of the SU(2) of isospin. In 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 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 <u>approximate</u> symmetry under <u>global</u> isospin rotations, which might, for instance, interchange all protons and neutrons. The gauge theory analogue would be much stronger: it would require <u>exact</u> symmetry under <u>locally</u> 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). <u>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.</u> 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 (2.1) through (2.11).

The matter field (the quark) is in the <u>3</u> representation of SU(3).

$$\psi_a(x) \in \underline{3} \qquad a = 1, 2, 3 \tag{2.1'}$$

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

$$A^\mu_\alpha(x) \in \underline{8} \qquad \alpha = 1, 2, \ldots, 8 \tag{2.2'}$$

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

$$\psi_a(x) \longrightarrow e^{i(Q_\alpha)_{ab}\Lambda_\alpha(x)}\,\psi_b(x) \tag{2.3'}$$

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

$$\psi_a(x) \longrightarrow \psi_a(x) + i\,(Q_\alpha)_{ab}\,\Lambda_\alpha(x)\,\psi_b(x) \tag{2.4'}$$

The statements that $\psi_a \in \underline{3}$ and $A_\alpha \in \underline{8}$ 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

$$\bar{\psi}_a\,\gamma_\mu \frac{\partial}{\partial x_\mu}\,\psi_a \tag{2.5'}$$

is not invariant under (2.3′) because of the action of the derivative on $\Lambda_\alpha(x)$.

Again we define a covariant derivative

$$\mathcal{D}^{\mu}_{ab} = \partial^{\mu}\delta_{ab} - ig\,(Q_\alpha)_{ab}\,A^{\mu}_{\alpha} \qquad (2.6')$$

where δ_{ab} is the Kronecker delta and g, the analogue of e in (2.6), is the strong interaction coupling constant. But compared to (2.7) the gauge transformation of the gluon field has an extra term (which I show to first order in small Λ_α)

$$A^{\mu}_{\alpha} \rightarrow A^{\mu}_{\alpha} + \frac{1}{g}\frac{\partial}{\partial x_\mu}\Lambda_\alpha + i\,(Q_\gamma)_{\alpha\beta}\,\Lambda_\gamma\,A^{\mu}_{\beta} \qquad (2.7')$$

Now we have successfully duplicated Pauli's trick since we can replace (2.5′) by

$$\overline{\Psi}_a\,\gamma_\mu\,\mathcal{D}^{\mu}_{ab}\,\Psi_b \qquad (2.8')$$

which is gauge invariant. As before the second term in the transformation of the gluon field, (2.7′), cancels the noninvariance of (2.5′).

The new feature, the third term in (2.7′) arises because the gluon carries color so that it too rotates in color space just as the quark does in (2.4′). Because of this extra term in (2.7′) the gluon field strength tensor contains an extra term not found in (2.9) which is bilinear in the gluon field,

$$F^{\mu\nu}_{\alpha} = \frac{\partial A^{\nu}_{\alpha}}{\partial x_\mu} - \frac{\partial A^{\mu}_{\alpha}}{\partial x_\nu} + ig\,(Q_\alpha)_{\beta\gamma}\,A^{\mu}_{\beta}A^{\nu}_{\gamma} \qquad (2.9')$$

and which is required so that $F_\alpha{}^{\mu\nu}$ rotates correctly (covariantly) under the transformation (2.7′):

$$F^{\mu\nu}_{\alpha} \longrightarrow F^{\mu\nu}_{\alpha} + i\,(Q_\gamma)_{\alpha\beta}\,\Lambda_\gamma\,F^{\mu\nu}_{\beta} \qquad (2.10')$$

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

$$\mathcal{L}(x) = \overline{\Psi}_a \left(i\gamma_\mu \mathcal{D}^\mu_{ab} - m\,\delta_{ab} \right) \Psi_b - \frac{1}{4} F_\alpha^{\mu\nu} F_{\alpha\mu\nu} \tag{2.11'}$$

Apart from the color subscripts which decorate (2.11′) the QCD Lagrangian looks just like the QED Lagrangian (2.11). The crucial difference is in the last term of (2.11′), where it is hidden by the compact notation. Because $F_\alpha^{\mu\nu}$ contains terms linear <u>and</u> bilinear in the gluon field, the F^2 term in (2.11′) 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 (2.8) and (2.8′). 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 (2.11′).

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 opposite 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 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,

$$|\text{meson}\rangle = \frac{1}{\sqrt{3}} \sum_{a=1}^{3} |\bar{q}_a q_a\rangle \tag{2.12}$$

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

$$|\text{glueball}\rangle = \frac{1}{\sqrt{8}} \sum_{\alpha=1}^{8} |g_\alpha g_\alpha\rangle \tag{2.13}$$

Equation (2.13) suggest 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 V.

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

III. Glueballs

This section begins with a brief review of the present theoretical understanding of the glueball spectrum and dynamics. The conclusion of this review is that there are still no reliable, quantitative predictions of glueball masses or decays. This conclusion contributes significantly to the difficulty of the experimental search. However there is good news for BEPC in two conclusions about which there is general agreement: (1) the lightest glueballs lie between 1 and 2 GeV, and (2) radiative ψ decay is the best way to search for them. We may also be encouraged that lattice calculations have good prospects to yield more reliable predictions for the pectrum by the end of the decade when BEPC begins operation.

The second part of this section concerns how, despite the theoretical uncertainties, we can proceed experimentally to identify at least some of the glueball states. The emphasis is on the need to develop a complete picture of all the particles in a given partial wave, including especially the ordinary $\bar{q}q$ mesons and also the possible new $\bar{q}qg$ and $\bar{q}\bar{q}qq$ particles discussed in Sections IV and V. The glueball discussion is illustrated with the example of the I, $J^{PC} = 0,0^{-+}$ channel, which contains three established resonances — $\eta(549)$, $\eta'(958)$, $\iota(1440)$ — and two possible new resonances — the $\zeta(1270)$ seen in $\pi p \to \eta\pi\pi n$ and the $\rho\rho$ structure seen in $\psi \to \gamma\rho\rho$ between 1.6 and 1.9 GeV.

A. A Brief Theoretical Review

The glueball spectrum has been studied chiefly with three approaches: lattice calculations, potential models, and bag models.

Of these approaches, it is clear that the lattice method has the best chance to eventually give accurate quantitative results. It is equally clear that accurate quantitative results are not yet available today. The progress and difficulties are discussed in two recent reviews.[12] They describe three major sources of uncertainty in the present calculations:

(1) All hadron masses and, in particular the glueball masses, are proportional to the QCD scale parameter Λ. Lattice calculations typically determine Λ as a function of the string tension K, which is in turn estimated empirically from the slope of the Regge trajectories on the Chew-Frautschi plot or from the linear terms in charmonium and bottomonium potentials. A calculation reported last year on a larger lattice gave a value of Λ/K which was twice as big as earlier values. Until the discrepancy is resolved, we must acknowledge a factor 2 uncertainty in these lattice computations of glueball masses. (The scale may be set in other ways, e.g., by the chiral or gluon condensates or non-glueball hadron masses, but these methods have their own uncertainties.)

(2) The lattice spacing "a" must be much less than the typical hadronic scale, $r_H \sim 1/2 - 1$ fm. But the length of the lattice "universe", L = Na, must be much bigger than r_H to avoid spurious effects. So we require

$$1 \ll \frac{r_H}{a} \ll N \tag{3.1}$$

where it is sometimes even argued that N should be replaced by $N/2\pi$. Present computing methods are limited to N of order 10, so that (3.1) cannot be comfortably satisfied.

(3) Quantum corrections due to fermion loops are not included in the present calculations. For glueballs this means that the effect of mixing with $\bar{q}q$ mesons are not incorporated in the results.

With the development of more powerful computing methods and machines it should be possible to make considerable progress on all these problems. Therefore theory may be able to offer much more solid guidance to BEPC glueball searches at the end of the decade than it can today. For now the principal conclusions which we can safely draw from lattice studies are that the lightest glueballs lie between ~ 1 and ~ 2 GeV, with quantum numbers $J^{PC} = 0^{++}, 0^{-+}, 2^{++}$. The lightest "odd ball" glueball (a state with quantum numbers not attainable in the nonrelativistic $\bar{q}q$ spectrum) is typically found to be the $J^{PC} = 1^{-+}$ state, a little heavier than the 2^{++}.

The bag model and the nonrelativistic potential model are both useful though limited phenomenological guides to the light $\bar{q}q$ spectrum. Applied to the glueball spectrum they are likely to be even less reliable. In the potential models [13] neither the value of the assumed gluon "constituent mass" nor the strength of the confining potential are known from $\bar{q}q$ physics. There is also disagreement about whether the constituent gluon has a longitudinal spin mode. The bag model[14] has the advantage of being a relativistic approximation which treats the gluon spin unambiguously and requires no new parameters to fix the glueball spectrum to leading order in α_s, the strong coupling constant. Both approaches suffer from the likelihood of large quantum corrections due to the large spin and color Casimir values of the gluon. The convergence of the perturbation expansion has not been established in either approach.

Calculations done to $O(\alpha_s)$ in the bag model[15,16] agree with the lattice results that the lightest glueballs have $J^{PC} = 0^{++}, 0^{-+}, 2^{++}$ and are likely to be found between ~ 1 and ~ 2 GeV. Similar results are found in the potential models.

I will briefly describe the bag model calculation.[14] Free, massless (and therefore unambiguously transverse) gluons are confined to a static spherical cavity. Just as for the analogous problem in electromagnetism, the single gluon eigenmodes are determined by the boundary conditions to be TE (transverse electric), $P = (-1)^{L+1}$, or TM (transverse magnetic), $P = (-1)^L$. The energies of the three lowest models in terms of the cavity radius R are

$$\begin{array}{lll} TE_1 & L^P = 1^+ & E = 2.74/R \\ TE_2 & L^P = 2^- & E = 3.96/R \\ TM_1 & L^P = 1^- & E = 4.49/R \end{array} \qquad (3.2)$$

The ground state glueball is then constructed from two TE_1 modes. Since it is a color singlet Bose statistics requires the symmetric combination, therefore $J^{PC} = 0^{++}, 2^{++}$. The energy E(R) is then

$$E(R) = 2 \cdot \frac{2.74}{R} + \frac{4}{3}\pi R^3 B \tag{3.3}$$

where B is the bag constant, determined from meson and baryon spectrum. Minimizing with respect to R we find

$$E = \left(\frac{4\pi B}{3}\right)^{1/4}\left(\frac{8 \cdot 2.74}{3}\right)^{3/4} \cong 1 \text{ GeV}. \tag{3.4}$$

The first excited states are obtained from $TE_1 \times TE_2$ with $J^{PC} = (1,2,3)^{-+}$ and $TE_1 \times TM_1$ with $J^{PC} = (0,1,2)^{-+}$. However four of these six states are regarded as p-wave excitations of the $(0,2)^{++}$ ground state with respect to the artificially fixed cavity. This gives four "spurious" states with quantum numbers $(0,2)^{++} \times 1^{-} = 1(1,1,2,3)^{-+}$. Discarding these, the remaining two excited states with $J^{PC} = (0,2)^{-+}$ are found at ~ 1.3 GeV.

Three groups have calculated the $O(\alpha_s)$ corrections to these results.[15,16] The degeneracies are split, with the 0^{++} falling below the 2^{++} and the 0^{-+} falling below the 2^{-+}. The calculations are incomplete because the self energies of the gluon cavity modes are not known, though they are calculable in principle. Furthermore, none of the bag model calculations with gluon constituents satisfy the correct boundary conditions which really require nonspherical cavities) though there are indications this may not have a large effect on the masses. The convergence of cavity perturbation theory applied to glueballs cannot be known until the gluon self energies are computed, but the part of the $O(\alpha_s)$ corrections already computed is uncomfortably large. It is likely that the bag model will be no more than a qualitative guide to the glueball spectrum.

