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A REVIEW OF RADIATIVE TRANSITIONS AND DECAYS OF $\psi$ and $\psi'$

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October 6, 1976

ABSTRACT

This paper is a somewhat expanded and up-dated version of a "mini-rapporteur" talk presented at the Tbilisi conference. Theoretical ideas about the new particle spectrum and dynamics are compared with the experimental data for the radiative transitions and decays of $\psi'$ and $\psi$.

I. INTRODUCTION

The radiative transitions of $\psi$ and $\psi'$ provide us with the best window on the spectroscopy and dynamics of even charge conjugation states of hidden charm. The data presented to the conference this morning strongly suggest that the observed spectrum is just the spectrum predicted for quark-antiquark bound states. This illustrates in a dramatic way the value of the new particles in hadronic physics: while we still worry about the existence of the $A_1$ meson it is very likely that its $J^{PC} = 1^{++}$ counterpart in the new spectroscopy has already been discovered (at $M = 3.50$ GeV as argued below). Further study of the angular distributions of radiative transitions will before long allow us to confirm the assignment of quantum numbers which is now strongly suggested. In a few years we may well have learned more about the spectrum of hadrons from the new particles than we have learned in as many decades from the old particles.

There are also reasons to hope that the new particles will be of great value in the effort to understand hadron dynamics. Here the question is whether the new hadrons--because of their greater mass--are in some ways simpler dynamical systems than the old hadrons. In this connection two independent conjectures have been expressed, both under the heading "charmonium." First it is suggested that the states of hidden charm may be simply described as nonrelativistic bound states of a heavy quark-antiquark pair in a potential dominated by a Coulomb component at short distances and a linearly rising component at large distances. The second suggestion is that the annihilation of the heavy quarks in charmonium decays may take place at short distances and therefore be calculated perturbatively in an asymptotically free theory. I will review the most general consequences of these two conjectures in the light of the data presented this morning.

In a very different theoretical approach to the new particles, dual model and Regge pole ideas have been used with qualitative success to describe the phenomenology of nonradiative O.I.Z. rule violating decays. In this approach the vector meson dominance (VMD) hypothesis is used to make contact with radiative transitions and decays. Because of the uncertain validity of the VMD hypothesis for the new particles, radiative phenomena are not the best testing ground for this approach. Consequently I will not say very much about the dual model approach or about other models (e.g., SU(8)$_c$) which rely on VMD. Instead I will limit myself to a brief discussion of the formidable complications which must be disentangled in order to test the $\psi$-VMD hypothesis.
Some of the effects I will mention will be discussed more fully by Professor Ioffe in his talk in this session.

The plan of the talk is as follows. First I will discuss the spectrum of even charge conjugation states and the most probable assignment of quantum numbers. Next I will compare with the data the conjectures that the new particles are bound states of heavy quarks in a potential and that their decays are controlled by the short distance behavior of the gauge theory of color SU(3). Finally I will comment on the problems which must be understood in order to test the validity of the VMD hypothesis for the $\psi$.

II. SPECTRUM OF EVEN CHARGE CONJUGATION STATES

As we heard this morning, $\psi'$ and $\psi$ reveal five even charge conjugation states. This is just what we would naively expect for a quark-antiquark spectrum, $\bar{\psi}c$, which would probably have below the $\psi'$ two pseudoscalars, $\eta_c$ and $\eta_c'$, and three $L=1$ states, $J^P = 0^+, 1^+, 2^+$ (the fourth $L=1$ state has $C=-$ and will be hard to detect). Two of the five $\bar{\psi}c$ states must have natural parity and indeed two of the five observed states are reported to decay into $\pi\pi$ and $KK$: the 3.41 and the 3.55.

The photon angular distribution (with respect to the beam) for $\psi' \rightarrow \gamma + 3.50$ is
\[
1 + (1.1 \pm 0.4) \cos^2 \theta,
\]
consistent with the required $1 + \cos^2 \theta$ if the 3.50 has spin zero. We therefore assign 3.41 to the scalar, $J^P = 0^+$, and the remaining natural parity state at 3.55 to the tensor, $J^P = 2^+$. The state observed$^3$ decaying into two photons, to which I nominally assign the mass 2.85 GeV, is presumably the $\eta_c$, $J^P = 0^-$, and the remaining problem is to match 3.45 and 3.50 to $1^+$ and $\eta_c'$.

The DESY-Heidelberg group has measured the photon angular distribution$^4$ in $\psi' \rightarrow \gamma + 3.50$ to be
\[
1 - (1.1 \pm 0.5) \cos^2 \theta,
\]
which is six standard deviations from the $1 + \cos^2 \theta$ distribution required if the 3.50 has spin zero. This implies the assignments 3.50 = $1^+$ and $3.45 = \eta_c'$. But the LBL-SLAC group reports$^1$ $1 + (0.26 \pm 0.5) \cos^2 \theta$ for the same distribution, which is only $1/2$ standard deviations from the distribution for spin zero. It is therefore significant that two additional arguments$^5$ suggest that the correct assignment is indeed 3.50 = $1^+$ and 3.45 = $\eta_c'$:

1. If we choose 3.50 = $\eta_c'$ and if-as is true in a variety of potential models$^6$-- $| W(\psi \rightarrow \eta_c') | > | W(\psi' \rightarrow \eta_c') |$
then for the branching ratios we have
\[
B(\psi \rightarrow \eta_c') \geq 3B(\psi' \rightarrow \eta_c') = (0.27 \pm 0.09),
\]
which is an order of magnitude above the experimental upper limits$^1, 2$ for $\psi \rightarrow \gamma \pi$.

