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NEW PARTICLE SPECTROSCOPY *)

J. D. Jackson **)

CERN - Geneva

A B S T R A C T

The spectroscopy of the new mesonic states with masses above 2.8 GeV is discussed in terms of the charmonium picture based on the hopes of QCD. The status of our understanding of the level scheme, multiplet splittings, transition amplitudes and hadronic annihilation rates is reviewed. Brief mention is given of dynamics above the charm threshold. Despite absence of a firm theoretical basis for some of its basic assumptions, the charmonium picture is remarkably successful. Only details of the spin dependence and/or the pseudoscalars seem serious problems.

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**) On leave from the University of California, Berkeley, 1976-77.

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NEW PARTICLE SPECTROSCOPY

J.D. Jackson*), CERN, Geneva, Switzerland

1. INTRODUCTION

This paper is a review of the heavy mesonic states with "hidden charm" ($M \simeq 3.5 \pm 1.0$ GeV) discovered since the summer of 1974, mostly in e^+e^- annihilation. It does not cover, except in passing, the overtly charmed hadrons (D, D*, Λ_c, ...) or the "baryonium" states, or the recently discovered resonance at 9.6 GeV. For these topics, the reader should consult the reports, respectively, of Litke and Goldhaber, of Hemingway, and of Lederman in these Proceedings.

In the space available, no justice can be done to the beautiful experimental results from both SPEAR and DORIS. Only the briefest sketch of the deductions can be given, with an occasional illustrative example. The experimental conclusions are given an "anschaulich" interpretation with the quark model based on quantum chromodynamics (QCD); areas of success and of difficulties are spelled out.

The *basic* experimental facts, accumulated over 2½ years, are[1]:

i) In spite of their surprising relative stability against decay into the "old", light particles, the new particles or resonances are hadrons that conserve P, C, G, I, and S in their interactions.

ii) The first discovered states, $\psi(3095)$ and $\psi'(3684)$, have $J^{PC} = 1^{--}$, isospin I = 0, and normal vector-meson couplings to e^+e^-. There is evidence that $\psi(3095)$ is predominantly a SU(3) singlet.

iii) The χ states at 3413 MeV, 3508 MeV, and 3552 MeV, formed from $\psi'(3684) \to \to \gamma\chi$ have C = +1 and I = 0.

iv) The $\psi'(3684)$, with a total width of 228 keV, decays dominantly (53%) to $\psi(3095)$ via $\pi\pi$ or η emission, and significantly (22%) to the χ states via γ emission.

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*) On leave from the University of California, Berkeley, Calif., 1976-77.

v) The newly discovered $\psi(3772)$ state[2], 88 MeV above $\psi'(3684)$, decays predominantly into $D\bar{D}$, even though only 40 MeV above the $D\bar{D}$ threshold.

The conclusion from these facts seems inescapable. These states are hadrons with "hidden charm" in the sense of Bjorken and Glashow[3]. In the language of the quark model, the new particles are bound or resonant states of a fourth, massive quark (c) and its antiparticle ($\bar{c}$). The c quark differs from the familiar light quarks u, d, s of SU(3), not only in mass but also in possessing a non-zero value of a new additive hadronic quantum number called charm ($\mathbb{C}$). The ($c\bar{c}$) states possess "hidden charm", having $\mathbb{C} = 0$ like ordinary hadrons, but consisting of charmed quarks.

2. THEORETICAL FRAMEWORK

2.1 Quantum chromodynamics

The spectroscopy of the new particles is described in terms of confined, spin-½ quarks[4]. The formal theoretical underspinning is provided by QCD, a non-Abelian gauge theory of strong interactions in which flavourful and colourful quarks interact via exchange of an octet of colourful, massless, vector gluons. The theory is thus

$$[SU(3)]_{colour} \times [SU(n)]_{flavour}$$

where n = 2, 3, 4, ... With u, d, s, c quarks, n = 4. The "colour" symmetry is not developed; only colour singlet states are supposed to exist in nature[5].

An essential feature of QCD is its "asymptotic freedom". Crudely speaking, asymptotic freedom means that the renormalized coupling "constant" is not a constant, but varies with the energy or momentum flowing through the vertex and in particular, decreases towards zero as the momentum increases towards infinity. High momentum corresponds to short distances. Thus, as $r \to 0$, we expect weak coupling and hope that for close separations of quarks we can use perturbation theory in analogy with QED. In the opposite limit of large separations (small momenta) there will be very strong coupling, so strong that quarks will be confined. This domain of the theory is extremely complicated. Quark confinement is an unsolved problem, but for the spectroscopy guide posts are provided by dual string models, monopole models with gauge-field flux tubes between quarks, etc. These, as well as some lattice calculations in QCD, indicate that the "potential energy" between quark and antiquark increases linearly at large separations, with an energy per unit length independent of the quark masses or flavours.

2.2 Zweig's rule

An important aspect of the new particles is their amazing reluctance to decay into ordinary hadrons [basic experimental facts (iv) and (v)]. This is attributed to Zweig's rule[6], whereby a process is forbidden if the quarks in the initial (final) state must annihilate (materialize) rather than flow through and be part of the final (initial) hadrons. Figure 1a shows a Zweig-allowed process, relevant for $\psi(3772) \to D\bar{D}$, while Fig. 1b shows a Zweig-forbidden process, the decay of a heavy ψ particle into ordinary, light hadrons (annihilation of $c\bar{c}$ and materialization of light $q\bar{q}$ pairs). These diagrams are not quite Zweig's rule diagrams because they contain gluon as well as quark lines. This is to illustrate the QCD "explanation" of Zweig's rule: the Zweig-forbidden decay of a heavy meson (e.g. $c\bar{c}$) necessarily involves the emission of "hard" gluons, because the large mass of the meson is transmitted through a small number of gluons. Because of asymptotic freedom, the coupling is weaker, the harder the gluon. Therefore the decay of massive states (made of new, different quarks) into particles consisting of old, light quarks is relatively inhibited.

There are logical loopholes in this argument[7] and it is only qualitative at best. Zweig's rule is a reality, nevertheless. The new particles below the $D\bar{D}$ threshold are remarkably stable. This makes the spectroscopy richer experimentally and simpler to interpret theoretically.

3. MODEL FOR ψ SPECTROSCOPY

3.1 Basic assumptions

The ideas of QCD apply equally, in principle at least, to the "old" mesons and baryons as well as the new ones, but the smaller masses (and therefore stronger couplings and relativistic motion) and many open channels make difficult implementation of a programme of spectroscopic calculations. For the new particles (the ψ sector) we deal with a simpler situation. The conventional picture contains the following ingredients:

a) massive I = 0 c quarks with $m_c \approx 1.3$-1.7 GeV, $e_Q = 2/3$;

b) non-relativistic dynamics as a lowest-order approximation;

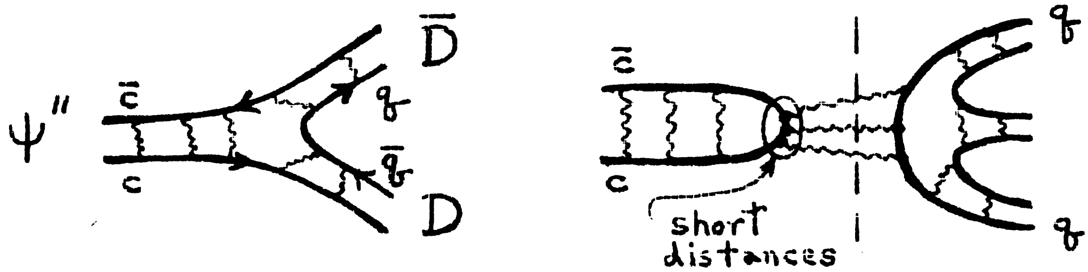
c) small (?) relativistic corrections;

d) one-gluon exchange potential at short distances;

e) confinement potential at large distances (in considering the relativistic corrections there is a question of the Lorentz-group properties of the confining potential -- see below);

f) coupled channels are only important close to, or above, the charm ($D\bar{D}$) threshold (Fig. 1a shows the Zweig-allowed coupling of a ψ state to $D\bar{D}$);

g) annihilation (Zweig-forbidden decays to light particles) is a short-distance phenomenon ($\Delta r \approx \hbar/m_c c$) that can be treated in lowest order QCD.