This is a discouraging description of the present state of theoretical knowledge of the glueball spectrum. Dynamical properties are even harder to understand, such as decay widths, branching ratios, and mixing with $q\bar{q}$ states.

For example, consider the question of the decay widths of glueball states. A statement sometimes made in the literature is that typical glueballs should have a width which is the geometric mean of OIZ allowed and OIZ suppressed meson decay widths.[17] This estimate of glueball widths is based on the observation that in perturbative QCD, OIZ violating amplitudes are mediated by intermediate gluons. The initial state quarks annihilate to gluons which then create the final state quarks. For glueball decay only the latter occurs so we expect a suppression which is the square-root of full OIZ suppression. This estimate ignores the distinction between the two and three gluon channels which is phenomenologically important: the large deviation from ideal mixing of the light pseudoscalars shows that the OIZ rule is not honored in the J = 0 two gluon channel at ~ 1 GeV.

There is also a more general difficulty. The estimate follows from the tacit assumption that the intermediate gluons in the Feynman diagram of an OIZ violating amplitude implies glueball dominance of the real intermediate states. For instance, in a glueball pole model $\phi \rightarrow G \rightarrow \rho\pi$ meson-glueball couplings appear twice, yielding the estimate $\Gamma_G \sim (\Gamma_{OIZ\ allowed} \cdot \Gamma_{OIZ\ forbidden})^{\frac{1}{2}}$. But the identification of intermediate gluons with intermediate glueballs overlooks the existence of what may be the dominant intermediate states. For instance $\phi \to \rho\pi$ can proceed via the OIZ allowed $\bar{K}K$ intermediate state, $\phi \to \bar{K}K \to \rho\pi$, since $\phi \to \bar{K}K$ and $\bar{K}K \to \rho\pi$ are both OIZ allowed. This raises a "paradox", which for $\phi \to \rho\pi$ is most crisply formulated with the unitarity equation,

$$\mathrm{Im} \langle \phi | \rho\pi \rangle = \langle \phi | \bar{K}K \rangle \langle \bar{K}K | \rho\pi \rangle + \cdots \tag{3.5}$$

The left side is OIZ suppressed though neither factor on the right side is. Cancellations are not possible, since the other intermediate states are OIZ suppressed, so $\langle \bar{K}K | \rho\pi \rangle$ must be small though it is OIZ allowed. This shows that the OIZ rule is not self-contained, in the sense that some other dynamical principle is needed to make it consistent with unitarity. My view is that the small K:π ratios seen in the central region in hadron-hadron and e^+e^- collisions embody the physics of this unstated principle.

The real part of $\langle \phi | \rho\pi \rangle$ has contributions from intermediate glueballs and from OIZ allowed channels like $\bar{K}K$. If the real part is small and/or if it is saturated by the $\bar{K}K$ contribution, then the glueball couplings could be much smaller than the simple estimate. Or, if there were big cancellations, the glueball couplings could be much larger. Another uncertainty is the distance to the relevant glueball poles, which is probably large in this example but in general would vary greatly from case to case. (An important related point is that the qualitative expectation that $\psi \to \gamma X$ is a glueball channel is not affected by these considerations because the $\bar{D}D$ threshold is well away from the glueball masses being considered.)

My conclusion is that we do not know how broad glueballs are or even that there is a single scale which characterizes the width of the ordinary "garden-variety" glueball. To the contrary, as we already know for ordinary mesons, both kinematical and dynamical considerations may cause different glueballs to have widely varying widths.

Another often repeated statement is that glueballs are SU(3) flavor singlets and can therefore be identified by their flavor symmetric decays.[18] Again for both kinematical and dynamical reasons, this statement may be very unreliable. I will illustrate this with three examples.

First consider a spin zero glueball containg two "valence" gluons. The two lowest order decay diagrams are shown in figure (1). For a $J = 0$ initial state, Figure (1a) vanishes for massless quarks $m_q = 0$, while for massive quarks the amplitude is proportional to m_q. This is a consequence of the familiar argument based on helicity conservation which explains $\Gamma(\pi \to \mu\nu) >> \Gamma(\pi \to e\nu)$. It applies both to $J = 0$ glueballs and to the pion because in both cases the interactions are helicity conserving (V and V-A respectively). Therefore Figure (1a) favors $\bar{s}s$ over $\bar{u}u$ and $\bar{d}d$ by a factor which is at least ~ 3 for constituent quark masses and could be as large as ~ 400 for current quark masses. As discussed below, if $\iota(1440)$ is a glueball this argument could be at least a partial explanation of why it decays predominantly to $\bar{K}K\pi$, contrary to what would naively be expected for a flavor singlet state.[19,20]

Second, the bag model suggests a dynamical mechanism which causes certain glueballs (and also $\bar{q}qg$ meiktons) to decay predominantly to final states containing two or more strange particles.[21] In cavity perturbation theory the vertices are proportional to the overlap integrals of the cavity mode eigenfunctions. The lowest gluon mode, TE, has roughly flavor symmetric s-channel couplings to $\bar{u}u$, $\bar{d}d$, and $\bar{s}s$, but the TM mode couples much more strongly to $\bar{s}s$ (by ~ 5 in the amplitude). The TE mode has $J^{PC} = 1^{+-}$ while the TM gluon has $J^{PC} = 1^{--}$, so a $J^{PC} = 0^{-+}$ glueball is constructed from one TE and one TM mode. Therefore for a 0^{-+} glueball, such as $\iota(1440)$ might be, this contributes an additional enhancement of $\bar{s}s$ pairs in Figure (1a) and assures that one of the $\bar{q}q$ pairs in Figure (1b) is predominantly $\bar{s}s$. The same mechanism causes $\bar{d}sg_{TM}$ and $\bar{u}s\,g_{TM}$ meiktons to decay to final states with three strange particles, and $g_{TM}g_{TM}$ glueballs and $\bar{s}sg_{TM}$ meiktons to decay to final states with four strange particles (such as $\phi\phi \to \bar{K}K\bar{K}K$).

It is true these arguments rely heavily on perturbation theory, and, in the second instance, on details of the bag model, so they may not be entirely dependable. But at the very least they demonstrate how kinematics and dynamics could create large violations of flavor symmetry. They show in particular that we need not be surprised if we find a pseudoscalar glueball which decays predominantly to $\bar{K}K\pi$.

A third example of flavor symmetry breaking is suggested by the data[22] on another glueball candidate, $\theta(1700)$. The θ decays to $\bar{K}K$ much more than to $\pi\pi$,[23,6,7] contrary to flavor symmetry for a flavor singlet state. To show how this could occur, I will consider a simple model, not because I think it is a really adequate model of the physics, but just because it illustrates a point: that flavor symmetric QCD dynamics need not imply flavor symmetry of the exclusive final states. The model shows that $\pi\pi$ could be a smaller mode than $\bar{K}K$ even if θ is a $J = 2$ glueball and a flavor singlet.

The model is just Figure (1a). That is, I assume the decay begins with the flavor symmetric annihilation of the gluons to a single $\bar{q}q$ pair, given by $1/\sqrt{3}(\bar{u}u + \bar{d}d + \bar{s}s)$, which subsequently hadronizes to form the observed final states (if θ has $J = 0$ the $\bar{q}q$ pair would instead be mostly $\bar{s}s$ as discussed above for ι). Now we must consider how the $\bar{u}u + \bar{d}d$ and $\bar{s}s$ pairs hadronize. I will assume that no additional $\bar{s}s$ pairs are formed in the process of hadronization (a conservative assumption for the purpose at hand). Then the $\bar{u}u + \bar{d}d$ pairs will materialize as $(\pi\pi)_2$, $(\eta\eta)_2$, $(\eta\pi\pi)_3$, $(\pi\pi\pi\pi)_2 = (\rho\rho)_0$, and $(\pi\pi\pi\pi\pi\pi)_4 = (\omega\omega)_0$. The subscripts denote the least possible units of orbital angular momentum and I have indicated the dominant resonant combinations of the 4π and 6π states. Similarly the $\bar{s}s$ pairs materialize as $(\bar{K}K)_2$, $(\eta\eta)_2$, $(\bar{K}K\pi)_3 = (\bar{K}^*K)_2$, and $(\bar{K}K\pi\pi)_2 = (\bar{K}^*K^*)_0$.

The point is that for the $\bar{s}s$ decays the three and four body final states are heavily penalized by phase space as is the d-wave $\bar{K}^*K$ decay while the corresponding quasi-two-body channel $\bar{K}^*K^*$ is actually forbidden. But for the $\bar{u}u + \bar{d}d$ decays the four and six body decays proceed with no inhibition in the quasi-two-body s-wave channels $\rho\rho$ and $\omega\omega$, while $\eta\pi\pi$ has more available phase space than the corresponding $\bar{K}K\pi$. Therefore simply because of the available phase space we expect a much larger fraction of the $\bar{s}s$ decays to hadronize to $\bar{K}K$ than $\bar{u}u + \bar{d}d$ to $\pi\pi$. Flavor symmetry at the level of the quarks is not incompatible with the flavor symmetry breaking observed for the exclusive final states. (If this argument really applies to $\theta(1700)$ then some large nonstrange decay modes of θ should be found, such as $\rho\rho$, $\omega\omega$, or $\eta\pi\pi$; if not, another interpretation discussed in Section V is that θ is a $\bar{q}\bar{q}qq$ state or even two such states.)

The degree of mixing between glueball states and $\bar{q}q$ mesons (or with $\bar{q}qg$ and $\bar{q}\bar{q}qq$ states) is among the most difficult theoretical questions and may make the interpretation of the experimental data more complex in at least some cases. For mixing between two states the mixing angle is

$$\tan 2\theta = \frac{2\langle A | B \rangle}{m_A - m_B} \tag{3.6}$$

The numerator depends on the dynamics of wave functions we do not understand and the denominator is a matter of chance. My guess is that we will find at least a few examples where the mixing is not of order 1, that is, states which are pure glueball to a good approximation. This guess is based on the fact that the ordinary mesons and baryons are so well described in terms of their valence $\bar{q}q$ and qqq configurations. As discussed in Section IV, this may be due to a surprising weakness in the strength of the strong interactions that govern hadron structure. This in turn suggests that the numerator of eq. (3.6) may be small on the scale that controls hadron dynamics, so that θ may not be too big unless $m_A - m_B$ is rather small, say < 100 MeV.

The possibility of mixing is also closely related to the use of electromagnetic decays to try to distinguish glueballs from mesons. For instance, the decay of a pure glueball state to two photons would have to go by an intermediate quark loop, so we would expect it to be suppressed relative to the $\gamma\gamma$ decay width of a ground state $\bar{q}q$ meson. The actual amplitude would depend on a factor like the numerator of eq. (3.6) and for a mixed state would obviously also depend on the angle θ. In addition we must remember that the glueball candidates are likely to be in the 1-2 GeV region which contains excited $\bar{q}q$ states. Many of these (for instance, the radially excited pseudoscalars) will also tend to have suppressed $\gamma\gamma$ couplings relative to what we would estimate from the known $\gamma\gamma$ couplings of the ground states.