2. If $3.45 = 1^+$ and the $\psi' \rightarrow \gamma \{\psi, \eta_c \}$ transitions are predominantly electric dipole, then $B(\psi' \rightarrow \gamma + 3.45) =$
\[
3 \left( \frac{220}{264} \right) \frac{3}{4} B(\psi' \rightarrow \gamma + 0^+) = (0.15 \pm 0.04), (\text{Ref. 1}),
\]
or $= (0.18 \pm 0.06), (\text{Ref. 2})$, which is again well above the experimental upper limit.$^2$

For all these reasons the most probable assignment is 3.50 = $1^+$ and $3.45 = \eta_c'$. This assignment is also supported by the apparent similarities of 2.85 and 3.45: nearly equal hyperfine splitting from $\psi$ and $\psi'$ and formation rates from $\psi$ and $\psi'$ decays which are significantly smaller than the rates for the $\psi'$ transitions to the other three $C$-even states.

It is important to keep in mind that we cannot now exclude the possibility that some of the five observed states do not correspond to
the five expected \( \bar{c}c \) found states. For instance, Okun and Voloshin consider the possible existence of hadronic "molecules," \( \bar{D}D \) systems bound by exchange of \( \pi, \rho, \) and \( \omega \). Another possibility within the charm framework is bound states of a diquark, \( \bar{c}q \), with an antidiquark, \( \bar{q}c \).

Fortunately we do not have to rely on guesswork to establish the proper assignment. By measuring angular distributions in the decay sequence \( \psi' \rightarrow \gamma X \rightarrow \gamma \gamma \gamma \rightarrow \gamma \mu^+\mu^- \) the spin of the \( X \) states can be determined.\(^9,10,11\) To determine the parity of the \( X \) states it is necessary to measure photon polarizations in the \( \psi' \rightarrow \gamma X \rightarrow \gamma \gamma \gamma \) events or to study particular hadron final states in the sequence \( \psi' \rightarrow \gamma X \rightarrow \gamma + \text{hadrons} \).

For the remainder of this talk I assume the validity of the assignment given above. In particular I am taking completely seriously the existence of the states at 3.45 and 2.85, which are not as well confirmed as the other three. If either of these states is not confirmed or if they are more "exotic" than \( \bar{c}c \) bound states,\(^7,8\) then much of what I have just said about the spectrum and will now say about dynamics may be called into question.

III. DYNAMICS

The central question is whether the dynamics of the new physics will be more simple than the old. In particular, it is conjectured that the new states may be simply described as bound states of heavy quarks in a potential and that certain decay rates may be governed by calculable short distance physics. I will try to compare the most model independent features of these conjectures with the emerging experimental situation. I will also very briefly discuss \( \psi \) vector meson dominance in connection with models which use \( \psi \)-VMD to make contact with radiative phenomena.

A. Heavy Quarks in a Potential

Since the new quarks are heavy it is natural to conjecture that the \( \psi, \chi_c \), and \( \eta_c \) states may be treated for some purposes in nonrelativistic approximation and to consider the possibility of a description in terms of an effective potential. For heavy quarks in a potential the ratio of \( M_2 \) to \( E_1 \) transitions is \( \sim 0(k_\gamma^2/m_Q^2) << 1 \) and a more careful estimate\(^12\) using charmonium parameters gives 0.005. \( E_1 \) dominance is best tested by measuring angular distributions,\(^9-12\) e.g., for \( \psi' \rightarrow \gamma J^+ \) and \( \psi' \rightarrow \gamma 2^+ \) the photon should be distributed with respect to the beam direction like \( 1 - \frac{1}{3} \cos^2 \theta \) and 

\[ 1 + \frac{1}{3} \cos^2 \theta, \]

respectively. These predictions are consistent with one set of data\(^7\) but differ by two standard deviations from another measurement\(^4\) for the 3.50. In more sophisticated models, not all transitions are expected to be \( E_1 \) dominated. For instance, in an approach\(^13\) combining current algebra and \( \psi \) vector dominance, the transitions \( \psi' \rightarrow \gamma 2^+ \) and \( \psi' \rightarrow \gamma J^+ \) are predicted to have comparable contributions from \( M_2 \) and \( E_1 \) moments. (In this approach it is also predicted that \( \psi' \rightarrow \gamma J^+ \rightarrow \gamma \gamma \gamma \) should be dominated by the \( 2^+ \) state; experimentally the 3.50 dominates, which we have taken to be the \( 1^+ \) state.)