(a) (b)

Fig. 1

In applications, the annihilation process is treated as a two-step affair. QCD is used to calculate the $c\bar{c}$ annihilation into the *minimum* number of gluons (to the left of the dashed line in Fig. 1b). Then unitarity is invoked to argue that the width so computed is approximately the total width for the decay of the $c\bar{c}$ state into old, light hadrons. As already indicated, this hypothesis does not have a secure basis. We nevertheless adopt it tentatively and look for confirmatory or otherwise evidence in the experimental data. As is well known by now[8], application to the hadronic width of the $\psi(3095)$, $\Gamma_h(\psi) \simeq 50$ keV, gives the QCD "fine structure constant", $\alpha_s \simeq 0.19$, for states of $M \approx 3$ GeV.

3.2 Qualitative aspects of the spectrum of states in a confining potential

In the static limit the effective potential between quark and antiquark is assumed to be of the form,

$$V(r) = -\frac{4\alpha_s/3}{r} + V_c(r) \tag{1}$$

where the first term is the one-gluon-exchange potential, projected onto colour-singlet states, and the second is the confinement potential, assumed to increase linearly with r for large separations. On general grounds of continuity one expects the energy levels to exhibit a pattern intermediate between the spectrum of the hydrogen atom ($V_C = 0$) and the isotropic harmonic oscillator ($\alpha_s = 0$, $V_C \propto r^2$) -- see Fig. 2. The expected order of the lowest five states (or multiplets, when spin is included) is therefore

$$E(1s) < E(1p) < E(2s) < E(1d) < E(2p) .$$

This ordering for the Schrödinger equation has been proved for a wide class of potentials V(r), Eq. (1), by Grosse and Martin[9]. The constraints on the confining potential V_C are

(A) $$\frac{d^3}{dr^3}\left(r^2 V_c\right) > 0 \quad \text{for all r, and}$$

$$\lim_{r\to 0}\left[2rV_c + r^2\frac{dV_c}{dr}\right] = 0 ;$$

(B) $$\frac{d}{dr}\left[\frac{1}{r}\frac{d}{dr}\left(2V_c + r\frac{dV_c}{dr}\right)\right] < 0 \quad \text{for all r.}$$

The first inequality follows instantly from the presence of the repulsive centrifugal barrier for $\ell = 1$; the second and fourth require (A) alone, while the third depends on both (A) and (B). The conditions (A) and (B) are not overly restrictive. For a power law potential, $V_C \propto r^\alpha$, for example, the constraint is merely $0 < \alpha < 2$. The order of appearance of successive levels with increasing energy is thus established to be insensitive to the details of the potential, provided only that it is confining.

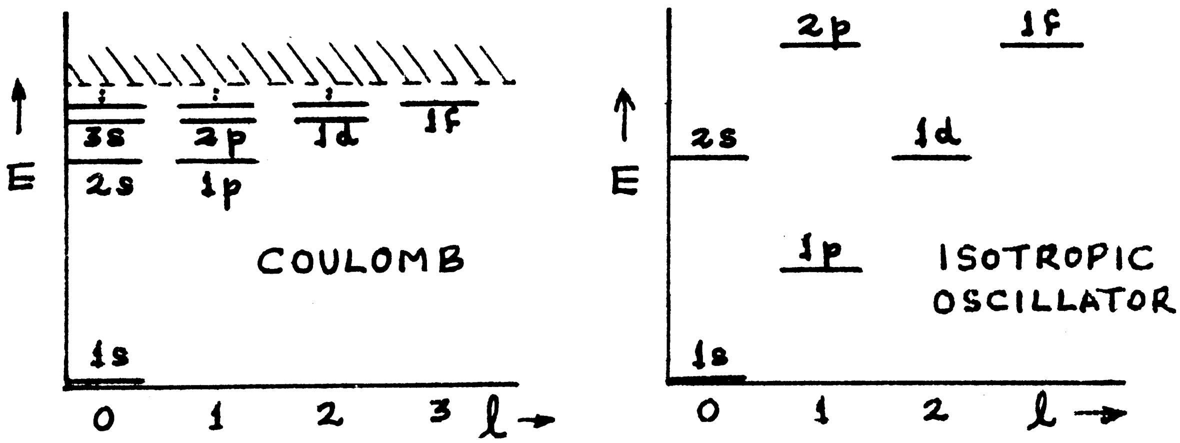


Fig. 2

A result of a different nature[10] concerns the relative magnitudes of the 1s and 2s wave functions at the origin, of interest for the coupling of $c\bar{c}$ to photons or gluons (Γ_e, Γ_h, $\Gamma_{\gamma\gamma}$, ...): If $V'' \gtrless 0$ for all r, then

$|\psi_{2s}(0)/\psi_{1s}(0)| \gtrless 1$. The popular potential, $V = -(k_1/r) + k_2 r$, $k_1, k_2 > 0$, satisfies the lower inequality and therefore has $\Gamma_e(\psi') < \Gamma_e(\psi)$, in agreement with experiment.

The general aspects of the ordering of energy levels is given concrete realization in Fig. 3, where the levels for various potentials are displayed, all scaled to have the 1s-2s interval in agreement with the ψ'-ψ mass difference. As noted by Gottfried[11], the level structure is remarkably stable, with large differences developing only above 4 GeV. The triplet and singlet spin multiplets from each Schrödinger level are indicated by the arrows (directions not significant). The longest arrows represent $J^{PC} = 1^{--}$ states; the next longest, states with C = +1 that can be reached from the 1^{--} states by single photon emission; the shortest arrows correspond to states with C = -1, but $J^P \neq 1^-$. The latter are difficult to observe in e^+e^- annihilation. They are formed only by 2γ cascades or by hadronic decays from higher mass states.

The observed energy levels are indicated on the left side of Fig. 3. In the 1 GeV interval from 2.8 to 3.8 GeV, eight states are now known. This

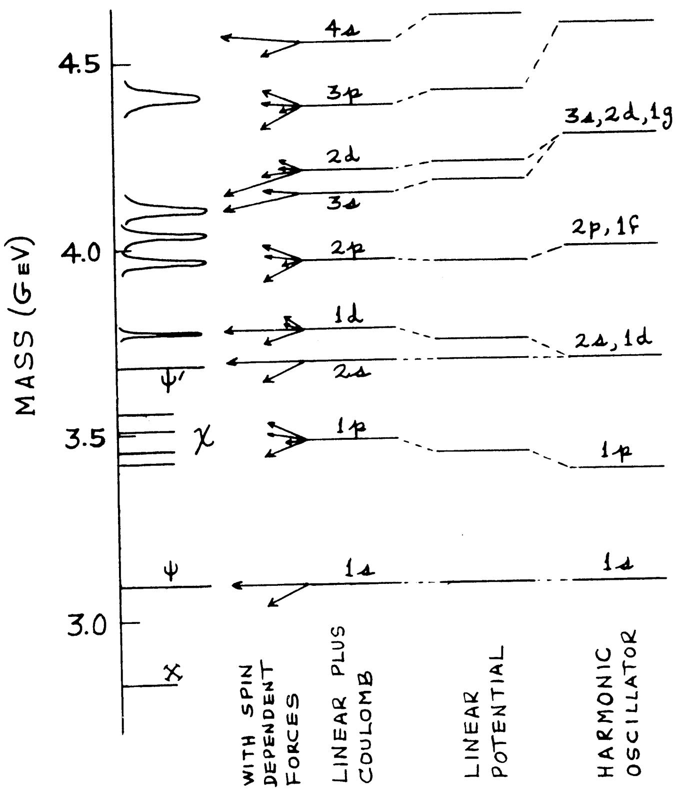


Fig. 3

number corresponds almost exactly to the number of longest and next longest arrows associated with the 1s, 1p, 2s, 1d levels, nine in all[12]. This alone makes one entertain very seriously the simple $c\bar{c}$ bound state picture.