Some work has been done to model the mixing of glueballs with $\bar{q}q$ mesons.[24] This is interesting theoretical and phenomenological work, but the uncertainties in the assumptions and approximations mean that the results cannot be taken as reliable guides to the problems of interpreting the experimental data. With a clearer understanding of the experimental picture we will be better able to test the validity of such theoretical models. On the other hand, the experimental data cannot be interpreted without some theoretical framework. Progress requires that we pull ourselves up by our bootstraps, with continual give and take between theory and experiment. In the next subsection I will review some of the existing data, to give an idea of how we might begin to identify some examples of the glueball spectrum while sing only the most simple and general features of the theoretical framework.

B. Looking for Glueballs

We want to use the most general and simple theoretical ideas. Two such important properties are

A) Glueballs do not fit in the $\bar{q}q$ nonets of the quark model.

B) Glueballs are copiously produced in hard gluon channels, the best example being radiative ψ decay.

A high statistics source of mesons is crucial to study these two properties, which is why BEPC could be the ideal instrument for glueball physics. Let's discuss A) and B) in turn.

Property A) is among the safest statements you will ever hear a theorist make. It applies even when glueballs are mixed with $\bar{q}q$ mesons: in any case, there will be too many particles to be classified in the $\bar{q}q$ spectrum. But to make use of property A) we must understand the qq spectrum very well, to be able to recognize the particles it does not contain. Therefore we will depend heavily on the progress being made in meson spectroscopy at fixed target accelerators. Even here BEPC will be crucial, since high statistics studies of $\psi \to \gamma X$ will help us understand precisely those $\bar{q}q$ mesons which have the same quantum numbers as glueballs and with which they can mix. In applying property A) we must of course be aware that if a particle is not an ordinary $\bar{q}q$ meson, it is not necessarily a glueball. Other possibilities could be the new particles discussed in Sections IV and V or new kinds of $\bar{q}q$ excitations (discussed below in connection with "oddballs") which are not found in the spectrum of the nonrelativistic $\bar{q}q$ model.

Property B) is not as safe as property A) since it follows not just from pure logic but requires some physics as well. It is especially likely to be correct if glueballs can be described in terms of valence gluons, as they are in bag and potential models. In lattice calculations there are no valence gluons in the strong coupling limit but there is evidence that they may emerge as the continuum limit is approached. Radiative ψ decay is the perfect example of a hard gluon source, since in perturbation theory the leading mechanism is $\psi \to \gamma gg$. The two gluons are in a color singlet so they could naturally resonate to form a glueball composed of two valence gluons. As the photon energy E_γ changes we can produce glueballs of mass

$$M = \sqrt{M_\psi^2 - 2E_\gamma M_\psi} \qquad (3.7)$$

In lowest order perturbation theory the dominant partial waves of the two gluons in $\psi \to \gamma gg$ are $J^{PC} = 0^{++}, 0^{-+}, 2^{++}$, which corresponds exactly to the quantum numbers of the lightest glueballs expected in the theoretical calculations discussed above!

It has been suggested[25] that a good way to identify a glueball is to find the $J^{PC} = 1^{-+}$ oddball-glueball, which has exotic quantum numbers not found in the nonrelativistic $\bar{q}q$ spectrum. It would certainly be interesting to study the 1^{-+} channel, though the following points should be kept in mind:

1) To the extent that lowest order perturbation theory is a good guide, we do not expect much $J^{PC} = 1^{-+}$ production in $\psi \to \gamma X$.
2) The 1^{-+} states are expected at bigger masses, where continuum backgrounds may be more severe.
3) $J^{PC} = 1^{-+}$ is not a unique signal for a glueball. For example, the four ground state $\bar{q}qg$ nonets, discussed in Section IV, include a 1^{-+} nonet. The isoscalars from these nonets could be confused with glueballs. Another possibility is the cavity fluctuations of $\bar{q}q$ mesons whose possible existence is suggested by the bag model.[26,27] These states are charge conjugation reflections of both orbital and radial excitations. For example the 0^{-+} and 1^{--} radially excited nonets

could have C-parity reflected states (at some higher mass) with exotic $J^{PC} = 0^{--}$ and 1^{-+}.

I will illustrate the possibilities and the difficulties of glueball searches by discussing as an example the iota particle, $\iota(1440)$, which may well turn out to be the first discovered glueball. Interest in this particle began precisely because it satisfied property B) above. That is, it was discovered in 1980 in $\psi \to \gamma \bar{K}K\pi$ at very large rate, now thought to be about $4 \cdot 10^{-3}$ of all ψ decays or about 5% of all radiative ψ decays.[4] It was at first confused with the E(1420), a predominantly $s\bar{s}$ meson in the $J^{PC} = 1^{++}$ nonet of the A_1 meson. Some theoretical considerations together with an analysis of the experimental history of E(1420) suggested that the particle seen at SPEAR was not the 1^{++} E but rather a 0^{-+} state.[19,20] To make matters even more confusing, this analysis implied that the resonance which was first called E, discovered at CERN in 1965 in $\bar{p}p$ annihilation (and named E, the first resonance discovered in Europe), is not the 1^{++} E(1420) but is instead the same particle seen 15 years later at SPEAR. This analysis was supported when the Crystal Ball group[5] established that the SPEAR $\bar{K}K\pi$ resonance is indeed a $J^{PC} = 0^{-+}$ state and suggested it be called iota 1440 or $\iota(1440)$.

The next question is whether the iota is a member of the radially excited 0^{-+} nonet. This nonet contains a broad $\pi'(1300)$ now seen by several groups, a K'(1400) seen by the LASS and ACCMOR groups, and an isoscalar which I call $\zeta(1270)$ that has been seen so far only by one high statistics experiment[28] in $\pi^- p \to \eta\pi^+\pi^- n$ (the only experiment with enough statistics to do the partial wave analysis needed to see it). The iota could then be the ninth member of this nonet, the second isoscalar. I do not think this is a likely explanation, although the question is not conclusively settled. I will explain briefly why I do not favor the interpretation of iota as a member of the π' nonet and will then discuss the evidence we would like to have to reach a more definite conclusion.

There are two striking facts which must be explained in any interpretation of $\iota(1440)$. These are (1) the very large branching ratio for $\psi \to \gamma\iota$ and (2) the dominance of the decay $\iota \to \bar{K}K\pi$. I am impressed that both facts are easily explained if iota is a $J^P = 0^-$ glueball. For the glueball interpretation the explanations have already been stated in the preceding subsection: (1) glueballs should be abundantly produced in $\psi \to \gamma X \sim \psi \to \gamma gg$ and (2) a pseudoscalar glueball will decay strongly to $\bar{K}K\pi$ because of the mass enhancement favoring $gg \to s\bar{s}$ in fig. (1c) and perhaps also for the enhancement of the bag model vertex $g_{TM} \to s\bar{s}$ which contributes in both figures (1a) and (1b).

According to the Mark III data[6] the branching ratio for $\psi \to \gamma\iota$ is

$$B(\psi \to \gamma\iota) \cdot B(\iota \to \bar{K}K\pi) = (5.6 \pm 0.4 \pm 1.3) \cdot 10^{-3} \qquad (3.8)$$

which gives a lower limit

$$B(\psi \to \gamma\iota) > (5.6 \pm 0.4 \pm 1.3) \cdot 10^{-3} \qquad (3.9)$$

Taking

$$B(\psi \to \gamma + \text{hadrons}) \sim 6-10\% \tag{3.10}$$

we see that $\psi \to \gamma\iota$ represents at least 5 to 10% of all radiative ψ decay, a very large fraction for such a complicated channel with many different possible final states. For comparison the second largest known resonance in the channel is the long established $\eta'(958)$, with

$$B(\psi \to \gamma\eta') = (3.6 \pm 0.5)\cdot 10^{-3} \tag{3.11}$$

This implies $B(\psi \to \gamma\iota) > 2\,B(\psi \to \gamma\eta')$, which for two reasons is contrary to what we would naively expect if ι were the radial excitation of η'. First, the phase space factor for $\psi \to \gamma\iota$ is smaller by ~ 2 than for $\psi \to \gamma\eta'$. Second, the matrix element for a qq meson M to couple to two gluons, $\langle gg|M\rangle$, is proportional to the meson wave function at the origin, $\psi_M(0)$, which is smaller for a radial excitation than for the ground state. (For instance, the ratios $\Gamma(\psi' \to e^+e^-)/\Gamma(\psi \to e^+e^-)$ and $\Gamma(\rho' \to e^+e^-)/\Gamma(\rho \to e^+e^-)$ are both consistent with $|\langle M|e^+e^-\rangle|^2$ being half as big for the radial excitation as for the ground state.) So if ι were the radial excitation of η' I would expect $\Gamma(\psi \to \gamma\iota)$ to be a few times smaller than $\Gamma(\psi \to \gamma\eta')$ rather than at least two times bigger.

If ι is a $\bar{q}q$ meson, it becomes even harder to understand the large rate for $\psi \to \gamma\iota$ when we consider in addition the dominance of the $\iota \to \bar{K}K\pi$ decay mode. To understand the $\bar{K}K\pi$ decay we might assume that ι is predominantly an ss state rather than an approximate flavor singlet like η'. This would imply that ι is 1/3 flavor singlet, which both a naive application of flavor symmetry and the data for $\Gamma(\psi \to \gamma\eta)/\Gamma(\psi \to \gamma\eta')$ would suggest implies a further reduction of $\Gamma(\psi \to \gamma\iota)$.

These arguments against a predominantly $\bar{q}q$ interpretation of ι are plausible but they are not completely conclusive. Bethe-Salpeter calculations[29] give larger values than my naive estimate of $\psi \to \gamma + \eta'$ (radial excitation), though still not as large as the ratio of eqs. (3.9) to (3.11). Furthermore the weak-binding approximation needed for these calculations is very badly violated (of order one), so their applicability is not clear. It has also been suggested that ι could be a mixture of $\bar{q}q$ ground state and radial excitation,[30] though how this could explain the total experimental picture is not clear to me. If η' and ι contained appreciable $\bar{c}c$ components this would complicate the rate estimates for $\psi \to \gamma\eta'$ and $\psi \to \gamma\iota$. However since the amount of $\bar{c}c$ would be proportional to the flavor singlet components in η' and ι, it remains hard to understand both the inequality $\Gamma(\psi \to \gamma\iota) >> \Gamma(\zeta\psi \to \gamma\zeta)$ and the predominance of $\iota \to \bar{K}K\pi$. Finally the symmetry breaking effects discussed above for Fig. (1a) could also enhance $\psi \to \gamma + (\bar{s}s)_{0^-}$, though the observed ratio $\Gamma(\psi \to \gamma\eta)/\Gamma(\psi \to \gamma\eta')$ implies this is not a dominant effect.

The preceding paragraph is too brief for the details to be understandable, but I hope the main point is clear: the important properties of iota are simply understood if it is predominantly a glueball but the possibility of complicated qq interpretations is not completely excluded. So we are left with the central question: what further experimental studies are needed to decide the issue?

The theoretical answer is very easy to state but the necessary experimental program may be very difficult. If iota is a glueball, then there must be one more pseudoscalar in the mass region < 1.6 GeV to complete the π' nonet. It is essential to understand the π' nonet and especially the two isoscalar partners. Therefore the existence of $\zeta(1270)$ must be verified and, if ι is a glueball, the ninth member of the nonet must be found. Part of this work can be done with fixed target experiments. But high statistics studies of radiative ψ decay are most essential, because they are best able to tell us the degree of mixing between the $\bar{q}q$ isoscalars and the glueball states. What we would like, for example, is enough statistics for a partial wave analysis of channels like $\psi \to \gamma\eta\pi\pi$ and $\psi \to \gamma\bar{K}K\pi$ which would be sufficiently sensitive to see $\psi \to \gamma\zeta$ and $\psi \to \gamma\zeta'$ (where ζ' is the missing ninth member of the nonet) even if ζ and ζ' are purely radially excited $\bar{q}q$ states.