For electric dipole transitions we would naively expect formation rates proportional to \( \Gamma(\psi' \rightarrow \gamma J^+) \propto (2J + 1)k_\gamma^3 \), i.e.,

\[ 0^+ : 1^+ : 2^+ = 1.6 : 1.5 : 1.0, \]

to be compared to experimental branching ratios announced this morning of \( 0.08 \pm 0.03 : 0.09 \pm 0.03 : 0.08 \pm 0.03 \). These measured ratios agree with the \( E_1 \) ratios at the
one standard deviation level. But such ratios are not as general a
test of $E_1$ dominance as the angular distributions, because the wave
functions of the $L = 1$ states may be distorted by mixing. A more
sophisticated treatment which attempts to take $D\bar{D}$ mixing into
account gives $2.3 : 1.2 : 1$.

The hypothesis that strong interactions are mediated by the
massless gluons of the color SU(3) gauge theory suggests that the
effective nonrelativistic potential is coulombic at short distances
and linearly growing at large distances. For a vector potential with
the nonrelativistic limit $V = \lambda r + \frac{\beta}{r} + \gamma$ ($\lambda$, $\beta$, and $\gamma$ are
constants) the $L = 1$ mass splitting computed in Born approximation
must satisfy $0.8 \leq (M_2 - M_1)/(M_1 - M_0) \leq 1.4$ while the above
assignment of states gives $\sim 0.5$. However a solution of the
Salpeter equation with relativistic kinematics using a combination of
Coulomb potential plus a scalar linear potential gives values between
0.3 and 0.5 (to obtain values near 0.5 a momentum dependent coupling
constant as in asymptotic freedom is assumed). It is important to
understand the difference in the results of Refs. 15 and 16. Are
relativistic effects important? Is the Born approximation not
adequate? Or is the difference due chiefly to the Lorentz character
assumed for the linear potential?

The magnitudes of $E_1$ and $M_1$ transition rates are evidently
rather model dependent. In some potential models the estimated
rates for $\psi' \rightarrow \gamma P_{\gamma}$ are within a factor of two of the experimental
values reported this morning; other estimates are in more serious
disagreement with the data. The suppressed $M_1$ decays, $\psi' \rightarrow \eta_c \gamma$ and $\eta_c \rightarrow \psi \gamma$, may be the most model dependent of all:

relativistic effects are so important that the ratio

$$\Gamma(\psi' \rightarrow \eta_c \gamma)/\Gamma(\eta_c \rightarrow \psi \gamma)$$

may vary by four orders of magnitude depending on the Lorentz transformation properties of the binding potential.

This is a reminder of how difficult life can be when relativistic
effects are important, as they must surely be in the light bound states
of the "old physics."

(Since the conference a paper by Jackson has appeared
containing useful tests of the nonrelativistic potential hypothesis.
He observes that in the Russell-Saunders limit $F_1$ transition rates
must satisfy the Thomas-Reiche-Kuhn sum rule of atomic physics. This
implies upper bounds for $\Gamma(\psi' \rightarrow \gamma \eta)$ which depend on the quark mass
and charge but are independent of the details of the binding potential.
Lower bounds can be obtained from a second atomic physics sum rule, at
the cost of more restrictive assumptions about the nature of the
potential.)

B. Decay Rates as Short Distance Physics

In the color SU(3) gauge theory (QCD : quantum chromo-
dynamics) the creation of a heavy quark pair occurs at short distances
for which quark-gluon interactions are thought to be weak because of
asymptotic freedom. Thus Appelquist and Politzer conjectured (and,
to their credit, before November of 1974) that just below the charm
threshold there might be a narrow $c\bar{c}$ state, orthocharmonium,
alogous to the orthopositronium states of quantum electrodynamics.
In this view the small hadronic width of the $\psi$ (and the "Okubo-
izuka-Zweig rule") is explained by the fact that the heavy quarks
must annihilate at short distances, on the order of their Compton
wavelength, so that the interaction is weak by asymptotic freedom and
the total hadronic decay width is dominated by the lowest order
Feynman diagrams.
It is important to add that the theoretical basis for this approach is rather fuzzy: the argument is not much more than the heuristic remarks of the previous paragraph. It is certainly not as well founded as the predictions of asymptotic freedom for deep-inelastic lepton-nucleon scattering. The most complete discussions can be found in the work of Appelquist and Politzer,\textsuperscript{21} which the interested reader should consult.

In their analysis two levels of prediction emerge. On the deeper level, one calculates the value of the charmonium wave function at the origin and then predicts the magnitudes of various decay rates. Simple estimates\textsuperscript{21} show that these calculations cannot be quantitatively reliable until we have produced quarks with masses which are even beyond PEPI-PEP2 energies. (The necessary condition is 
\[ \alpha_s \ln \frac{2\alpha_s}{3} \ll 1 \] where \( \alpha_s \) is the effective quark-gluon fine structure constant. For instance, if \( \alpha_s (E_{CM} = 3 \text{ GeV}) \approx 0.4 \), then
\[ \alpha_s \ln \frac{2\alpha_s}{3} \bigg|_{E_{CM} = 30 \text{ GeV}} \approx -0.4 \]

Less ambitiously, for charmonium we must settle for calculating ratios of decay rates which do not depend on the wave function at the origin. I will discuss the predictions for the following ratios:

(a) \( \frac{\Gamma(\psi \to \text{hadrons})_{\text{direct}}}{\Gamma(\psi \to e^+e^-)} \)
(b) \( \frac{\Gamma(\eta_c \to \text{hadrons})}{\Gamma(\eta_c \to \gamma\gamma)} \)
(c) \( \frac{\Gamma(\eta_c \to \gamma\gamma)}{\Gamma(\psi \to e^+e^-)} \)
(d) \( \frac{\Gamma(\psi \to \gamma + \text{hadrons})}{\Gamma(\psi \to \text{hadrons})_{\text{direct}}} \)
(e) \( \frac{\Gamma(0^+ \to \text{hadrons})}{\Gamma(1^+ \to \text{hadrons})} : \frac{\Gamma(2^+ \to \text{hadrons})}{\Gamma(0^+ \to \text{hadrons})} \)

In the lowest order diagram for \( \psi \to \text{hadrons} \), the \( \bar{c}c \) pair annihilates into three "real" on-mass-shell color gluons (one and two gluons are forbidden by color SU(3) symmetry and charge conjugation). As in the parton model, we ignore the "final state interaction" which occurs at large distances (with strong coupling) and converts the "confined" gluons to hadrons. The total rate for \( \psi \to \text{hadrons} \) is computed by squaring the amplitude and computing the three gluon phase space. For \( \psi \to e^+e^- \) we have \( \bar{c}c \) annihilation via a virtual photon. In the ratio of the two rates the \( \bar{c}c \) wave function cancels and we have\textsuperscript{21,22}

\[ \frac{\Gamma(\psi \to \text{hadrons})_{\text{direct}}}{\Gamma(\psi \to e^+e^-)} = \frac{5 (m^2 - 9) \alpha_s^3}{16\pi \alpha^2}. \] (1)

The numerator is the direct hadronic decay width, i.e., with the contribution (~12 keV) from \( \psi \to \gamma + \text{hadrons} \) subtracted. Experimentally the left-hand side is ~47/4.8 ~ 10 and we use Eq. (1) to determine \( \alpha_s \) at \( q^2 = m^2_\psi \), obtaining

\[ \alpha_s (m^2_\psi) \approx 0.19. \] (2)

This small value for \( \alpha_s \) is already reason for some nervousness. Naively, to fit \( R \approx 2.5 \) for \( 2.4 < \sqrt{s} < 3.4 \text{ GeV} \) we need \( \alpha_s \approx 0.4 \), which is also the kind of value that typically emerges from more careful analyses\textsuperscript{23} of \( \sigma(e^+e^- \to \text{hadrons}) \). But the theoretical and experimental\textsuperscript{24} uncertainties are large enough that we can provisionally regard \( \alpha_s \approx 0.2 \) as a possibility. For instance, the SPEAR measurements of \( R \) contain a ~10\% normalization uncertainty and a possible smooth variation of ~15%; these errors are not shown in the error bars in the usual displays of \( R \).\textsuperscript{24} For consistency we must of
course use this small value of $\alpha_s$, Eq. (2), in predicting the ratios (b), (d), and (e).

(b) The calculation of $\Gamma(\eta_c \rightarrow \text{hadrons})/\Gamma(\eta_c \rightarrow \gamma\gamma)$ proceeds just as in (a). We compare the contribution of the two photon amplitude with the two photon amplitude obtaining

$$\frac{\Gamma(\eta_c \rightarrow \text{hadrons})}{\Gamma(\eta_c \rightarrow \gamma\gamma)} = \frac{9}{6} \left( \frac{\alpha_s}{\alpha} \right)^2 \approx 760. \quad (3)$$

According to the DASP data presented this morning,3 the branching ratio $B(\psi \rightarrow \gamma \eta_c \rightarrow \gamma\gamma)$ is $\approx (1.6 \pm 0.5) \times 10^{-4}$. For $B(\psi \rightarrow \gamma \eta_c)$ upper bounds have been reported25 of $\lesssim 0.039$ and $\lesssim 0.03$. Taking $B(\psi \rightarrow \gamma \eta_c) \lesssim 0.03$ we have $\Gamma(\eta_c \rightarrow \text{hadrons})/\Gamma(\eta_c \rightarrow \gamma\gamma) \lesssim 200 \pm 60$, a serious disagreement with the prediction, Eq. (3). If this result and the assignment $\chi(2s) = \eta_c$ are verified by future experiments, then it will be clear that the short distance hypothesis cannot be applied to charmonium dynamics on a quantitative level. But if the upper end of the inequality is saturated so that $\Gamma(\eta_c \rightarrow \gamma\gamma)$ is not much more than a factor $\sim 3$, then the qualitative connection between short distance physics and the narrow width of the $\psi$ might still turn out to be correct.

(c) The prediction21,22

$$\Gamma(\eta_c \rightarrow \gamma\gamma) = \frac{4}{3} \Gamma(\psi \rightarrow e^+e^-) \quad (4)$$

is in one way more general than Eqs. (1) and (3) since it does not depend on the strong interactions of the OZI rule. But it is more model dependent in another way, since it assumes that the value of the wave function at the origin is equal for both $\psi$ and $\eta_c$. Given the unexpectedly large hyperfine splitting, $m_\psi - m_{\eta_c} \sim 250 \pm 50$ MeV, we should probably be wary of this assumption.