3.3 Spin-dependent forces (relativistic corrections)

As is well known in atomic physics, the static (electric) interaction is augmented by spin-dependent magnetic interactions (dipole-dipole, dipole-motional magnetic field) whose specific forms depend on the 4-vector nature of the coupling of the particles to the photon. In QCD, the short-range part of the potential (one-gluon exchange) should have spin-dependent relativistic corrections of a completely analogous character. For the longer-range, confining effective potential, it is not so clear. Some authors[13] assume that the entire potential has 4-vector origins, while others assume that the confining potential [V_C of Eq. (1)] comes from a Lorentz scalar coupling. These different assumptions give rise to different spin-dependent forces, as indicated for an equal-mass two-fermion system in Table 1. With $\mathcal{V}(r) = -e^2/r$, the 4-vector column yields the familiar results of the lowest-order corrections for positronium (to which must be added the annihilation contribution, $e^+e^- \to \gamma_V \to e^+e^-$). Note particularly that the Lorentz 4-scalar coupling gives a spin-orbit interaction of opposite sign to that of the 4-vector [inverted multiplets in nuclei[14]]. The relative amount of spin-orbit force versus tensor and $\vec{\sigma}_1 \cdot \vec{\sigma}_2$ forces can evidently be adjusted by mixing the two types of relativistic couplings.

Table 1: Static and quasi-static interactions

Type of interaction	Lorentz property: 4-vector $\gamma_\mu \otimes \gamma_\mu$	Lorentz property: 4-scalar $1 \otimes 1$
Static potential	$\mathcal{V}(r)$	$\mathcal{S}(r)$
Spin-orbit	$\frac{3}{2m^2}\frac{1}{r}\frac{d\mathcal{V}}{dr}\vec{L}\cdot\vec{S}$	$-\frac{1}{2m^2}\frac{1}{r}\frac{d\mathcal{S}}{dr}\vec{L}\cdot\vec{S}$
Tensor force	$\frac{1}{12m^2} S_{12}\left[\frac{1}{r}\frac{d\mathcal{V}}{dr} - \frac{d^2\mathcal{V}}{dr^2}\right]$	0
Fermi hyperfine	$\frac{\vec{\sigma}_1\cdot\vec{\sigma}_2}{6m^2}\nabla^2\mathcal{V}$	0

4. COMPARISON OF EXPERIMENT WITH MODELS

4.1 Experimental results

The wealth of experimental detail and inferences accumulated in the recent past can be inferred from Fig. 4, patterned after Feldman and Perl[1]. It summarizes the presently known facts on transitions and branching ratios among the states from 2.8 to 3.8 GeV. The one- or two-digit numbers are branching ratios in per cent. Typically these have relative uncertainties of order 10%, sometimes as much as 50%. Dotted lines indicate doubt, in the spin-parity assignments of X(2830) and χ(3455), for example.

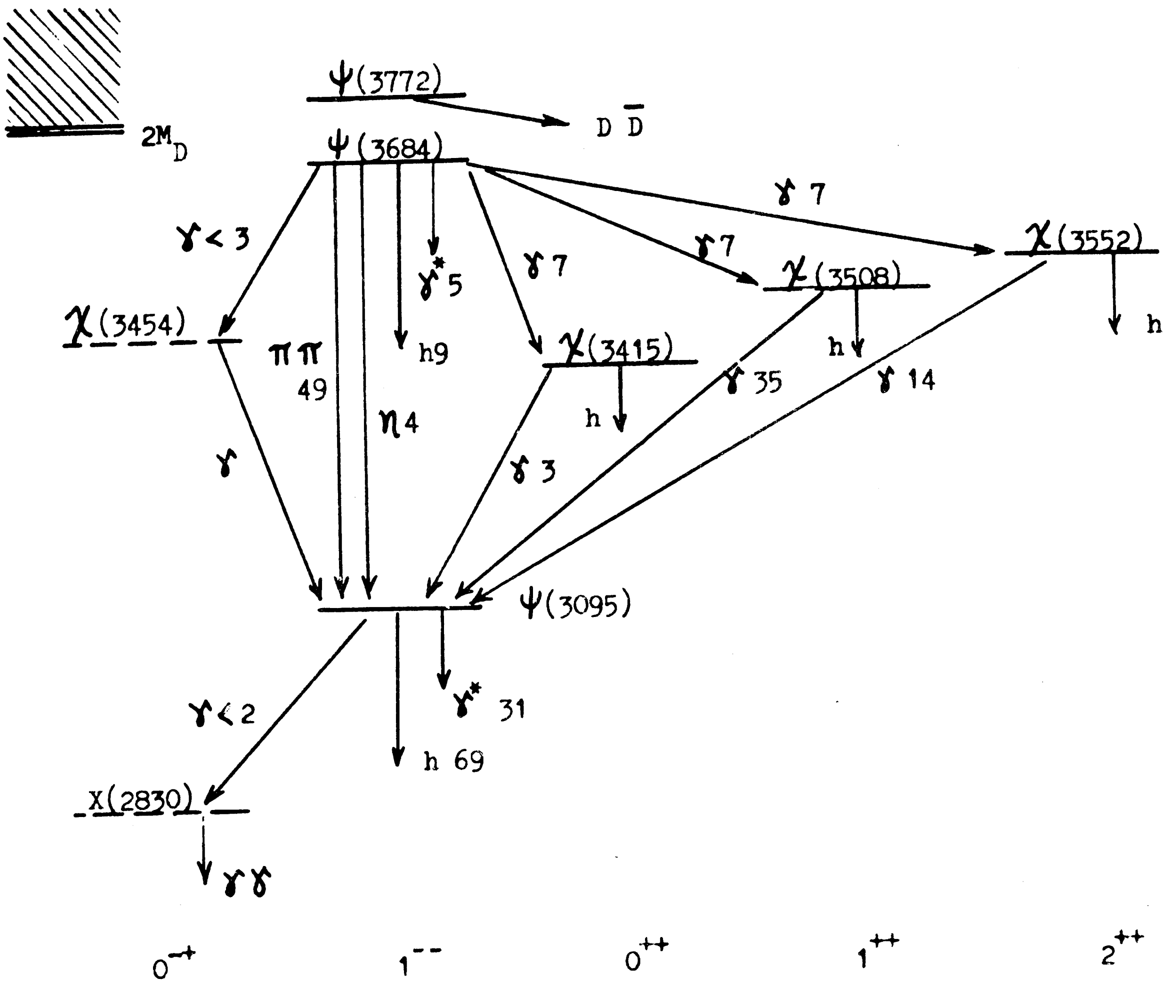


Fig. 4

Space limitations prevent discussion of the process by which many different experimental observations are interconnected and interpreted to yield the results of Fig. 4 (and many other conclusions). One example will suffice to give the flavour and to indicate the firmness of some of the deductions. Consider the question of the isospin and spin-parity assignments of the χ(3415) and χ(3552). The logical steps of observation and inference are as follows:

a) $e^+e^- \to \gamma_V \to \psi, \psi'$ with parity conservation,

$\therefore$ ψ and ψ' have $J^{PC} = 1^{--}$, I = 0, 1;

b) $\psi \to$ odd number of pions in direct decay to hadrons,

$\therefore$ G is odd and I = 0 for $\psi(3095)$;

c) $\psi'(3684) \to \eta\, \psi(3095)$ (also $\pi^0\pi^0$, $\pi^+\pi^-$ in ratio 1:2)