To have a rough idea of the possible signals I have made some simple-minded estimates based on the assumption that $\zeta(1270)$ and $\zeta'(?)$ are pure radially excited $\bar{q}q$ states. I consider two assumptions for the ζ-ζ' mixing angle: (1) 1-8 mixing, as is approximately true for η and η' and (2) ideal mixing as for ϕ and ω. I assume, as discussed below eq. (3.11), that the square of the matrix elements $|\langle M|gg\rangle|^2$ and $|\langle M|\gamma\gamma\rangle|^2$ are smaller by 2 for the radial excitation than the ground state (up to the simple corrections due to electric or color charges). I then estimate the $X \to \gamma\gamma$ and $\psi \to \gamma X$ widths for the radial excitations by comparing with the η and η' widths, corrected for phase space and for the appropriate $\gamma\gamma$ and gg coupling factors. Thus I neglect the possible complications listed above.

Suppose first that approximate 1-8 mixing holds, so that

$$\zeta \cong \zeta_8 = \frac{1}{\sqrt{6}}\left(\bar{u}u + \bar{d}d - 2\bar{s}s\right)$$

$$\zeta' \cong \zeta_1 = \frac{1}{\sqrt{3}}\left(\bar{u}u + \bar{d}d + \bar{s}s\right) \tag{3.12}$$

In this case the ζ-ζ' system is controlled by the same nonperturbative physics that controls the "U(1) problem" and the η' mass, and it is very hard to estimate the ζ' mass. I will simply guess $m_{\zeta'} = 1500$ MeV and put $m_\zeta = 1270$ MeV from ref. (28). Then naively (3.12) implies $\Gamma(\psi \to \gamma\zeta') >> \Gamma(\psi \to \gamma\zeta)$ and the above assumptions give

$$B(\psi \to \gamma\zeta') \sim 2 \cdot 10^{-3} \tag{3.13}$$

If for example ζ-ζ' mixing were just like η-η' mixing, so that ζ were exactly the excitation of η, my assumptions would give

$$\frac{B(\psi \to \gamma\zeta)}{B(\psi \to \gamma\eta)} \sim \frac{1}{2}\left[\frac{m_\psi^2 - m_\zeta^2}{m_\psi^2 - m_\eta^2}\right]^3 \tag{3.14}$$

or $B(\psi \to \gamma\zeta) \sim 3 \cdot 10^{-4}$, an estimate meant only to suggest the possible order of magnitude. In this case we expect ζ and ζ' to both appear in $\eta\pi\pi$ and $\bar{K}K\pi$, and ζ' may also have a sizeable decay to $\eta'\pi\pi$. Similarly I find

$$\Gamma(\zeta' \to \gamma\gamma) \sim 14 \text{ keV} \tag{3.15}$$

$$\Gamma(\zeta \to \gamma\gamma) \sim 5 \text{ keV} \tag{3.16}$$

and using the OIZ rule

$$\sigma(\pi p \to \zeta' n) \sim 2\,\sigma(\pi p \to \zeta n) \tag{3.17}$$

The ideal mixing assumption is

$$\zeta = \frac{1}{\sqrt{2}} \left(\bar{u}u + \bar{d}d \right) \tag{3.18}$$

$$\zeta' = \bar{s}s \tag{3.19}$$

In this case the nonperturbative dynamics of the U(1) problem presumeably does not contribute and we can estimate $m_{\zeta'}$ from the ideal mixing mass formula

$$m_{\zeta'} = 2m_{K'} - m_{\zeta} \sim 1600 \text{ MeV} \tag{3.20}$$

Then the results are

$$B(\psi \to \gamma\zeta') \sim 6 \cdot 10^{-4} \tag{3.21}$$

$$B(\psi \to \gamma\zeta) \sim 2 \cdot 10^{-3} \tag{3.22}$$

$$\Gamma(\zeta' \to \gamma\gamma) \sim 3 \text{ keV} \tag{3.23}$$

$$\Gamma(\zeta \to \gamma\gamma) \sim 13 \text{ keV} \tag{3.24}$$

$$\sigma(\pi p \to \zeta n) \gg \sigma(\pi p \to \zeta' n) \tag{3.25}$$

In this case ζ' decays predominantly to $\bar{K}K\pi$ and ζ to $\eta\pi\pi$.

Of course none of these numbers should be taken very seriously. They are only meant to suggest the range of orders of magnitudes that could occur depending on the ζ-ζ' mixing. If ζ and/or ζ' also contained glueball components, we would expect correspondingly smaller $\gamma\gamma$ widths and larger yields in radiative ψ decay.

It is interesting to compare $\Gamma(\psi \rightarrow \gamma X)$ with $\Gamma(X \rightarrow \gamma\gamma)$ since in lowest order these probe respectively the Xgg and $X\gamma\gamma$ couplings. An interesting quantity, with phase space effects removed, is the "stickiness coefficient"

$$S_X = \frac{\Gamma(\psi \rightarrow \gamma X)}{\Gamma(X \rightarrow \gamma\gamma)} \cdot \frac{LIPS(X \rightarrow \gamma\gamma)}{LIPS(\psi \rightarrow \gamma X)} \tag{3.26}$$

where LIPS means Lorentz Invariant Phase Space. S_X is useful because the wave function at the origin $|\psi_X(0)|^2$ cancels, so that stickiness measures more directly the relative strengths of the Xgg and $X\gamma\gamma$ couplings. S_X is also useful experimentally because if X is detected in the same final state $X \rightarrow AB \ldots$ in both $\psi \rightarrow \gamma AB \ldots$ and $\gamma\gamma \rightarrow AB \ldots$ then the branching ratio $B(X \rightarrow AB \ldots)$ cancels and need not be known to measure S_X. It is particularly useful to consider the ratio of ratios, S_X/S_Y, which is a figure of merit for the relative "glueness" versus "quarkness" of the states X and Y. For instance if G and M are respectively a glueball and $\bar{q}q$ meson of the same quantum numbers we would certainly expect

$$S_G \gg S_M \tag{3.27}$$

In fact we are almost in a position to measure $S_\theta/S_{f'}$, since $S_{f'}$ is already known[6,7] and only $\gamma\gamma \rightarrow \theta \rightarrow \bar{K}K$ must be measured to determine S_θ. We would also obviously like to measure S_ι, S_ζ, and $S_{\zeta'}$. $S_{\eta'}$ is known while there is a lower limit on S_ι (based on an upper limit for $\gamma\gamma \rightarrow \iota \rightarrow \bar{K}K\pi$). Substituting the experimental values I find

$$S_\iota / S_{\eta'} \gtrsim 4\tfrac{1}{2} \tag{3.28}$$

which is consistent with the glueball interpretation of iota.

To aid in understanding which states are glueballs it would also be helpful to study systemically the relative excitation cross sections in hadronic reactions. For instance, the OIZ rule predicts the ratio $\sigma(\pi^- p \rightarrow \eta n)/\sigma(\pi^- p \rightarrow \eta' n)$ as a function of the η-η' mixing angle. The experimental value is in reasonably good agreement with the value ~ 1 expected for $\theta_{\eta\text{-}\eta'} = -11^0$. Similarly the ratio $\sigma(\pi p \rightarrow \zeta n)/\sigma(\pi p \rightarrow \zeta' n)$ will help us to check our understanding of the π' nonet. It would also be useful to measure or bound production of glueball candidates in hadronic channels such as $\pi p \to Gn$. In particular a bound or measurement for $\pi p \to \iota n$ would be very interesting. We could then consider the "hadron coefficient",

$$H_X = \frac{\Gamma(\psi \to \gamma X)}{\sigma(\pi p \to X n)} \tag{3.29}$$

The significance of H_X is perhaps less clear than S_X but it is again likely that we would have

$$H_G \gg H_M \tag{3.30}$$

for G a pure glueball and M a pure meson in the same channel.

This most recent data on iota from the Mark III succeeds in raising more questions than it answers. The Mark III sees a somewhat broader and heavier iota than did the Mark II and Crystal Ball, because of different cuts on the $\bar{K}K$ mass. The $\rho\gamma$ signal in the iota region is also interesting and confusing. If the entire signal is attributed to the iota, then vector dominance arguments imply

$$\Gamma(\iota \to \gamma\gamma) \sim 19 \text{ keV} \cdot B(\iota \to \bar{K}K\pi) \tag{3.31}$$

with large errors, whereas the direct upper bound from TASSO is

$$\Gamma(\iota \to \gamma\gamma) < 7 \text{ keV}/B(\iota \to \bar{K}K\pi) \tag{3.32}$$

To reconcile (3.31) with (3.32) we would need $B(\iota \to \bar{K}K\pi) < 2/3$, implying substantial signals to $\iota \to \eta\pi\pi$, $\eta'\pi\pi$, or $\pi\pi\pi\pi$, which have not yet been seen. The available data for $\eta\pi\pi$ and 4π make large rates for $\iota \to \eta\pi\pi$ and $\iota \to 4\pi$ unlikely, but no data has been presented yet for $\eta'\pi\pi$. If $B(\iota \to \eta'\pi\pi)$ is also small, then it is most probable that the entire $\rho\gamma$ signal is not due to iota. It is also conceivable that the Mark III $\bar{K}K\pi$ signal in the iota region is due to more than one state. This can be tested (again with more statistics!) by studying the $\bar{K}K\pi$ Dalitz plot as a function of $\bar{K}K\pi$ mass.

Other remarkable new results from the Mark III include the structure(s) in $\eta\pi\pi$ at 1380 and perhaps also 1270, the possible $J^P = 0^-$ $\rho\rho$ and $\omega\omega$ signals at 1.6 - 1.9 GeV, and the $\xi(2220)$. Again we have more questions than answers. The answers will require a more complete program of partial wave analysis than has been attempted so far. The difficulty and importance of this program is the guarantee that there will be interesting and challenging physics for BEPC at the end of the decade.

IV. Meiktons

The gluonic degrees of freedom might also be observed by finding the mixed $\bar{q}qg$ states[16,31,32] which I will call meiktons, the classical Greek term for a mixed object. I will briefly describe the bag model predictions for the ground state meikton nonet[16,32] and for a class of excited nonets[21] which would have characteristic, experimentally distinguishable decays. If these meiktons are found it would confirm the existence of valence gluons in the very particular form required by the bag model.

The idea of valence gluons is a controversial one. 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 $\bar{q}q\bar{q}q$, $\bar{q}q\bar{q}q\bar{q}q$ ... 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.[27] 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 semiquantitative 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[33] 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 $J^P = 1/2^+$ 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 qq pair, either the spin singlet with $J^{PC} = 0^{-+}$ or 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 have now computed the masses of these nonets through $0(\alpha_s)$ in cavity perturbation theory and are in agreement except for differences in the treatment of quark and gluon self energies.[16.32] The results from Reference (16) are shown in Table 1 for three values of the ratio of gluon mode self energies $C_{TE}/C_{TM} = 2, 1, 1/2$. This ratio is fixed if we assume that $\theta(1700)$ is the TE^2 glueball, with $C_{TE}/C_{TM} \sim 1/2$ if θ has spin 2 and $C_{TE}/C_{TM} \sim 2$ if θ has spin 0. Table 2 shows the predicted glueball spectrum from the same calculation.