Combining Eqs. (2), (3), and (4) we can compute the widths of the direct hadronic decays of $\eta_c$ and $\eta_c'$:

$$\Gamma(\eta_c \rightarrow \text{hadrons}) \approx 10^3 \Gamma(\psi \rightarrow e^+e^-) \sim 5 \text{ MeV}$$

$$\Gamma(\eta_c' \rightarrow \text{hadrons}) \approx (900) \Gamma(\psi \rightarrow e^+e^-) \sim 2 \text{ MeV}. \quad (5)$$

With the assignment $\eta_c' = \chi_{3s1\zeta}$, Eq. (5) is probably too large by at least one order of magnitude.5 Experimentally we have $B(\psi' \rightarrow \gamma \eta_c' \rightarrow \gamma\gamma) = 0.008 \pm 0.004$ and conservatively taking the absence of a signal for $\psi' \rightarrow \gamma \eta_c'$ to mean that $B(\psi' \rightarrow \gamma \eta_c') \lesssim 0.06$, we have

$$\Gamma(\eta_c' \rightarrow \text{all}) \lesssim (8 \pm 4) \Gamma(\eta_c' \rightarrow \gamma\gamma). \quad (6)$$

It is very difficult to give a reliable estimate for the right-hand side of (6). First, as discussed in Section IIIA, the suppressed $M_{\perp}$ decays are extremely model dependent6 since relativistic effects are numerically of the leading order. Second, with the $\eta_c$ and $\eta_c'$ masses which we are assuming, the hyperfine splitting is very large; thus spin dependent forces are very important in the charmonium system and may cause the $^1S_0$ wave functions to be very different from the $^3S_1$. Estimates of $\Gamma(\eta_c' \rightarrow \gamma\psi)$ (and of $\Gamma(\psi' \rightarrow \eta_c \gamma)$) are therefore doubly suspect. However even a very conservative bound $\Gamma(\eta_c' \rightarrow \gamma\psi) \lesssim 20$ keV (which overestimates the model calculations of Ref. 6 by a factor $\sim 20$) implies that $\Gamma(\eta_c' \rightarrow \text{all})$ is at least an order of magnitude smaller than predicted in Eq. (5). We should however keep in mind that a large part of this discrepancy could be the effect of the hyperfine splitting rather than of the short distance physics: in this case the wave function at the origin must
be smaller for \( \eta_c \) than for \( \psi \). (Note that the more severe bound on \( \Gamma(\eta'_c \to \text{all}) \) obtained in Ref. 5 by comparing \( \eta'_c \to \gamma \psi \) and \( \psi' \to \gamma \eta_c \) may be unreliable, as discussed in Section IIIA and in Footnote 19.)

The experimental disagreement with Eq. (5) is held to an order of magnitude because of the generous estimate \( \Gamma(\eta'_c \to \gamma \psi) \approx 20 \text{ keV} \). The disagreement increases as this bound is decreased. Large spin-spin forces could save the day (qualitatively if not quantitatively) if they were to increase \( \Gamma(\eta'_c \to \gamma \psi) \) while at the same time causing \( \eta_c \) and \( \eta_c' \) to have smaller wave functions at the origin than \( \psi \) and \( \psi' \).

To compute the rate for \( \psi \to \gamma + \text{hadrons} \), one of the three gluons in (a) is replaced by a real photon. The result\(^{26} \)

\[
\frac{\Gamma(\psi \to \gamma + \text{hadrons})}{\Gamma(\psi \to \text{hadrons})_{\text{direct}}} = \frac{16 \alpha}{5 \alpha_s} \approx 0.12 \tag{7}
\]

which would be a sizeable signal. Furthermore the single photon energy spectrum should be linearly rising.\(^{27} \) These predictions have not yet been tested since it is much more difficult to measure the single photon continuum than a monochromatic photon signal.

The \( 2^+ \) and \( 0^+ \) states decay into hadrons via two massless gluons with the predicted ratio\(^{28} \)

\[
\frac{\Gamma(2^+ \to \text{hadrons})_{\text{direct}}}{\Gamma(0^+ \to \text{hadrons})_{\text{direct}}} = \frac{4}{15}. \tag{8}
\]

The \( 1^+ \) state is forbidden to decay into two massless gluons;\(^{29} \) it may decay into three massless gluons or into one massless gluon and a quark-antiquark pair. The gluon-quark-antiquark decay dominates because of a logarithmic dependence on the binding energy, with the result\(^{30} \)

\[
\frac{\Gamma(2^+ \to \text{hadrons})_{\text{direct}}}{\Gamma(1^+ \to \text{hadrons})_{\text{direct}}} = \frac{3\pi}{15} \frac{1}{\alpha_s} \frac{m_c}{\epsilon} \approx 4 \tag{9}
\]

if \( m_c/\epsilon \approx 10 \) (\( m_c \) is the \( c \) quark mass and \( \epsilon \) the binding energy of the \( 1^+ \) state).