$\therefore$ I = 0 for $\psi'(3684)$;

d) $\psi'(3684) \to \gamma\, \chi(3415, 3508, 3552)$,

$\therefore$ C = +1, I = 0, 1 for $\chi(3415, 3508, 3552)$;

e) $\chi(3415, 3508, 3552) \to$ even number of pions,

$\therefore$ G is even and I = 0 for these three χ states;

f) $\chi(3415, 3552) \to \pi^+\pi^-$, K^+K^-,

$\therefore$ $J^{PC} = 0^{++}, 2^{++}, 4^{++}, \ldots$ for $\chi(3415, 3552)$;

g) Photon angular distribution in $e^+e^- \to \psi' \to \gamma\, \chi(3415)$ is $(1 + \alpha \cos^2 \theta)$, with $\alpha \simeq 1$, as expected for a J = 1 to J = 0 (E1 or M1) radiative transition,

$\therefore$ $J^{PC} = 0^{++}$ for $\chi(3415)$ is consistent, but not proved;

h) Photon angular distribution for $\chi(3552)$ is *not* $(1 + \cos^2 \theta)$ at the level of two or three standard deviations,

$\therefore$ $J^{PC} = 2^{++}, 4^{++}, \ldots$ for $\chi(3552)$.

The logical chain of reasonably firm experimental inferences stops here. Comparison with theoretical expectations is the basis for more specific assignments. As already mentioned, the number and ordering of the multiplets are exactly as expected for a confined fermion-antifermion system. Furthermore, if the dominant interaction has 4-vector coupling, the ordering *within* the multiplets should be "normal", i.e. increasing J with increasing energy. Three of the four χ states are evidently candidates for the ^3P_J states, with $J^{PC} = 0^{++}, 1^{++}, 2^{++}$. Since the $\chi(3552)$ cannot have J = 0, while $\chi(3514)$ can, the *assignments* of $J^{PC} = 2^{++}$ and 0^{++}, respectively, are consistent with theoretical expectations[15].

Either the $\chi(3508)$ or $\chi(3454)$ could be the $^3P_1$ state. Here and for the above assignments, simple considerations of the radiative transitions favour a particular assignment. Transitions from the ψ' ($^3S_1$, $J^{PC} = 1^{--}$) to the ^3P_J states should be predominantly E1, with roughly the same transitional dipole moment. The branching ratios should thus be proportional to $(2J + 1)k^3$.

Experimentally, the branching ratio for $\psi' \to \gamma\ \chi(3454)$ is less than 2.5%, while those for the other states are all about 7%. The rough equality of radiative rates for the states at 3415, 3508 and 3552 MeV is consistent with the $(2J + 1)k^3$ recipe provided the "normal" ordering applies, the increase of J with mass balancing the decrease in photon energy (cubed). The inverted ordering is in gross disagreement with experiment. The smaller branching ratio to $\chi(3454)$ cannot be understood easily within the framework of E1 transitions to the ^3P_J states. Since this state has other peculiarities (see Section 5.1), it is not assigned to the ^3P_J multiplet. Assignment of $J^{PC} = 1^{++}$ to the $\chi(3508)$ is consistent with its lack of decay into $\pi^+\pi^-$ and the evidence from photon angular distributions that $J \neq 0$. We are thus left with the χ spin-parity assignments of Fig. 4, plausible or better to the believer, far from absolutely established for the sceptic[16].

4.2 Schrödinger wave function properties

In a comparison of models with experiment there are several levels. The first is at the lowest order of a static potential, for which we ask whether the gross energetic and spatial properties (obtainable from the Schrödinger equation) can be "understood", i.e. plausibly correlated. These properties include the spacing and ordering of the spinless energy levels, already discussed (Fig. 3), the values of the s-state wave functions at the origin, deduced from the formula,

$$\Gamma_e = 16\pi \frac{\alpha^2 e_Q^2}{M^2} |\psi(0)|^2 \qquad (2)$$

for the coupling of $J = 1^{--}$ s-states to e^+e^- or $\mu^+\mu^-$, and the dipole moments for E1 transitions. The answer is yes, the static properties are well understood, at least up to the charm threshold. This can be seen in a quick and dirty way by use of an oscillator "template" for estimation of matrix elements[5], but more careful calculations have been performed[17]. A few pieces of data are needed to fix the parameters (α_s, m_c, dV_C/dr); all else follows. A potential of the form of Eq. (1) with $V_C = r/a^2$ and parameters $\alpha_s = 0.2$, $a^{-2} = 0.2\ \text{GeV}^2$, and $m_c = 1.4$ GeV, does a fair job of the gross spectrum and spatial properties: ψ'-ψ mass difference and $\Gamma_e(\psi)$, fitted; $\Gamma_e(\psi') \simeq 1.5\ \Gamma_{exp}$; centre of gravity of p-states $\sim$ 50-60 MeV too low; E1 matrix elements within factors of 1.2-1.5 [18]; 1d level predicted at 80-90 MeV above the 2s state.

4.3 Spin-dependent properties

The (relativistic) corrections to the spinless static limit appear most obviously in two places: one is in spin-dependent contributions to the interactions (see Table 1) and the other is in spin-flip (M1) radiative transitions.

4.3.1 Multiplet fine structure

From Table 1 we learn that the spin-dependent part of the $c\bar{c}$ interaction is

$$\Delta V = A(r)\vec{s}_1\cdot\vec{s}_2 + B(r)\vec{L}\cdot\vec{S} + C(r)S_{12} \qquad (3)$$

For the s-states this gives the splitting between the vector ($V = {}^3S_1$) and pseudoscalar ($PS = {}^1S_0$). If the tensor force contribution to the ${}^3S_1$ binding is neglected, only the short-range part of the potential Eq. (1) is kept, and the level splitting is treated in first-order perturbation theory, one finds a connection between Γ_e, Eq. (2), and the V-PS mass difference, $\Delta M \simeq \simeq 2\,\alpha_s\,\Gamma_e(\psi)/\alpha^2$. With $\Gamma_e(\psi) = 4.8$ keV, one obtains $\Delta M \simeq 34$ MeV. This is to be compared with the ψ-X(2830) mass difference of 265 MeV, assuming the X(2830) is the ψ's pseudoscalar partner. Other estimates have been made[19,20]. It is difficult to obtain splittings larger than 100-150 MeV. For example, with the parameters given at the end of Section 4.2 for V(r) and the assumption that it is derived from 4-vector couplings, the $\vec{\sigma}_1\cdot\vec{\sigma}_2$ interaction of Table 1 gives 130 MeV, with the short-range contribution being 49 MeV. If the $\chi(3454)$ is the n = 2 pseudoscalar, the same story holds -- a V-PS mass difference of 230 MeV compared with theoretical estimates of $\lesssim$ 100 MeV.

For the splittings inside the ${}^3P_J$ multiplet, both the spin-orbit and tensor forces operate. Assuming perturbation theory again, the observed splittings lead to the empirical expectation values, $\langle B\rangle \simeq 34$ MeV, $\langle C\rangle \simeq \simeq 10$ MeV. Straightforward calculations in a model with 4-vector couplings ($\mathcal{S} = 0$ in Table 1) yield theoretical values roughly twice and half of experiment for $\langle B\rangle$ and $\langle C\rangle$, respectively. From Table 1 we see that B and C are both proportional to m_c^{-2}. Thus an increase of the quark mass can force the $\langle B\rangle$ value closer to experiment, but at the expense of a still smaller tensor force contribution[20]. Decrease of the quark mass does the reverse. This dilemma leads one to consider the flexibility provided by a potential that is partly Lorentz 4-vector and partly 4-scalar. As already discussed in Section 3.3, the relative proportions of spin-orbit and tensor interactions can thereby be adjusted. A relativistic calculation equivalent to taking $\mathcal{V}(r) = -4\alpha_s/3r$

and $S(r) = V_C = r/a^2$ has been performed[21]. Good agreement with the p-state splittings is obtained, along with a generally acceptable spectrum and spatial properties of the various states. An unfortunate consequence of this model, at least from the point of view of asymptotically-free QCD, is the necessity in the potential of $\alpha_s \simeq 0.4$, twice as large as the value found from $\psi \to$ $\to$ hadrons.