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 two qq 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 singet $\bar{q}q$ 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 as discussed above). It is therefore particularly interesting to look for the states of the 1^{-+} nonet. The quantum numbers of these four ground state nonets, 1^{--} and $(0,1,2)^{-+}$, are a specific test of the bag model because they follow from the axial vector quantum numbers of the TE gluon mode.

These $\bar{q}qg$ states are likely to decay by formation of a $\bar{q}q$ pair from the gluon, $g \rightarrow \bar{q}q$, 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 $L=1$ meson in a relative s-wave,

$$\bar{q}_s q_s g_{TE} \rightarrow \begin{cases} (\bar{q}q)_s + (\bar{q}q)_s & L=1 \\ (\bar{q}q)_s + (\bar{q}q)_p & L=0 \end{cases} \tag{4.1}$$

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

$$\text{"}\omega\text{"}\,(1^{-+}) \rightarrow \begin{cases} \eta\eta' & L=1 \\ \pi A_1 & L=0 \end{cases} \tag{4.2}$$

Notice that $\eta\eta'$ in a p=wave uniquely signals the 1^{-+} quantum numbers.

Since the TE gluon s-channel couplings to qq are approximately flavor symmetric, (see Table 1 of Ref. (16)), the $\bar{q}qg_{TE}$ meiktons may have characteristic multi-kaon and apparent OIZ violating decays. As for the $\bar{q}qg_{TM}$ states discussed below, but to a lesser extent, the qqg_{TE} states may have decays such as "ρ" $(1^{-+}) \rightarrow \pi E$, $\bar{K}K^*$; "ρ"$(1^{--}) \rightarrow \pi\phi$, $\bar{K}K$: and "ρ"$(2^{-+}) \rightarrow \pi f'$, $\bar{K}K^*$. The latter are of particular interest for the A_3-A_3' candidate discussed below.

The TM(transverse magnetic) gluon mode has vector quantum numbers, $J^{PC} = 1^{--}$, and mode energy 4.49/R. For R ~ 1 fm., the $\bar{q}qg_{TM}$ nonets should be a few hundred MeV heavier than the $\bar{q}qg_{TE}$ nonets. They are of special interest because, as seen in Table 1 of Ref. (16), the s-channel coupling $g_{TM} \rightarrow \bar{s}s$ is bigger by ~ 5 in amplitude than $g_{TM} \rightarrow \bar{u}u, \bar{d}d$. In taking this result seriously we are escalating our reliance on the spherical cavity approximation to the bag model but with a potentially great reward: if the predicted enhancement is even qualitatively correct then many $\bar{q}qg_{TM}$ meiktons will have spectacular decay signatures by which they can be clearly distinguished from $\bar{q}q$ mesons of the same quantum numbers. As already discussed in Section II, the dominance of $\iota \to \bar{K}K\pi$ is consistent with the TM $\to \bar{s}s$ enhancement and the interpretation of ι as a $J^{PC} = 0^{-+}$ TE-TM glueball.

In Ref. (21) the spectum of $\bar{q}qg_{TM}$ meiktons and TM^2 glueballs was computed to $0(\alpha_s)$ in cavity perturbation theory, using the same approximations and parameters that were applied in Ref. (16) to the $\bar{q}qg_{TE}$, TE^2, and TE-TM states. The results for the spectrum are shown in Table 3, as a function of C_{TE}/C_{TM} as before. There are four $\bar{q}qg_{TM}$ nonets with the same quantum numbers as the p-wave $\bar{q}q$ states, $J^{PC} = 1^{+-}, (0,1,2)^{++}$ (mixing with the $\bar{q}q$ p-wave is incorporated to $0(\alpha_s)$ and is small). Their masses range between 1.8 and 2.5 GeV for $C_{TE}/C_{TM} = 1/2$ and between 1.4 and 2.2 GeV for $C_{TE}/C_{TM} = 2$.

Some "signature" decay modes are shown in Table 4. The ϕ-like TM meiktons decay to final states with four K's including $\phi\phi$, so the "ϕ"(2^{++}) might be identified with one of the Brookhaven $\phi\phi$ candidates.[34] The strange qqg_{TM} states decay to three kaon final states, including ϕK and ϕK^*; these are the natural prey of high statistics kaon beam experiments such as LASS. The isovectors and ω-like isoscalars decay to final states containing a KK pair. The KK pair may materialize as a ϕ meson, either by final state interaction or directly by soft gluon emission from the color octet $\bar{s}s$ pair created by the $J^{PC} = 1^{--}$ TM gluon. These decays, such as "ρ"$(1^{+-}) \to \phi\pi$ or "ρ"$(0,1,2)^{++} \to \phi\rho$ are unmistakeable, since they would be OIZ forbidden decays for $\bar{q}q$ isovectors. Similarly "ω"$(0,1,2)^{++} \to \omega\phi$ would be an OIZ forbidden decay for any $\bar{q}q$ isoscalor.

Taking perturbation theory as a guide, the $I = 0$, $C = +$ meiktons are produced in radiative ψ decay with a rate intermediate between that of glueballs and $\bar{q}q$ mesons. We may hope to identify at least some of them by the unique signature decay modes discussed above. As with the $\bar{q}q$ mesons, identification will also depend on understanding the total picture of the $\bar{q}qg$ nonets, for which we rely on fixed target experiments and perhaps also (as discussed below) on <u>non</u>-radiative ψ decays.

There are several meikton candidates in the experimental literature, which I will now briefly discuss:

(1) A possible assignment for $\xi(2220)$, seen by the Mark III in $\psi \to \gamma K^+K^-$, is as a $\bar{q}qg_{TM}$ meikton, with $J^{PC} = 0^{++}$ or 2^{++}. As shown in Table 3 for $C_{TE}/C_{TM} = 1/2$ (corresponding to θ being the 2^{++} TE^2 glueball) the estimated masses are 1900 MeV for "ω"(0^{++}) and 2300 MeV for "ω"(2^{++}). The signature decays of Table 4 include "ω"$(0^{++}) \to \bar{K}K$ while "ω"(2^{++}) does not decay to $\bar{K}K$ in lowest order but can by single gluon exchange (a kind of color M-1 transistion, $\bar{K}_8{}^*K_8{}^* \to \bar{K}_1K_1$, where the subscripts denote color representations). For either spin we also expect $\xi \to \bar{K}^*K^*$, not a very striking prediction. However we also expect in Table 4 the very peculiar decay $\xi \to \phi\omega$. This assumes, beyond the lowest order decay mechanism in which $gq\bar{q} \to q\bar{q}q\bar{q}$ which "falls apart", that the ϕ forms either by final state interaction or by soft gluon exchange, $\phi_8\omega_8 \to \phi_1\omega_1$. The decay $\xi \to \phi\omega$ would be an OIZ suppressed decay if ξ were a $\bar{q}q$ state, since both $\bar{u}u + \bar{d}d \to \phi\omega$ and $\bar{s}s \to \phi\omega$ are OIZ suppressed. A rough estimate, based on cavity perturbation theory, of the widths of $\bar{q}qg_{TM}$ meiktons is consistent with a width of order 30 MeV. It would be very interesting to search for the decay $\psi \to \gamma\omega\phi$.

(2) The C(1430) is seen[35] as an 8σ signal in $\pi^-p \to Cn \to \pi^0\phi n$. $C \to \pi\phi$ is one of the signature decay modes discussed above. If the π and ϕ are in an s-wave, C(1430) could be the "ρ" of the $J^{PC} = 1^{+-}$ $\bar{q}qg_{TM}$ nonet, none of which could be produced in $\psi \to \gamma X$ because of C-parity. It could however be searched for in the

direct ψ decay,
$\psi \to C\pi$ which would be an s-wave decay if $J^{PC}(C) = 1^{+-}$.

(3) The 2^{++} $\phi\phi$ structures seen at Brookhaven[34] could be identified with the "ϕ"(2^{++}) $\bar{s}sg_{TM}$ meikton or with the $g_{TM}g_{TM}$ 2^{++} glueball. In either case they should be produced in $\psi \to \gamma\phi\phi$.

(4) The 0^{++} K_sK_s resonance[36] at 1770 MeV might be identified with "ω"(0^{++}), a qqg_{TM} state, which could occur in $\psi \to \gamma K\bar{K}$.

(5) Together with E(1420) the recently claimed $1^{++}K^*\bar{K}$ resonance,[37] D′(1526), makes too many states for the A_1 nonet. The D′(1526) could be the "ω"(1^{++}) $\bar{q}qg_{TM}$ state.

(6) The ACCMOR experiment,[8] with their sample of 600,000 $\pi^+\pi^-\pi^-$ events, confirmed the existence of the $A_3(1700)$, a $J^{PC} = 2^{-+}$ isovector state presumed to be the $\bar{q}q$ d-wave. However they also saw a second possible I = 1, $J^{PC} = 2^{-+}$ structure at 1850 MeV, too small a mass splitting for the second to be the radial excitation of the first. A possible interpretation is that the $A_3(1700)$ and $A_3'(1850)$ (assuming there really is a resonance at 1850) could be mixtures of the qq d-wave and the "ρ"(2^{-+}) $\bar{q}qg_{TE}$ meikton. If the mixing occured by the $f\pi$ s-wave intermediate channel, one of the resulting eigenstates would decouple from the $f\pi$ s-wave, as is indeed observed for the structure at 1850 MeV. The isocalar members of this nonet could be produced in radiative ψ decay in a p-wave, though $J^{PC} = 2^{-+}$ is not present in the lowest order amplitude for $\psi \to \gamma gg$. Possible channels to consider are $\psi \to \gamma\bar{K}K^*$, $\psi \to \gamma\eta f$, and $\psi \to \gamma\pi A_2$.

Direct ψ decays may be useful to search for some of the meikton states. I have already mentioned

$$\begin{array}{l} \psi \to \text{"}\rho\text{"}(1^{+-})\,\pi \\ \qquad\quad \hookrightarrow \phi\pi \end{array} \tag{4.3}$$

in connection with C(1430) above. Another interesting example is to search for the exotic "ρ"(1^{-+}) $\bar{q}qg_{TE}$ state in its characteristic $\pi\eta$ decay mode. This mode provides a beautiful signature since $\pi\eta$ in a p-wave is uniquely an exotic $J^{PC} = 1^{-+}$ isovector. A promising channel is the electromagnetic (but <u>not</u> radiative) ψ decay

$$\begin{array}{l} \psi \to \gamma^* \to \text{"}\rho\text{"}(1^{-+})^{\pm}\,\pi^{\mp} \\ \qquad\qquad\qquad \hookrightarrow \eta\pi^{\pm} \end{array} \tag{4.4}$$

with "ρ" and π in a relative p-wave. This decay must proceed by a virtual intermediate photon because of the positive G-parity of the final state. Recall that about 20% of the hadronic decays of ψ proceed via a virtual photon, and 3/4 of these produce I = 1, G = + final states.[38]

Very little work has been done so far using non-radiative hadronic ψ decays to search for new particles. This could be a promising new area for future high statistics studies.

V. $\bar{q}\bar{q}qq$ Exotics

It was possible to discover quarks from the hadron spectrum known in the early 60's because the known states could be identified with simple $\bar{q}q$ mesons and qqq baryons. Many years later and with many more states discovered, the simple classification scheme is still remarkably successful. Success is always gratifying but this success is also puzzling. What about more complicated states, such as the four quark exotics made of $\bar{q}\bar{q}qq$? Does QCD predict their existence or not? Exotic quantum numbers, such as Q = 2 or S = 2, would make them easy to detect. Is it a success or a failure that they have not yet been found?