In the approximation \( \Gamma(X \to \text{all}) \approx \Gamma(X \to \text{hadrons})_{\text{direct}} + \Gamma(X \to \gamma \psi) \) and assuming \( E_1 \) dominance we may use the identity\(^{5} \)

\[
\frac{\Gamma(X \to \text{hadrons})}{\Gamma(X \to \gamma \psi)} = \frac{\Gamma(X \to \gamma \psi)}{\Gamma(X \to \gamma \psi)} \frac{B(X \to \gamma \psi)}{B(X \to \gamma \psi)} - 1 \tag{10}
\]

to compare the predictions (8) and (9) with the available data. With the experimental values\(^{31} \) \( B(1^+ \to \gamma \psi) = 0.30 \pm 0.08 \) and \( B(2^+ \to \gamma \psi) = 0.125 \pm 0.075 \), we compute

\[
\frac{\Gamma(2^+ \to \text{hadrons})}{\Gamma(1^+ \to \text{hadrons})} \approx 4 \tag{11}
\]

which is the predicted result of Eq. (9). However the experimental errors are so large that for the moment it is premature to claim a real quantitative success. It is even more premature to attempt to compare Eq. (8) with the data, since at SLAC only one and at DESY only two candidates for \( 0^+ \to \gamma \psi \) have been observed. Equation (8) is on more solid theoretical ground than Eq. (9) since it does not depend on the poorly known parameter \( m_c/\epsilon \) and since the corrections from non-leading diagrams are probably smaller for Eq. (8) than for Eq. (9).

Although it is premature to claim the agreement of Eqs. (9) and (11) as a quantitative victory, it is certainly not premature to
note the qualitative success of this short distance prediction. If the radiative widths for $\psi' \rightarrow \gamma \chi'_j$ and $\chi'_j \rightarrow \gamma \psi$ are comparable (as both simple $E_1$ dominance and the available data suggest they are) then the short distance argument predicts that the decay sequence $\psi' \rightarrow \gamma \chi'_j \rightarrow \gamma \bar{\psi}$ will be strongest for the $1^+ \chi'_j$ state, as the data clearly indicates. This qualitative success is especially striking because the prediction is a very unique consequence of the short distance Q.C.D. hypothesis. It is the one indication in the available data that the short distance hypothesis may be correct on at least a semiquantitative level.

C. Other Approaches: Dual Models and $\psi$ Vector Meson Dominance

The experimental data as interpreted in Sections II and IIIB suggests that the hadronic widths of $\eta_c$ and $\eta'_c$ are not an order of magnitude larger than the widths of $\psi$ and $\psi'$. This is probably bad news for some $\eta - \eta' - \eta_c$ mixing models and for some WMD (vector meson dominance) models. Models in which the old quark content of $\eta_c$ and $\eta'_c$ approach the one percent level would require hadronic decay widths on the order of a few MeV or more. However dual models with mass-dependent mixing angles could accommodate appreciable $\bar{c}c$ mixing in $\eta$ or $\eta'$ yet have a much smaller light quark mixture in $\eta_c$ and $\eta'_c$. WMD models for $\psi \rightarrow \eta \gamma$ and $\psi \rightarrow \eta'_c \gamma$ which assume that the photon is dominated by a $\bar{c}c$ vector meson such as $\psi$ or $\psi'$ are also called into question, since part of the rationale for these models is the hypothesis that the charge-conjugation-even $\bar{c}c$ system should have more copious hadronic decays than the charge-conjugation-odd system (the present experimental value for $\Gamma(\psi \rightarrow \eta'_c \gamma)/\Gamma(\psi \rightarrow \eta \gamma)$ is smaller than these models predict). The WMD models also have other uncertainties which I will discuss below.

Many models use $\psi$ to make contact with radiative decays and transitions: dual models, current algebra models, and pure WMD models. In particular the dual models have developed a systematic semiquantitative treatment of OIZ rule phenomenology which is in reasonable agreement with the data. Since these models are purely hadronic in nature they do not confront radiative phenomena directly but only through the expedient of vector meson dominance. Because of the uncertain validity of the WMD hypothesis for the new particles, the predictions for radiative phenomena are not the best testing ground for these models. Therefore I will confine my remarks to the possible dangers of $\psi$-WMD and to the formidable difficulties involved in testing the $\psi$-WMD hypothesis experimentally. My remarks will be brief; Professor Ioffe will discuss some of these issues more extensively in his talk in this session.

The most obvious danger of $\psi$-WMD is the uncertainty of the tremendous extrapolation from $k^2 = 0$ to $k^2 = m^2_{\psi} \sim 10 \text{ GeV}^2$. Rho dominance, which requires an extrapolation from $k^2 = 0$ to $k^2 = 0.6 \text{ GeV}^2$, works extremely well. But $\phi$-dominance seems to fail by a factor of two, and for the $\phi$ we are only extrapolating from $k^2 = 0$ to $k^2 = 1 \text{ GeV}^2$. In the absence of any theoretical insight into the reliability of the extrapolation, we cannot be surprised if we fail to obtain even the correct order of magnitude when we extrapolate to $10 \text{ GeV}^2$.