4.3.2 Magnetic dipole transitions

The radiative transitions between vector and pseudoscalar states are M1 and in the quark model involve quark spin-flip. The effective magnetic dipole moment is $\mu_{eff} = \frac{1}{2}[(e_1/m_1) - (e_2/m_2)]$. The transition amplitude is the product of this moment and a spatial overlap integral. The M1 transitions can be "favoured" (large overlap integral between singlet and triplet s-states with the same radial quantum number) or "hindered" [small overlap integral because of different radial quantum number, rate very uncertain because of comparable coherent contributions from relativistic corrections to the wave functions[22]]. From the light quark sector it is known that the simple quark model estimates of M1 rates are only good to a factor of two or three.

The experimental situation and the theoretical estimates of the rates are shown in Table 2. It has been assumed that the X(2830) and $\chi(3454)$ are the pseudoscalar partners of the ψ and ψ', respectively. Three out of the four transitions have been observed, but so far only upper limits for the widths exist from the inclusive photon spectra[23]. Given the poor reliability of the theoretical estimates, serious disagreement exists at present only for the transition $\psi(3095) \to \gamma$ X(2830), where the theoretical value is more than a factor of 20 greater than the present upper limit.

The status of the M1 transitions and the mass splittings within multiplets indicates that our understanding of spin effects in charmonium is minimal. The

Table 2: M1 transitions

Initial state	Final state	Favoured/ Hindered	Seen/ Not seen	Widths (keV)	
				Experimental	Theoretical
$\psi(3095)$	X(2830)	F	Seen in 3γ	< 1.3	30
$\psi'(3684)$	$\chi(3454)$	F	Seen	< 6	17
$\psi'(3684)$	X(2830)	H	Not seen	< 2	($\sim$ 8)
$\chi(3454)$	$\psi(3095)$	H	Seen	-	($\sim$ 0.1)

mass splittings imply that *gluonic* "magnetic" couplings are larger than expected, while the M1 transitions imply the opposite tendency for the *photonic* magnetic couplings. What does it mean? Because most of the problems involve the pseudoscalar states, some physicists are content to remind us that the pseudoscalars (η, η') have always been a problem! (A fact, but not a comfort, not an answer.) To be fair, it should be said that some theorists claim to "understand" the M1 transitions for both new and old mesons in terms of a broken SU(4) symmetry scheme for the electromagnetic current[24].

4.4 The new state at 3772 MeV

Just at a time when doubts about the simple $c\bar{c}$ bound state picture begin to accumulate [because of the problems just described and/or doubts about the assignments for X(2830) and χ(3454) -- see Section 5.1], the model comes through with flying colours! As detailed in the paper by Litke, these Proceedings, a new direct-channel resonance in e^+e^- annihilation has been observed at M = 3772 MeV, just 88 MeV above the ψ'(3684). It decays strongly into $D\bar{D}$ (thresholds at 3730 and 3738 MeV for $D^0\bar{D}^0$ and D^+D^-, respectively) and has a total width, $\Gamma_t \simeq 28$ MeV. Its magnitude is only about 2 units in R, corresponding to an electronic width, $\Gamma_e \simeq 0.37$ keV, about one sixth of the Γ_e of its neighbour, ψ'(3684).

In the simple $c\bar{c}$ model, this state is the long-expected $^3D_1$ ($J^{PC} = 1^{--}$), a member of the 1d multiplet (see Fig. 3). As a state with $\ell = 2$, its wave function vanishes at the origin. Use of Eq. (2) therefore gives $\Gamma_e = 0$, contrary to observation. That is just the freshman physics answer. First of all, the tensor force mixes $^3S_1$ and $^3D_1$ states, as does the coupling to other channels even without tensor forces. With only 88 MeV separating ψ'(3684) and ψ''(3772), there is undoubtedly considerable mixing, the 3684 MeV state being predominantly $^3S_1$ ($\Gamma_e = 2.1 \pm 0.3$ keV) and the 3772 MeV state ($\Gamma_e \simeq 0.37$ keV) being mostly $^3D_1$. Even within the single ($c\bar{c}$) channel description, there are complications in deducing the mixing from the ratio of electronic widths. It is not true that $\Gamma_e = 0$ for a d-state. For a mixed $^3S_1$-$^3D_1$ state, one finds

$$\Gamma_e = \frac{4\alpha^2 e_Q^2}{M^2}\left| a R_s(0) + \frac{b}{M^2} R_d''(0)\right|^2 \qquad (4)$$

where $R_\ell(r)$ is the radial wave function for orbital angular momentum ℓ and a,b are numerical coefficients. A pure d-state thus has a non-vanishing Γ_e

from the second derivative of its radial wave function. Simple estimates give values of Γ_e for a pure d-state of one tenth the experimental width, indicating the need for mixing, but the presence in the amplitude of more than the first term (admixed s-state) complicates the evaluation[25].

Since the state lies above the $D\bar{D}$ thresholds any serious theoretical description must include the coupling of $c\bar{c}$ to $(c\bar{q})(\bar{c}q)$ and $(c\bar{c})(q\bar{q})$. (The use of the terminology "1d" or "$^3D_1 + \varepsilon\ ^3S_1$" must then be understood to refer only to the $c\bar{c}$ component of its wave function.) More than a year ago the Cornell group predicted the value of Γ_e to within a factor of two, the dominant decay into $D\bar{D}$ (rather than $\psi\pi\pi$, for example), and approximately the correct total width (given the mass at 3772 MeV and the D^0, D^+ masses)[26,27]. These are non-trivial accomplishments.

4.5 Above the charm threshold

The comparison of models with experiment above 3.8 GeV is much less complete than below. There are at least two reasons. One is that meaningful theoretical calculations are difficult because of the existence of many coupled channels. The other is that the experimental situation is quite complicated. Despite heroic efforts and brilliant accomplishments, outlined in the paper by Goldhaber, these Proceedings, very little is known about dynamics. The experimental concern so far has been focused, with good reason, on the masses, strong and weak decay modes, spin and parity assignments of the charmed particles (D^0, D^+, D^{*0}, D^{*+}, ...). There is no space to go into the phenomenology of these considerations. We content ourselves with remarks on one specific point -- the initially rather surprising relative production of $D\bar{D}$, $D^*\bar{D}$+c.c., $D^*\bar{D}^*$ at W = 4.028 GeV.