A neat solution to this puzzle has been given in the Bag model.[39,40] The solution has two parts:

1) The lowest-lying $\bar{q}\bar{q}qq$ states do not have exotic quantum numbers, but form nonets with the same net quantum numbers as $\bar{q}q$ nonets — they are called "crypto-exotic" nonets.

2) Most of the $\bar{q}\bar{q}qq$ states — both the truly exotic and the cryptoexotic — can "fall apart" into two constituent color-singlet $\bar{q}q$ mesons and are consequently too broad to detect as S-matrix poles.[41] The existence of the low-lying crypto-exotic nonets is implied by the hyperfine splitting due to single gluon exchange, the same approximation which gives a good qualitative description of the L + 0 hadrons. In this approximation, it is not hard to see[42] why the qqqq ground state turns out to be a $J^{PC} = 0^{++}$ scalar nonet.

The quark eigenmodes are classified by the group $SU(3)_{color} \times SU(2)_{spin} \times SU(3)_{flavor}$. It is useful to consider $SU(6)_{color\text{-}spin}$ which contains $SU(3)_{color} \times SU(2)_{spin}$ and to classify states by $SU(6)_{color\text{-}spin} \times SU(3)_{flavor}$. Where λ and σ denote the eight color and three spin matrices, the energy-shift due to single gluon exchange is

$$\Delta E = -K \frac{\alpha_s}{R} \sum_{i,j} \langle \lambda_i \cdot \lambda_j \; \vec{\sigma}_i \cdot \vec{\sigma}_j \rangle \tag{5.1}$$

K is a flavor-dependent constant and the sum is over all qq, $\bar{q}q$, and $\bar{q}\bar{q}$ pairs (i,j). In analogy to the SU(2) relation for a $\bar{q}q$ bound state

$$\vec{s}_1 \cdot \vec{s}_2 = -\frac{1}{2}\left[\vec{S}^{\,2}_{total} - \vec{s}_1^{\,2} - \vec{s}_2^{\,2} \right] \tag{5.2}$$

the expectation value in Eq. (5.1) may be rewritten in terms of $SU(6)_{color\text{-}spin}$ Casimir operators

$$\Delta E = K \frac{\alpha_s}{R} \left[\frac{1}{2} C_6(TOT) - C_6(qq) - C_6(\bar{q}\bar{q}) + \dots \right] \tag{5.3}$$

For simplicity I have displayed only the largest terms in Eq. (5.3); contributions of $SU(2)_{spin}$ and $SU(3)_{color}$ Casimir operators are omitted. C_6 is the sum of the squares of the 35 SU(6) generators, the analogue of $S(S + 1) = \Sigma\, \sigma_i^2$ for SU(2). C_6 dominates Eq. (5.3) just because SU(6) has more generators than SU(3) and SU(2).

The quantum numbers of the ground state are easily obtained from Eq. (5.3) and Fermi statistics. Since $C_6(qq)$ and $C_6(\bar{q}\bar{q})$ appear with a minus sign we want to maximize them. The largest Casimir for a diquark is obtained from the symmetric representation, the 21 in $6 \times 6 = 21 + 15^*$, [just like in SU(2), $2 \times 2 = 3 + 1$, where the triplet is symmetric and the scalar antisymmetric]. But if the diquark is symmetric under $SU(6)_{color\text{-}spin}$, Fermi statistics require that it be antisymmetric under $SU(3)_{flavor}$, *i.e.*, in the 3^* of $3 \times 3 = 6 + 3^*$. Therefore qq is in the flavor 3^*, $\bar{q}\bar{q}$ is in the flavor 3, and the ground state $\bar{q}\bar{q}qq$ is in a flavor nonet, $3 \times 3^* = 8 + 1$!

The spin of this nonet is determined by $C_6(TOT)$, the first term in Eq. (5.3). Since it contributes positively to ΔE we want to minimize it. This is achieved when the total state is an $SU(6)_{color\text{-}spin}$ singlet, in which case it is also a singlet of $SU(2)_{spin}$, that is $J = 0$. P and C are then positive because all four constituents are in an s-wave. The conclusion is that the lowest-lying $\bar{q}\bar{q}qq$ states have precisely the same quantum numbers as the $J^{PC} = 0^{++}$ nonet formed from $\bar{q}q$ in a p-wave!

Although this crypto-exotic nonet has the same net quantum numbers as the p-wave scalar nonet, its exotic quark content give it properties very different from the $\bar{q}q$ nonet. The quark content and estimated masses are shown in figure (2). Notice in particular the degenerate isoscalar and isotriplet at 1100 MeV., which are just the usual ideally mixed non-strange isoscalar and isotriplet plus an $\bar{s}s$ pair. Unlike the non-strange isoscalar and isotriplet of a $\bar{q}q$ nonet, these $\bar{q}\bar{q}qq$ states will couple strongly to $K\bar{K}$.

There is a good chance that at least some of the members of the 0^{++} cryptoexotic nonet have already been observed. A plausible hypothesis[39,40] is that the S*(975) and δ(980) are the $I = 0$ $\bar{s}s(\bar{u}u + \bar{d}d)$ and $I = 1$ $\bar{s}s\bar{u}d$, ... states of figure (2). This hypothesis explains in a simple way why the S* and δ are nearly degenerate yet very strongly coupled to $K\bar{K}$ (they are below $K\bar{K}$ threshhold but cause strong threshhold enhancements), properties which cannot be explained in a simple way if S* and δ are $\bar{q}q$ states. But then where are the predicted $\varepsilon = \bar{u}u\bar{d}d$ and $\kappa = \bar{u}s\bar{d}d$, ... states which are expected at lower masses? There is now no evidence for these states in ππ and Kπ phase shift analyses.

This question also has a simple answer,[41] one which suggests that most $\bar{q}\bar{q}qq$ states will not be observable as ordinary resonances. The point is that δ and S* at 980 MeV., 10% below the bag model estimate for their masses of 1100 MeV., are below the $K\bar{K}$ threshhold at 990 MeV. If they were above the $K\bar{K}$ threshhold, the four q's and q's would not be confined since they could pair off into a $K\bar{K}$ pair. They would therefore "fall apart", with a decay width of the order of their mass. Since the bag model estimates of the ε and κ masses are very far above their respective fall-apart ππ and Kπ threshholds, they presumeably do fall apart and are unobservable as S-matrix poles (though perhaps observable with the P-matrix" analysis[41]). In this picture how do S* and δ decay? For S* the only channel is S* → ππ (4π is possible but much suppressed by phase space), which requires an OIZ violation to annihilate the ss quark pair. OIZ suppression would then explain the narrow width found for the S* in a measurement[43] of ψ decay data, 14 ± 5 MeV., and an analysis[44] of ππ scattering which gives ~ 8

MeV. The δ has a component which can fall apart to $\eta\pi$; this could explain its larger width. In fact I would expect it to be broader[27] than the 50 MeV. width seen in $\delta \rightarrow \eta\pi$. There is an old suggestion[45] that 50 MeV. is not the true δ width but an effect of unitarity and analyticity and that the true width is much bigger than 50 MeV. This could be tested by a high statistics study of the I = 1 s-wave $\bar{K}K$ threshhold enhancement.

The lesson I would draw from this discussion of the light scalar mesons is that most $\bar{q}\bar{q}qq$ states will have uninhibited fall-apart decays, making them too broad to observe as ordinary resonances. The $\bar{q}\bar{q}qq$ states we may hope to see are those few, like S* and δ perhaps, which happen to lie below the threshholds of their principal fall-apart decay modes.

As B.A. Li and K.F. Liu have discussed in a series of papers,[47] some $\bar{q}\bar{q}qq$ states may be produced in $\gamma\gamma$ scattering and in radiative ψ decay. In perturbation theory $\bar{q}\bar{q}qq$ states are produced in $\psi \rightarrow \gamma X$ in the same order in α_s as $\bar{q}q$ mesons. To see which $\bar{q}\bar{q}qq$ states are most likely to be produced in $\psi \rightarrow \gamma X$ we examine $\alpha, \beta, \gamma, \delta$, the "recoupling coefficients", defined by[40]

$$|\bar{q}\bar{q}qq\rangle = \alpha P_1 P_1 + \beta P_8 P_8 + \gamma V_1 V_1 + \delta V_8 V_8 \tag{5.4}$$

Here we consider only the s-wave states discussed above. P_1 denotes a spin singlet (pseudoscalar) $\bar{q}q$ pair in a color singlet while P_8 denotes the color octet. V_1 and V_8 denote the analogous spin triplet (vector) $\bar{q}q$ pairs. Of course the normalization condition is

$$\alpha^2 + \beta^2 + \gamma^2 + \delta^2 = 1 \tag{5.5}$$

Eq. (5.4) is to be understood for a particular flavor ordering of the quarks. Then provided the flavor content is such that we can arrange the $\bar{q}\bar{q}qq$ state into two flavor neutral $\bar{q}q$ pairs, the probability to produce the state in $\psi \rightarrow \gamma X$ is proportional to δ^2. This is a kind of "gluon vector dominance". Li and Liu have done extensive calculations of the $\gamma\gamma$ excitation cross sections which are proportional to the γ^2 terms.

There is in fact a cryptoexotic candidate among the interesting new particles discovered in radiative ψ decay. This is the $\theta(1700)$, discovered initially in $\psi \rightarrow \gamma\eta\eta$ with

$$B(\psi \rightarrow \gamma\theta) \cdot B(\theta \rightarrow \gamma\gamma) \sim \frac{1}{2} \cdot 10^{-4} \tag{5.6}$$

The rate and the prominence of the two body decay mode immediately suggested[47] the possibility that θ could be a cryptoexotic state with a large value of γ^2 but below the threshhold for its V_1V_1 fall apart decay. Using Jaffe's bag model calculations there are three candidates which satisfy these conditions, all with the flavor content $\bar{s}s(\bar{u}u + \bar{d}d)$. All three states have large recoupling coefficients to fall apart to $\phi\omega$ or $\bar{K}^*K^*$ but at 1700 MeV would be below the threshholds to do so. For both states the principal OIZ allowed decays are $\eta\eta$ and $\bar{K}K$. Taking the usual η-η' mixing angle of $-11°$ I found

$$B(\theta \to \bar{K}K) \sim 2\, B(\theta \to \eta\eta) \gg B(\theta \to \pi\pi) \tag{5.7}$$

which has since been confirmed by Mark II[23] and Mark III[6] measurements of $\theta \to \bar{K}K$. (Depending on whether some of the $\eta\eta$ signal is attributed to f′, the experimental ratio of $\bar{K}K : \eta\eta$ might be somewhat larger than 2.)

Of the three candidates, one is a spin zero state [$C_s(9^*)$ in the notation of ref. (40)] with a small P_1P_1 recoupling coefficient, $\alpha^2 = .178^2 = .03$. The other two candidates are $J = 2$ states [$C_s(9)$ and $C_s(36)$] which are degenerate in mass in the approximation used in ref. (40). These states have no fall apart decays to two pseudoscalars, $\alpha^2 = 0$, simply because construction of an s-wave $J = 2$ $\bar{q}\bar{q}qq$ state requires that both $\bar{q}q$ pairs be in spin triplets. Therefore these two states can only decay to $\bar{K}K$ and $\eta\eta$ by gluon exchange. The lowest order mechanism is for $\bar{K}_8^*K_8^*$ to scatter by t-channel gluon exchange to $\bar{K}_1K_1$, a kind of color (double) M_1 transition.[47] Because of the flavor symmetry of gluon exchange, decays by this mechanism will also satisfy Eq. (5.7).