Furthermore it is possible that $\psi'$, $\psi''$, ... may introduce much larger errors into pure $\psi$ than do $\phi'$, $\phi''$, ... into pure
\[
\frac{r_p}{r_p'} = \frac{0.15}{0.42} \approx 0.15 \\
\frac{r_q}{r_q'} \approx 0.42.
\]

The effect may even be larger than these numbers suggest, since interference effects of unknown phase are characterized by the square roots of the above ratios. In old physics applications, \( r' \) corrections may be much smaller than \( \frac{1}{2} \) because of the small coupling of the \( r' \) to hadrons, which might be related partially to the small ratio \( \left( \frac{m_p}{m_{p'}} \right)^2 \approx 0.2 \). But for the new physics \( \left( \frac{m_q}{m_q'} \right)^2 \) is not a small number. Interference between \( q \) and \( q' \) could give rise to 100% corrections of the naive \( q \)-VMD estimates.

Given these grounds for caution, it is important to subject the \( q \)-VMD hypothesis to experimental test. This is possible in principle by photoproducing \( q \) from hydrogen and heavy nuclei. Using \( q \)-VMD the elastic photoproduction of \( q \) from hydrogen implies a value for the \( qN \) total cross section, which can be compared to the value obtained from the \( A \) dependence of \( q \) production from heavy nuclei. But in practice it may be very difficult to give these experiments an unambiguous interpretation. I will just mention the effects which must be disentangled:

1. The \( qN \) forward amplitude may not yet be diffractive, as assumed in the usual VMD interpretation of \( \gamma N \to qN \).

Then \( qN \to qN \) might be dominated not by Pomeron exchange but by elementary particle poles: Ioffe suggests that the t-channel pole of the \( J^P = 1^{++} \) state might dominate at present energies.

2. Even if the \( qN \) amplitude is already dominated by Pomeron exchange, the Pomeron-\( q \) virtual vertex may be dynamically suppressed compared to the Pomeron-\( q \)-real vertex.

3. Interfering contributions from virtual \( q' \), \( q'' \), ... may be substantial.

4. The \( q\gamma \) coupling may be very different at \( q^2 = 0 \) and \( q^2 = m_q^2 \).

To study these possible effects, in addition to \( q \) and \( q' \) photoproduction, it is important also to study the photoproduction of charmed particles. If the usual \( q \)-VMD hypothesis is correct then together with the CIZ rule it implies copious charmed particle production in \( \gamma N \) scattering:

\[
\sigma(\gamma N \to \bar{D} + \cdots)/\sigma(\gamma N \to qN) \approx 30 - 50.
\]

But if the effects discussed above are operating then this prediction will be modified.

As an example, suppose that effects (1) and (4) are important but that (2) and (3) are not. Define an extrapolation factor \( \epsilon \) by

\[
\left( \frac{\epsilon}{\epsilon_q} \right)^2 = \frac{M(qN \to qN)_{\text{E}}}{M(qN \to qN)_{\text{E}}}.
\]

and let \( \beta_E \) be the ratio of the real to imaginary part of the forward \( qN \) amplitude at incident "lab" energy \( E \):

\[
\beta_E = \frac{\text{Re } M(qN \to qN)}{\text{Im } M(qN \to qN)}
\]

Then the true \( qN \) total cross section (determined from the \( A \) dependence of \( q \) photoproduction from nuclei) is related to the naive
ψ-VMD cross section (determined from ψ photoproduction off hydrogen) by the expression

\[ \sigma_{\text{TOT}}^{\psi N} = \frac{1}{\sqrt{1 + \beta_E^2}} \sigma_{\text{VMD}}^{\psi N} \quad (13) \]

Using also the OIZ rule we can relate the total photoproduction cross section for ψ and charmed particles to \( \sigma_{\text{VMD}}^{\psi N} \):

\[ \sigma(\gamma N \to \psi + \cdots) + \sigma(\gamma N \to \text{charm} + \cdots) \]

\[ = \sqrt{\frac{\epsilon}{1 + \beta_E^2}} \left( \frac{e}{T_{\psi}} \right)^2 \frac{1}{m_\psi^2} \sigma_{\text{TOT}}^{\psi N} \quad (14) \]

Thus with enough data we could determine \( \epsilon \) and \( \beta_E \) from (13) and (14).

Of course we must also contend with effects (2) and (3) which were neglected in the above example. Professor Ioffe will discuss effect (2) in his talk (see also Ref. 49). To my knowledge no one has looked carefully at the possible importance of \( \psi-\psi' \) interference mentioned in (3).

IV. CONCLUSION

The crucial experimental issue on which much of the preceding discussion depends is whether the existence of the states at 2.8 and 3.45 can be confirmed and their quantum numbers determined. If the 2.8 is confirmed then it is very likely to be the \( \eta_c \), and that would tend to reinforce the assignment of the 3.45 as the \( \eta_c' \), given the similarities of the 2.8 and 3.45. If the existence of the 2.8 is not confirmed, then the 3.45 might still exist but it is less likely to be the \( \eta_c' \) and would probably have a more exotic interpretation. In this case, the hyperfine splitting of \( \eta_c - \psi \) and \( \eta_c' - \psi' \) might be small, \( \leq 100 \text{ MeV} \), as naively expected, and the allowed radiative transitions \( \psi \to \gamma \eta_c \) and \( \psi' \to \gamma \eta_c' \) would be severely suppressed by the limited phase space. The transition \( \psi' \to \gamma \eta_c \) might then be one of the best ways to discover the \( \eta_c' \), though its rate is extremely model dependent and could also be very small.