Immediately after the discovery of the $D^0(1864)$, and $D^{*0}(2006)$ in the recoil spectrum, theorists[28,29] indicated the expectations for the recoil spectrum at a given energy, the structure in R as a function of energy and other details. The early experimental finding of negligible $D\bar{D}$ production, compared to $\bar{D}D^*$+c.c. and $D^*\bar{D}^*$, at W = 4.028 GeV was explained qualitatively by a simple statistical weight argument that yields 1:4:7 for the relative abundances of the three final states. These must be modified, of course, by kinematic p-wave phase-space factors:

$$D\bar{D} : (\bar{D}D^* + c.c.) : \bar{D}^*D^* \simeq 1 : 4\left(\frac{p_2}{p_1}\right)^3 : 7\left(\frac{p_3}{p_1}\right)^3$$

At W = 4.028 GeV, the momenta for D^0, D^{*0} are p_1 = 0.763, p_2 = 0.558, p_3 = = 0.179, and the ratios become 1:1.6:0.09, the last so small because the c.m.s. energy is only 14 MeV above threshold. Experimentally, the ratios are 1: ∿10: ∿10. This implies an incredible enhancement of the $D^*\bar{D}^*$ mode over naïve expectations and led to the idea of "molecular charmonium"[30), i.e. the presumed resonant state at 4.028 GeV (and others) being a 4-quark system ($c\bar{c}q\bar{q}$) that can easily organize itself into pairs of charmed mesons interacting via exchange of vector mesons and pions. (Recall that once upon a time the ρ resonance was thought to be a resonant state formed of two pions interacting via ρ-meson exchange!) The relative abundances of the different two-body (and more) channels is a detailed dynamical question that has yet to be worked out in detail. Reference 30 nevertheless contains numerous qualitative points about a rich "molecular" spectroscopy attendant on this picture.

An alternative (if not totally orthogonal) explanation is given by the Cornell coupled-channel approach[17). The radial nodes of the 3s wave function in the $c\bar{c}$ sector are held responsible for the drastic departure of the relative intensities of the two-body modes from the phase-space modified 1:4:7. This can be understood in simple terms of wave function overlap. The interaction in the Cornell model is

$$H_{int} = \frac{1}{2}\int d^3r \int d^3r' \rho(\vec{r}') U(|\vec{r}-\vec{r}'|) \rho(\vec{r}) \qquad (5)$$

where $\rho(\vec{r})$ is the quark colour density and U(r) is a universal potential, taken to be r/a^2. In the $c\bar{c}$ sector this gives the familiar model with a linear confining potential. The interaction couples to other channels as well, for example, $c\bar{c} \to D\bar{D}$. Suppose for the moment that the D meson is very compact compared to the extent of the $c\bar{c}$ 3s state (this is far from true!) and that we calculate the transition amplitude in perturbation theory. The amplitude will be some sort of convolution of the wave functions sketched in Fig. 5. Evidently the nodes in the 3s radial wave function will appear as corresponding nodes in the transition amplitude as a function of momentum, giving a modulation to the normal over-all decrease of the amplitude with increasing momentum. The actual calculations are not so simple -- the D, D^* wave functions are in fact as large as, or larger than, the 3s $c\bar{c}$ system, etc. But the basic interference effect is as indicated by the simple overlap argument.

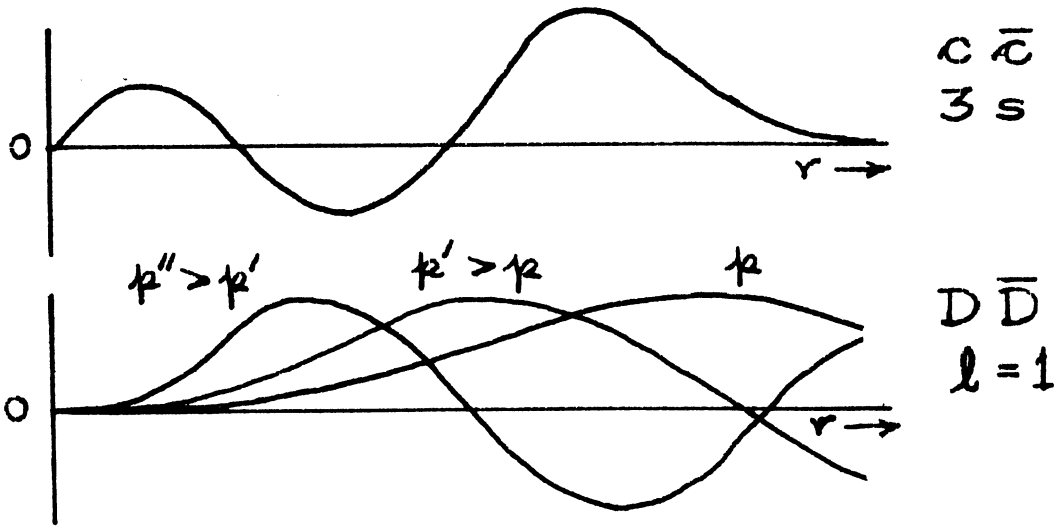


Fig. 5

The modulation occurs in the different channels at more or less the same momentum, but, because of different masses, at different c.m.s. energies. If data happen to be taken at an energy where the $D\bar{D}$ amplitude has a node, the production of that channel will be suppressed below naïve expectations. Figure 6 shows the calculations of ΔR of more than a year ago by the Cornell group[29]. The 3.77 GeV peak (decaying entirely into $D\bar{D}$) is shown on the left. The short-dashed, long-dashed, and thin solid curves give the contributions of the $D\bar{D}$, $\bar{D}D^*$+c.c., and $D^*\bar{D}^*$ channels, respectively. The canonical 1:4:7 ratio is present, but modified by the nodes. At $W \simeq 4.06$ GeV (higher than 4.03 because of an assumed D^* mass of 2.020 instead of 2.006), a node suppresses the $D\bar{D}$ production, while the $\bar{D}D^*$ and $D^*\bar{D}^*$ rates are equal, in rough agreement with experiment. Drastic enough "form factors"[28] for each channel can produce this result at one energy, of course. The signature of the modulation by the 3s radial nodes is a characteristic variation of the relative amounts of the three channels -- below *and* above 4.028 GeV the proportion of $D\bar{D}$ should increase relative to, say, $\bar{D}D^*$+c.c. With form factors to cut off the p^3 threshold rise, the amount of $D\bar{D}$ should fall monotonically above 4 GeV.

Given the success of the simple $c\bar{c}$ picture below the charm threshold, the Cornell model has a strong appeal to an old nuclear physicist like me. The elemental structure apparent in Fig. 6, assisted perhaps by final-state interaction effects (that is where contact is made with "molecular charmonium"), may well explain the complex, and as yet incompletely resolved, structure seen experimentally in R from 3.9 to 4.2 GeV[31].

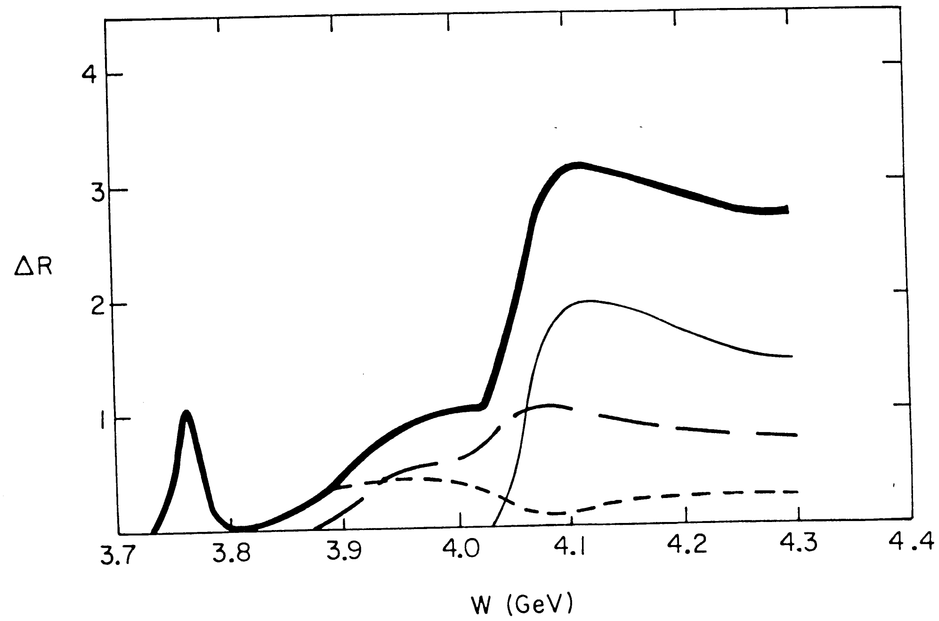


Fig. 6

5. QCD AND TOTAL WIDTHS

The discussion of the spectroscopy so far has been largely empirical, based on the non-relativistic quark model and with little dependence on the supposed underlying theory, QCD. Indeed what little input we have used (spin-dependent effects) has not proved very satisfactory at the quantitative level.