The observation[48] by the Mark II of a large signal for $\psi \to \gamma\rho\rho$ in roughly the region of the θ suggested that θ might be a glueball, since combining the $\bar{K}K$, $\eta\eta$, and possible $\rho\rho$ signals implied $B(\psi \to \gamma\theta) \geq 6 \cdot 10^{-3}$, a very large rate. However the recent analysis[7] of the angular distribution of the $\rho\rho$ signal by the Mark III shows it to be predominantly negative parity and probably $J^P = 0^-$, whereas the θ must have positive parity because of the $\eta\eta$ decay mode. This shows again the crucial importance of partial wave analysis. With the new result the only known decays are $\theta \to \eta\eta$ and $\bar{K}K$, and the cryptoexotic interpretation is again attractive. The $J = 2$ alternative is especially intriguing since it implies that the θ may be two states, which could give rise to interesting interference effects that could be different in $\psi \to \gamma\theta$ and $\gamma\gamma \to \theta$.

Continuing with the cryptoexotic hypothesis, the OIZ violating decays $\theta \to \rho\rho, \omega\omega$ occur to order α_s is amplitude, and will therefore be less strongly suppressed than the OIZ violating decays of ordinary $\bar{q}q$ mesons. The initial configuration $\phi_8\omega_8$ can become $\omega_8\omega_8$ by an s-channel gluon exchange, and the subsequent $\omega_8\omega_8$ can then fall apart to $\omega\omega$ and $\rho\rho$. In cavity perturbation theory one often finds that s-channel exchange amplitudes are smaller than the usually softer t-channel exchange amplitudes, so although $\theta \to \rho\rho, \omega\omega$ are of the same order in α_s as $\theta \to \bar{K}K, \eta\eta$ (for the $J = 2$ case), they probably occur with smaller branching ratios.

The cryptoexotic states are interesting in their own right, and radiative ψ decay is one of the best places to look for them. Considered as a background to the glueball search they are distinguished by the following considerations:

1) Most $\bar{q}\bar{q}qq$ states will be essentially unconfined and will not be visible as ordinary resonances with reasonably small decay widths.
2) Those few $\bar{q}\bar{q}qq$ states which lie below their principal fall apart threshholds will tend to have an unusually large fraction of two body PP and VV decays (and also PV in the case of spin 1 cryptoexotics which are not expected to be produced strongly in $\psi \rightarrow \gamma X$).
3) The rate to produce cryptoexotics in radiative ψ decay should be much less than the rate to produce the most prominent glueball states.

VI. Charmonium

Although most of this talk is concerned with the use of ψ decay to look for and study new particles, I will briefly discuss in this section some physics issues pertaining to the study of the charmonium spectrum itself. A fuller discussion is given by Rosner in his paper on BEPC physics.[49]

Although the nonrelativistic potential model description of the cc spectrum has been a qualitative and even a semi-quantitative success, there are two factors which limit its applicability to charmonium.[50] One is the relativistic corrections, controlled by v^2/c^2, which is variously estimated at between 0.2 and 0.4 for the cc states. This means that even a first order treatment of relativistic corrections (which is all that has been attempted) is not enough to yield precise quantitative predictions. These corrections are much smaller for $\bar{b}b$ and $\bar{t}t$ systems. The second problem is the effect of light $\bar{q}q$ pairs, which is outside the scope of the potential model description. The Cornell group has approximated these effects by estimating the mixing of the cc states with the virtual continuum states such as $\bar{D}D$. These corrections are of the same order as the (spin-independent) relativistic corrections for charmonium.

Neither of these corrections can be computed precisely for charmonium though both can be roughly estimated. This imposes a practical limitation, which is not likely to change by the end of the decade, on how precisely we should try to measure the masses, transitions rates and decay widths of the charmonium system. It suggests that we concentrate on those remaining questions which are of a qualitative nature. I will briefly discuss a few of these below.

One qualitative puzzle is the nature of the cc spectrum above 4.0 GeV. Vector meson cc structures are seen at 4030, 4160 and 4415 MeV. One interpretation[51] is that the 4030 and 4415 are the 3S and 4S radial recurrences while the 4160 is the 2D state. Other new states suggested for this region are string excitations[52] and the perhaps related possibility of ccg states.[53] The question of S versus D states can be studied by measuring the $\bar{D}D$, $\bar{D}D^* + \bar{D}^*D$ and $\bar{D}^*D^*$ cross sections as a function of energy.[51] Some of the observed structure could also be due to F and F* production which has not yet been carefully studied in this region.

The transitions $\psi \to \gamma\chi$ and $\chi \to \gamma\psi$ have been a particular problem for the potential models, with the widths $\Gamma(\psi \to \gamma\chi)$ typically overestimated by a factor ~ 2. This problem appears now to be solved by adding to the Cornell coupled-channel model the effect of the relativistic corrections on the wave function overlap integrals.[54] The agreement with experiment is now at the $\sim 30\%$ level, which is as much as might be expected given the approximations involved. This suggests that these transitions are as well understood as we can expect. This view is further supported by the experimental agreement of the ratios $\Gamma(\psi' \to \gamma\chi_2):\Gamma(\psi' \to \gamma\chi_1):\Gamma(\psi' \to \gamma\chi_0)$ with the $(2J + 1)k^3$ behavior expected for E_1 transitions. Measurements of angular correlations in $\psi' \to \gamma\chi_J \to \gamma\gamma\psi$ would provide a further check of E_1 dominance.

An outstanding challenge is to find the lightest, still undiscovered charmonium state, the spin singlet $^1P_1$, $J^{PC} = 1^{+-}$. The mass of this state provides qualitative information about the spin dependent potential. The mass splittings between ψ and η_c and between ψ' and η'_c are well accounted for by the short-distance Coulombic $S_1 \cdot S_2$ term in the Breit potential, so this term is unlikely to have a large long range component.[51] In the absence of any large long-range spin-dependent forces, the $^1P_1$ state is expected to be at approximately the center of gravity of the ^3P_J states or

$$\frac{1}{9}(m_0 + 3m_1 + 5m_2) = 3525 \text{ MeV} \tag{6.1}$$

If it is at this mass or lower, the best chance is the decay $\psi' \rightarrow \pi^0\ {}^1P_1$ which Rosner[49] has estimated at a branching ratio of order 10^{-3}. If the mass is too heavy for $\psi' \rightarrow \pi^0\ {}^1P_1$ to occur, then the charmonium ${}^1P_1$ state might never be observed.

VII. ξ(2220) as a Higgs Boson Candidate

The ξ(2220) is one of the most intriguing particles found in radiative ψ decay. I have discussed in a previous section the possibility that it could be a $\bar{q}qg_{TM}$ meikton or a $g_{TM}g_{TM}$ glueball. The most recent experimental statement on the ξ width is $\Gamma_\xi(TOT) < 40$ MeV at the 95% confidence level.[7] If $\Gamma_\xi(TOT)$ is in fact much less than this upper limit, then the glueball, meikton, or any other hadronic interpretation of ξ would be very unlikely. The possible extreme narrowness of ξ suggests that it could be a particle from the electroweak world, a Higgs boson.

The measured rate

$$B(\psi \to \gamma \xi) \cdot B(\xi \to K^+K^-) = (5.8 \pm 1.8 \pm 1.5) \cdot 10^{-5} \quad (7.1)$$

already rules out the possibility that ξ is the Higgs boson H of the standard model with a single complex scalar doublet, because we then would expect

$$B(\psi \to \gamma H) \cong 3 \cdot 10^{-5} \quad (7.2)$$

Since only isoscalars are copiously produced in $\psi \to \gamma X$, we can safely assume a rate equal to (7.1) for $\psi \to \gamma\xi \to \gamma K_s K_s$, consistent with what is observed. If ξ is a Higgs boson we also expect that $B(\xi \to \bar{K}^*K^*)$ is at least as large and perhaps a few times larger than $B(\xi \to \bar{K}K)$. So the discrepancy implied by (7.1) and (7.2) is at the level of an order of magnitude.

This has led to consideration of the possibility that ξ is a nonstandard Higgs boson. The simplest alternative is to consider models with two Higgs doublets.[55,56,57] This is not entirely artificial, since two doublet models are motivated theoretically by attempts to solve the strong CP violation problem and by supersymmetry. To avoid tree level flavor changing neutral currents, all fermions of a given charge must couple to one of the two Higgs doublets. The principal constraints imposed by the experimental data are then

1) the enhancement of $B(\psi \to \gamma\xi)$ relative to eq. (7.2).
2) the upper limit on $\xi \to \mu^+\mu^-$, currently $B(\psi \to \gamma\xi) \cdot B(\xi \to \mu^+\mu^-) < 7.3 \cdot 10^{-6}$ at the 90% confidence level.[7]
3) the failure to observe non-strange final states, such as the 90% confidence level upper limit[7] $B(\psi \to \gamma\xi) \cdot B(\xi \to \pi^+\pi^-) < 3 \cdot 10^{-5}$.
4) the upper limit[58] from the CLEO detector $B(\Upsilon \to \gamma\xi) \cdot B(\xi \to K^+K^-) < 2 \cdot 10^{-4}$.
5) upper limits on the strangeness-changing transistion $b \to s + \xi$.

The class of two doublet models which may satisfy these constraints have charge + 2/3 quarks coupled (with enhanced strength relative to the standard model) to the Higgs doublet containing the principal component of the ξ, while the charge − 1/3 quarks and charge − 1 leptons are coupled principally to the other Higgs doublet and very little to ξ. This construction trivially satisfies the experimental constraints 1), 2), and 4), but there is a potential problem with the prominence of $\xi \to K^+K^-$ and the experimental constraint 3). In this class of models the $\xi \to \bar{K}K$ decays must come principally from hadronization of the two gluon decay, $\xi \to gg$, which occurs via the enhanced c and t quark loop diagrams. The problem then is why the digluon would hadronize preferentially to the $\bar{K}K$ final state.

This problem is clearly related to the question of flavor symmetry in glueball decays that was discussed in Section III, and there is a similar possible solution. In the J = 0 channel, the $gg \to \bar{q}q$ amplitude of figure (1a) is proportional to m_q and will favor $\bar{s}s$ over $\bar{u}u + \bar{d}d$. There is no such enhancement for figure (1b), but to the extent that intermediate $g_{TM}g_{TM}$ glueball states are dominant, both figures (1a) and (1b) interpreted as cavity perturbation theory diagrams imply a predominance of $\bar{s}s$ pairs in the final state.

As the authors of ref. (57) correctly remark, there is no evidence for such an enhancement in the decays of the J = 0 $\bar{c}c$ state, $\chi(3410)$, which also proceeds in perturbation theory via two gluons. However this does not necessarily negate the previous argument, for two reasons. First, counting rule arguments imply the dominance of fig. (1b) for $gg \to \bar{K}K$ at a large mass like 3400 MeV but not for smaller masses where fig. (1) could be more important. Second, intermediate glueball states are less likely to dominate at 3400 MeV.

The conclusion is that the $\pi\pi/\bar{K}K$ ratio involves such difficult dynamical issues that it cannot be regarded as a decisive test of the model. Therefore, although it is very important to search for $\pi\pi$, $\bar{\mu}\mu$, and other decay modes, the results of those searches will not yield a definitive test of the Higgs hypothesis.

There are at least three tests which could be decisive. If the spin is measured to be greater than zero or if a hadron-scale width is established, the Higgs hypothesis would be excluded. The third test is to see whether the ξ coupling is proportional to the quark mass by observing the radiative decay of toponium. Depending on the accomplishments of the Mark III collaboration in the intervening years, the first two tests could remain as important tasks for BEPC to perform.