The nonrelativistic potential model gives a remarkably good qualitative picture of the radiative transitions to the P-wave levels. But even where it succeeds best, it remains to be seen whether the potential model can be made the basis of a real quantitative description. If the 2.8 and 3.45 are indeed the \( J^P = 0^- \) states, then the situation for the \( M_1 \) transitions is much worse; spin-spin forces are evidently very important and we do not have even a qualitative understanding of the hyperfine splitting and the \( M_1 \) transitions.

To me the most intriguing and ambitious aspect of potential model phenomenology is the possibility that it might bring detailed experimental facts to bear on the theoretical investigations of the infrared behavior of the quark-gluon gauge theory and quark confinement. For instance, as discussed in Section II, the mass splittings between the three observed P-wave states may provide a clue as to the Lorentz transformation properties of the effective potential. On the other hand, without knowing how to solve the gauge theory we might be able to deduce the Lorentz nature of the effective nonrelativistic potential. This is one example, of how we could hope for specific experimental confrontations of theoretical ideas whose present principal contact with experiment is just the nonexistence of free quarks.
Though familiarity may numb our awareness of the problem, we still cannot say that we understand the narrowness of the charmonium states and the origin of the OIZ rule. Dual models provide one qualitative line of explanation; I have not discussed them here because of their reliance on the uncertain $\psi$ vector dominance hypothesis to make contact with radiative phenomena. In a field theoretical framework the only explanation which has been advanced is the short distance hypothesis discussed in Section IIIB.

If the $2.8$ and $3.45$ are the $0^-$ states I have assumed them to be, then the present experimental evidence bearing on the short distance hypothesis is mixed: the prediction for $\Gamma(\eta_c \rightarrow \text{hadrons})/\Gamma(\eta_c \rightarrow \gamma \gamma)$ fails by at least a factor $\sim 3$ and the prediction for $\Gamma(\eta'_c \rightarrow \text{hadrons})$ fails by at least one order of magnitude, but the very striking prediction for the $P$-wave states is at least qualitatively successful (we cannot now say more because of limited experimental statistics). Since the grossest failure—the prediction for $\Gamma(\eta'_c \rightarrow \text{hadrons})$—could in part be due to the poorly understood effect of spin-spin forces, the present situation might yet be consistent with the validity of the short distance hypothesis at the level of a factor $\sim 3$. This would probably be good enough to claim that short distance physics is connected with the origin of the OIZ rule. Future experimental data should be decisive, in particular, elucidation of the states at $2.8$ and $3.45$, more statistics for $\psi' \rightarrow \gamma \, ^3P_2$, and measurement of the inclusive photon yield and energy spectrum in $\psi \rightarrow \gamma + \text{hadrons}$. If these predictions are so badly violated that the short distance hypothesis seems not even qualitatively correct, then, at least in a field theory context, we have to start again from the beginning to try to understand the narrowness of the $\psi$ and the other manifestations of the OIZ rule.

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FOOTNOTES AND REFERENCES

* Research supported by the U. S. Energy Research and Development Administration.


1. F. Pierre, report to this conference.
2. W. Vernon, report to this conference.
3. H. Martyn, report to this conference.
4. G. Wagner, report to this conference.
19. This means that the more restrictive bound on the width of $\eta_c$ obtained in Ref. 5 may be unreliable.
25. This value is reported in the lectures of G. Feldman, Proceedings of SLAC Summer Institute, 1976, to be published. The bound $\lesssim 0.039$ is obtained by the LBL-SLAC collaboration assuming $m_c = 2.8$. The bound $\lesssim 0.03$ is obtained by the MPS collaboration and is less dependent on $m_c$.
26. M. Chanowitz, Phys. Rev. D12, 918 (1975); Okun and Voloshin, ITEP-95, 1976. The calculation was first done in Ref. 17 (see their footnote 12) but with an incorrect result.
31. These values were reported in Ref. 29.
44. See A. Silverman, Ref. 24.
45. For instance, if we assume an unsubtracted dispersion relation for the virtual vector meson-hadron coupling and furthermore assume equal residues for the $\rho$ and $\rho'$ contributions, then $(\frac{m_{\rho}}{m_{\rho'}})^2$ is the number of merit.
46. B. L. Ioffe, ITEP-24, Moscow, 1975.
47. Zakharov, Ioffe, and Okun, JETP 48, 1635 (1975).
48. Sivers, Townsend, and West, Phys. Rev. D__
50. Another suggestion which has appeared just before this preprint goes to press is that the $3.45$ is the $^1D_2$ $\bar{c}c$ state--see

H. Harari, CERN-TH 2222, 1976. This suggestion requires large hyperfine splitting to explain the low mass of the $^1D_2$, so that the $\bar{c}c'$ mass is also expected to be low, say $\sim 200 - 300$ MeV below the $\psi'$. 
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