Where can the ideas of QCD be tested? One area is that of total widths, or rather widths Γ_h for annihilation of the new particles into ordinary hadrons. We recall from Section 3.1 and Fig. 1b that the assumption of annihilation at short distances leads to the hope that Γ_h can be estimated from the width for decay of the $c\bar{c}$ system into the minimum number of gluons (computed in lowest order perturbation theory). From $\Gamma_h/\Gamma_e \simeq 10$ for $\psi(3095)$ decay, one estimates $\alpha_s = 0.19$ by this approach[8]. Since $\alpha_s < 1$ at least, it is logically permissible to look for other data that will support or refute the QCD assumptions.

5.1 Pseudoscalars

The hadronic decay of a $c\bar{c}$ pseudoscalar state is estimated from its decay into two gluons. The calculation is identical, apart from a "colour"

Clebsch-Gordan coefficient, to the QED one for $^1S_0$ positronium states. In fact, with the QCD assumptions, but otherwise reliably, the ratio of PS → → hadrons to PS → γγ is[32])

$$\Gamma(PS \to \text{hadrons}) = \frac{9}{8}\left(\frac{\alpha_s}{\alpha}\right)^2 \Gamma(PS \to \gamma\gamma) \tag{6}$$

The electromagnetic (2γ) decay of the PS is closely related, in the non-relativistic $c\bar{c}$ model at least, to the decay of its vector partner, $V \to e^+e^-$. Neglecting mass and wave-function differences, one has

$$\Gamma(PS \to \gamma\gamma) \simeq \frac{4}{3}\Gamma(V \to e^+e^-) \tag{7}$$

As the partner of the ψ(3095), the X(2830) is estimated by these arguments to have a hadronic width of roughly 5 MeV. Its branching ratio into two photons is expected (somewhat more reliably) to be $B(X \to \gamma\gamma) \simeq 1.3 \times 10^{-3}$. For the χ(3454), if it is the partner of the ψ′(3685), one estimates $\Gamma_h \simeq$ ≃ 2.2 MeV.

The X(2830) has been detected so far only in the 3-photon final state by the chain, $\psi(3095) \to \gamma_1\, X(2830)$, $X(2830) \to \gamma_2\gamma_3$. The experimental product of the branching ratios is[33]) $B(\psi \to \gamma_1 X) \cdot B(X \to \gamma\gamma) = (1.2 \pm 0.5) \times 10^{-4}$. On the other hand, the inclusive photon spectrum shows no evidence of the transition, $\psi \to \gamma X$ and yields $B(\psi \to \gamma_1 X) < 1.7 \times 10^{-2}$ (90% confidence level)[23]). We can therefore conclude that $B(X \to \gamma\gamma) > 7 \times 10^{-3}$. This is five times larger than the expected 1.3×10^{-3}, but because of the experimental errors, is only two standard deviations away. There is thus no real disagreement yet. It would be desirable, nonetheless, to see independent evidence for the X(2830), especially in hadronic decay modes!

The X(2830) is perhaps a little peculiar as the pseudoscalar partner of the ψ. The χ(3454) is absolutely bizarre as the partner of the ψ′! This state is seen only as an intermediate step in the 2-photon cascade, $\psi' \to$ $\to \gamma_1\chi \to \gamma_1\gamma_2\psi$. The product of branching ratios is[34]) $B_1 \cdot B_2 = (0.8 \pm 0.4) \times$ $\times\, 10^{-2}$. Again, the inclusive photon spectrum shows no peak; $B_1 < 2.5 \times 10^{-2}$ (90% confidence level)[23]). The radiative decay $\chi(3454) \to \gamma_2\, \psi(3095)$ therefore has a branching ratio, $0.32 \pm 0.16 < B_1 < 1.0$. Such a large radiative branching ratio seems impossible to reconcile with the QCD expectations for this state. The first radiative transition should be a "favoured" M1, the second a "hindered" M1. We have already noted that $\Gamma_1 < 6$ keV is marginally within theoretical expectations (see Table 2). If we take the very unreliable estimate of $\Gamma_2 \sim 0.1$ keV at face value, the lower limit on B_2 implies that the total width of $\chi(3454)$ must be less than 0.3-0.6 keV. If we increase Γ_2 by a factor of *100* to allow for the uncertainties of relativistic corrections[22], we still find only 30 to 60 keV, compared with the expected 2 MeV! Alternatively, treat both transitions as *allowed* dipole transitions. Then $\Gamma_2/\Gamma_1 = 3(0.340/0.223)^3 = 10.7$. From the upper limit, $\Gamma_1 < 6$ keV, we have $\Gamma_2 < 64$ keV and $\Gamma_t < 200$ keV, still an order of magnitude smaller than QCD estimates. Another, probably better, estimate comes from comparison of the two "hindered" M1 transitions, $\psi' \to \gamma X$ and $\chi(3454) \to \gamma\psi$. The first has an upper limit on its width of 2 keV (Table 2). Using three times this upper limit for $\chi(3454) \to \gamma\psi$ (despite the much smaller Q value), we estimate $\Gamma_t < 18$-36 keV.

Because of the peculiar properties of $\chi(3454)$, it has been suggested[35] that the state is not $n = 2\ {}^1S_0$, but rather ${}^1D_2$ ($J^{PC} = 2^{-+}$). Given the large spin-dependent splittings, the presence of a d-state so low in energy is not totally unpalatable. The assignment as a d-state has, however, serious problems of its own[36].

5.2 P-states

Another place to check QCD expectations is the total (or hadronic) widths of the ${}^3P$ and ${}^1P$ states[37]. The 0^{++} and 2^{++} states can decay into two gluons, while the 1^{++} and 1^{+-} states must decay into at least three gluons or one gluon plus a light $q\bar{q}$ pair. Because $\alpha_s \simeq 0.2$, we have the qualitative expectation that the hadronic widths of the $J = 1^+$ states will be smaller than those of the 0^{++} and 2^{++} states. The branching ratios shown in Fig. 4 support this idea. Of the χ states, the $\chi(3508)$, assigned to be 1^{++}, has the largest radiative branching ratio of the three ${}^3P_J$ states.

The detailed QCD predictions fall into three classes of decreasing reliability. The firmest result is that $\Gamma_h(2^{++})/\Gamma_h(0^{++}) = 4/15$. Next in reliability is the ratio of the $J = 1$ state hadronic widths to that of 0^{++}:

$$\frac{\Gamma_h(1^{+\pm})}{\Gamma_h(0^{++})} \simeq N_\pm \alpha_s \ln\left(\frac{4m_c^2}{4m_c^2 - M^2}\right), \tag{8}$$

where $N_+ = 4/9\pi$ and $N_- = 10/27\pi$. Finally, least reliable because it depends on the derivative of the p-state wave function, is the actual magnitude of the hadronic width for the 0^{++} state,

$$\Gamma_h(0^{++}) = \frac{96\,\alpha_s^2}{M^4}|R'(0)|^2 \ , \tag{9}$$

where R(r) is the p-state radial wave function. With $\alpha_s = 0.19$ and an estimate of the logarithm as $\ln(1/\alpha_s^2)$ in Eq. (8), one finds the ratios of hadronic widths of the ^3P states, $\Gamma_h(0^{++}):\Gamma_h(1^{++}):\Gamma_h(2^{++}) = 15: \sim 1:4$. With the preferred wave functions of Ref. 37, one estimates $\Gamma_h(0^{++}) \simeq 2.2$ MeV. Other wave functions give somewhat smaller values.