VIII. Conclusion

The purpose of this talk was to discuss those aspects of J/ψ physics which are of great interest today and will continue to be of great interest after BEPC begins operation. The color force carried by the gluon is the central feature of the strong interaction. The most direct manifestations of the gluon are the particles – glueballs and perhaps meiktons – which have gluon constituents. These particles have no counterparts in electrodynamics; they can exist only because of the unique properties of the color force. We will not understand the strong interaction until we have found these particles and studied their properties. Because of the difficulty of the related theoretical and experimental problems, most, if not all, of this work will remain to be done after 1988. BEPC will be a unique world facility for these studies.

The difficulty is due in large measure to the great complexity of the spectrum in the 1 to $2\frac{1}{2}$ GeV region where the gluonic states are likely to occur. There is already evidence for 26 $\bar{q}q$ meson nonets which conform to the classification of the nonrelativistic quark model. There may in addition be cavity/string excitations of $\bar{q}q$ states and possibly cryptoexotic $\bar{q}\bar{q}qq$ states. These particles may overlap in mass and in some cases they will mix. The result is that the gluonic states cannot be treated in isolation. Rather it is necessary to understand the spectrum as a complete entity if we want to understand its components. An example is the discussion in Section III of the pseudoscalar glueball candidate ι(1440), which turns on understanding the π′ nonet of radially excited pseudoscalar $\bar{q}q$ mesons.

The problems are formidable but not insoluble. Not all states will be strongly mixed. There are characteristic features which may help to identify the different kinds of particles:

1) Glueballs should be the most prominent particles in radiative J/ψ decay. Their production rates in J/ψ decay, γγ scattering and hadron scattering will not be consistent with assignments in $\bar{q}q$ nonets. This requires a thorough understanding of the relevant qq nonets, obtained from many different kinds of experiments. The lightest glueballs are expected in $J^{PC} = 0^{++}, 0^{-+}, 2^{++}$.
2) The $\bar{q}qg$ meikton states may be distinguished by "signature" decay modes to multi-Kaon final states which would be OIZ suppressed decays of $\bar{q}q$ mesons. The lightest nonets are $J^{PC} = 1^{\pm -}, (0,1,2)^{\mp +}$.
3) Most cryptoexotic $\bar{q}\bar{q}qq$ configurations are unconfined and unobservably broad. The few which lie below their principal fall-apart threshholds may be distinguished by unusually large branching ratios to two body final states. The s-wave states occur in $J^{PC} = 0^{++}$, $1^{+\pm}, 2^{++}$.

The technical and scientific challenge of this program is considerable. In theoretical physics we must increase the power of the lattice computations and perhaps find new analytical methods. At BEPC the challenge of J/ψ physics is to cross a new frontier in statistics, to produce, observe, and analyze the unprecedented large number of events needed to understand the complex particle spectrum between 1 and $2\frac{1}{2}$ GeV.

This is first of all a challenge to approach the design specifications for luminosity and beam spread at 1.55 GeV per beam. The J/ψ production rate at BEPC will then surpass all previous storage rings by an order of magnitude.

The second challenge is the detection and analysis of an unprecedented number of events. My non-expert understanding is that this event rate would not overload the detector but would overload the proposed off-line computing facility which has the capability to analyze 5,000,000 J/ψ events per year. To extract the physics from the approximately 8% of radiative J/ψ decays, it would be necessary to devise a hard photon trigger and/or an efficient off-line pre-screening and/or to increase considerably the proposed computing capability. Further attention must be given to the computing requirements of a multi-amplitude partial wave analysis program with $\sim$ 100,000 events per channel as discussed in Section I.

With this level of statistics we will have a data sample that is appropriate to the difficulty of the problem. The largest available collection of analyzed J/ψ decays – 2.7 million from the Mark III – has succeeded primarily in teasing us with a richness of physics we are still unable to understand. It is clear that ten or perhaps twenty million decays are needed to answer the most straightforward questions raised by the present sample, such as the spin of $\xi(2220)$. Beyond that, experience in meson partial wave analysis suggests that we will require of order one hundred million J/ψ decays (or about eight million radiative decays) to resolve the complexity of structure up to 2 GeV.

It remains to be seen whether the Mark III will achieve the ten to twenty million event level, but there is no prospect in sight to obtain even higher statistics at SPEAR. If BEPC approaches its design specifications at the J/ψ energy, it will produce fifty to one hundred million per year. I can hardly wait to see the results from that data.

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Table 1. The meikton spectrum for Ref. 16 for C_{TE}/C_{TM} = 1/2, 1,2. Particles are labeled by analogy with the vector mesons. All masses are in GeV and all radii in GeV^{-1}.

J^{PC}	TYPE	$C_{TE}/C_{TM} = 1/2$		$C_{TE}/C_{TM} = 1$		$C_{TE}/C_{TM} = 2$	
		Mass	Radius	Mass	Radius	Mass	Radius
1^{--}	ρ/ω	1.64	6.10	1.83	6.35	2.02	6.56
	K^*	1.80	6.03	1.99	6.29	2.18	6.50
	ϕ	1.96	5.95	2.16	6.22	2.35	6.44
0^{-+}	ρ/ω	1.20	5.50	1.41	5.81	1.61	6.05
	K^*	1.41	5.42	1.62	5.74	1.82	5.98
	ϕ	1.61	5.34	1.82	5.67	2.03	5.91
1^{-+}	ρ/ω	1.41	5.80	1.61	6.05	1.80	6.31
	K^*	1.59	5.73	1.80	5.98	1.99	6.25
	ϕ	1.78	5.66	1.99	5.90	2.18	6.18
2^{-+}	ρ/ω	1.79	6.30	1.97	6.51	2.15	6.70
	K^*	1.94	6.24	2.13	6.45	2.13	6.65
	ϕ	2.09	6.17	2.28	6.39	2.47	6.59

Table 2. Predicted glueball masses from Ref. 16, for gluon self energy ratios C_{TE}/C_{TM} = 1/2, 1, 2 and for two different fits (I and II) to the mesons and baryons. Masses are in GeV. The 1.44 mass is an input parameter.

FIT	C_{TE}/C_{TM}	0^{++}	2^{++}	0^{-+}	2^{-+}
	1/2	0.67	1.75		
I	1	1.14	2.12	<u>1.44</u>	2.30
	2	1.56	2.47		
	1/2	0.65	1.74		
II	1	1.21	2.18	<u>1.44</u>	2.30
	2	1.70	2.59		

Table 3. Masses of TM^2 glueballs and $q_s q_s$ TM meiktons at $O(\alpha_s)$ using fit I of Reference 16. All masses are in GeV. The radii of the states are ~5-6 GeV^{-1}.

State		$C_{TE}/C_{TM} = 1/2$ ($C_{TM} = 2.16$)	$C_{TE}/C_{TM} = 1$ ($C_{TM} = 1.62$)	$C_{TE}/C_{TM} = 2$ ($C_{TM} = 1.08$)
TM^2	0^{++}	1.93	1.55	1.13
	2^{++}	2.64	2.30	1.94
1^{+-}	ρ/ω	2.13	1.95	1.76
	K^*	2.26	2.08	1.89
	ϕ	2.40	2.21	2.02
0^{++}	ρ	1.80	1.61	1.41
	ω	1.90	1.71	1.51
	K^*	1.98	1.79	1.59
	ϕ	2.20	2.01	1.81
1^{++}	ρ	1.94	1.76	1.56
	ω	2.04	1.86	1.67
	K^*	2.11	1.92	1.72
	ϕ	2.31	2.12	1.93
2^{++}	ρ	2.23	2.05	1.87
	ω	2.32	2.14	1.96
	K^*	2.35	2.17	1.99
	ϕ	2.51	2.33	2.15

Table 4. "Signature" decays of the $\bar{q}_s q_s$ TM meiktons into two L = 0 mesons in a relative s-wave, as expected from the decay mechanism discussed in the text.

	1^{+-}	0^{++}	1^{++}	2^{++}
"ρ"	$\phi\pi, K^*\bar{K}^*,$ $K\bar{K}^*, K^*\bar{K}$	$\phi\rho, K\bar{K}, K^*\bar{K}^*$	$\phi\rho, K\bar{K}^*, K^*\bar{K}$	$\phi\rho, K^*\bar{K}^*$
"ω"	$\phi\eta, \phi\eta', K^*\bar{K}^*,$ $K\bar{K}^*, K^*\bar{K}$	$\phi\omega, K\bar{K}, K^*\bar{K}^*$	$\phi\omega, K\bar{K}^*, K^*\bar{K}$	$\phi\omega, K^*\bar{K}^*$
"K^*"	$\phi K, \phi K^*$	ϕK^*	$\phi K, \phi K^*$	ϕK^*
"ϕ"	$\phi\eta, \phi\eta'$	$\phi\phi, \phi\omega^{\S}$	$\phi\omega^{\S}$	$\phi\phi, \phi\omega^{\S}$

§These decays may be suppressed relative to the others in the table since they involve the TM gluon coupling to $\bar{u}u$ and $\bar{d}d$, but they are included because they are not OZI suppressed for meikton decays unlike the corresponding decays of their ordinary meson counterparts.

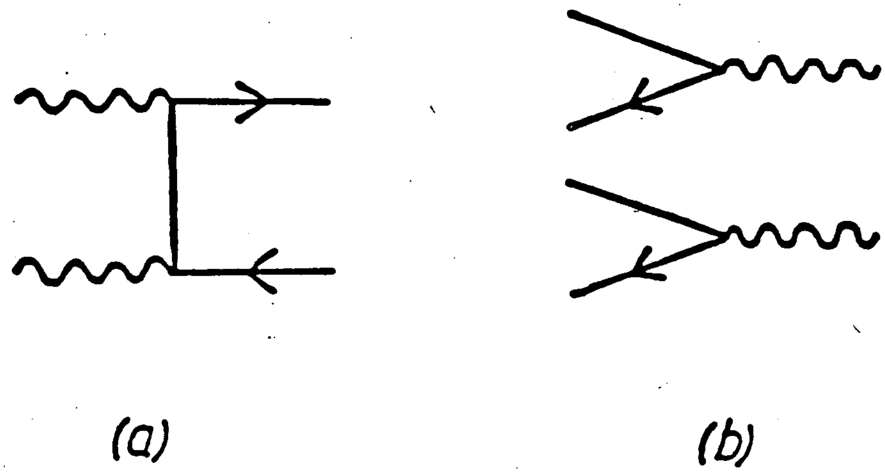

Figure 1

Lowest order glueball decay mechanisms.

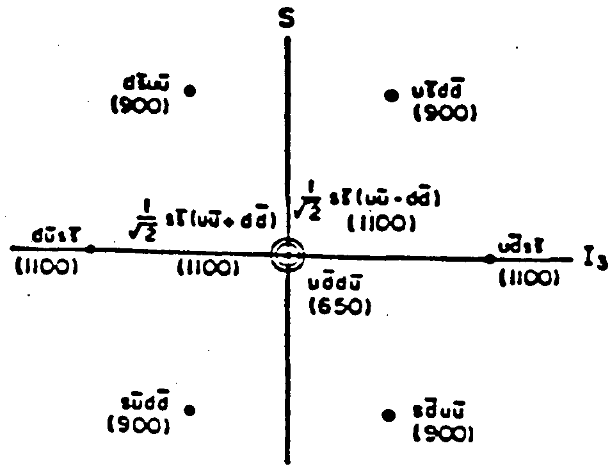

Figure 2, from Ref. (42). The lightest qqqq exotics: the $J^{PC} = 0^{++}$ crypto-exotic nonet. The quark content and masses are shown.

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