Direct comparison of these expectations with experiment is not possible because the ^3P states are formed only via radiative transitions from the ψ'. Only branching ratios are known (see Fig. 4). The *ratios* of the hadronic widths can be plausibly deduced from experiment by assuming that the $\chi_J \to \gamma\psi$ transitions are E1 with the same transitional dipole moment[38]. Then their radiative widths should be in the ratios of their photon energies cubed. Using the radiative branching ratios of 3, 35, and 14% for the 0^{++}, 1^{++}, and 2^{++} states, respectively, from Fig. 4, and the k^3 assumption one finds 8.3 : 1 : 4.4 for comparison with 15:1:4. The agreement is quite satisfactory (the discrepancy between 8.3 and 15 is only apparent -- with errors included, $8.3 \to 8^{+\infty}_{-4}$).

For the widths themselves it is necessary to know the actual magnitudes of the radiative widths. Model calculations can be used. A somewhat better way is provided by certain dipole sum rules that provide lower as well as upper bounds[39,40]. One finds $\Gamma_h(2^{++}) \simeq 1.5$-$1.8$ MeV and $\Gamma_h(1^{++}) \simeq 0.4$-$0.8$ MeV, for example, to be compared with 0.6 MeV and 0.15 MeV, respectively, from QCD[37]. Perhaps a discrepancy of a factor of three should not be taken too seriously in view of the theoretical uncertainties in the actual values of these widths.

The present situation on tests of QCD predictions for annihilation is mixed. The indirect conclusions about Γ_h for the ^3P states seem in fair agreement with QCD. The branching ratio of $X \to \gamma\gamma$ stands in 2σ disagreement with QCD. The total width of $\chi(3454)$ appears almost certainly an order of magnitude smaller than the QCD prediction. Note that the discrepancies in absolute magnitudes for the ^3P states are that the QCD estimates are too small, while for the pseudoscalars the opposite holds true.

All of this argues for some caution in accepting QCD as the ultimate theory of hadrons just yet. Theoretical motivations and support from other areas of physics apart, QCD is "not proven" by the spectroscopy by any means.

6. CONCLUSIONS AND A GLIMPSE AT THE FUTURE (PERHAPS)

There seems little doubt that the new particles are manifestations of charm[3], within a badly broken SU(4), and that the qualitative features of the spectroscopy are described by the charmonium picture derived from QCD. Successful predictions, some from SU(4) alone and some from QCD, include:

i) Narrow resonances below the charm threshold, together with a rapid rise in R and structure just above threshold, with asymptotic approach to $\Delta R_{\text{hadronic}} = 4/3$ from above[8].

ii) Number and ordering of energy levels according to a "confined" positronium spectrum.

iii) I = 0 for all new states, at least below the charm threshold.

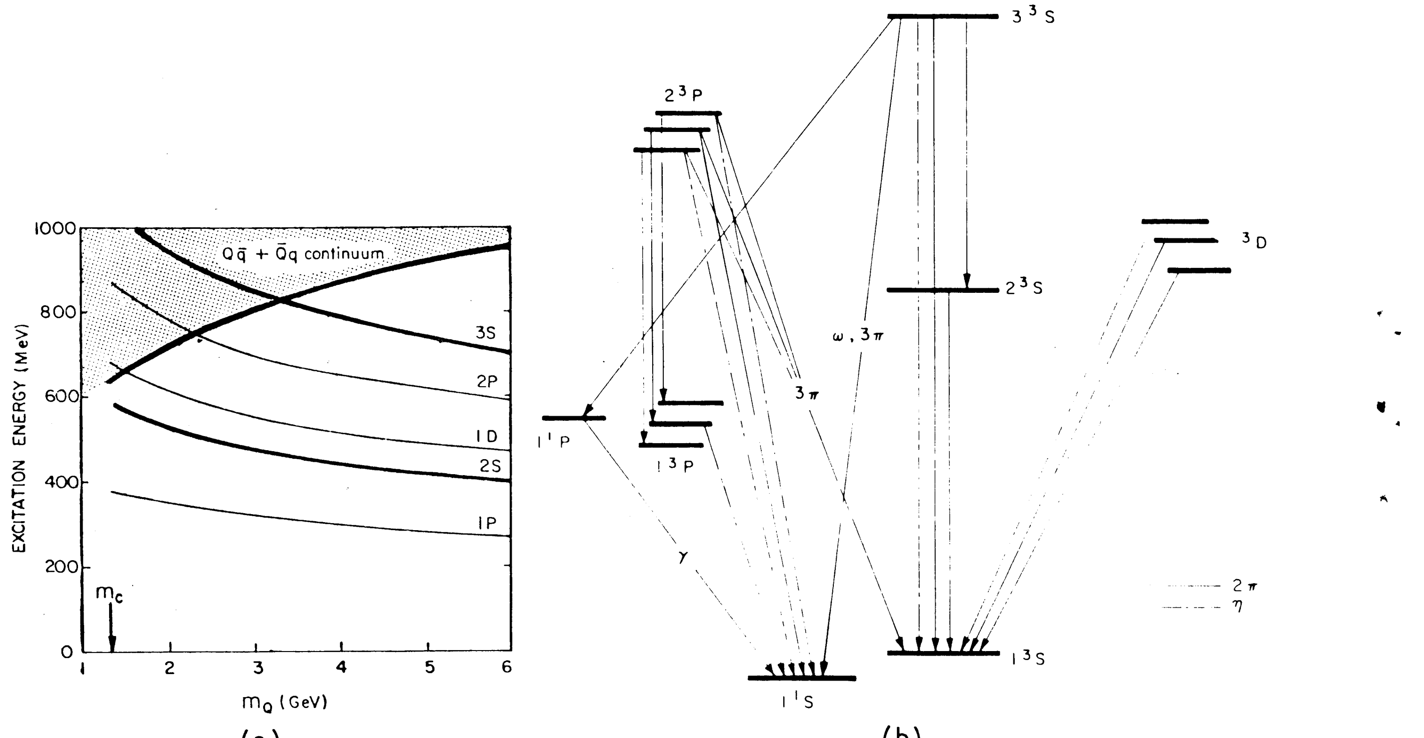
iv) Strong decay of states formed in e^+e^- annihilation above charm threshold into new $I = \frac{1}{2}$ particles (D^0, D^+; $\bar{D}^0$, D^-) the lightest of which decay by weak interactions.

[A further prediction of SU(4) is the existence of a new I = 0, charmed strange particle F^+ and its antiparticle, with a mass about 200 MeV greater than the lightest $I = \frac{1}{2}$ doublet.]

At the more quantitative level of mass differences, transition rates, and multiplet splittings, the successes still outnumber the failures, with perhaps some small cause for concern. Gross structure and spatial properties of the states are handled well enough by the models, but spin-dependent effects are given poorly or worse. The pseudoscalar states are troublesome from every angle. On the other hand, the new state at 3772 MeV must be counted a success for the theory; features of the strong (Zweig-allowed) decays of states above charm threshold are probably "understood". Significant tests of QCD are few and far between, with some hopeful evidence (Γ_h for ^3P states) and some not [especially Γ_h for $\chi(3454)$].

By any absolute standard, the new spectroscopy is truly a phenomenon! On the experimental side, newer discovery has followed new discovery in unbelievably rapid succession. On the theoretical side, the observations have been easily fitted into the charmonium picture, with no really rough edges, apart from that devil, $\chi(3454)$! If charmonium works so well at W = 3-4 GeV, what can we expect if there are still heavier quarks of different flavour, call it "beauty", with masses of, say, 5 GeV?

If the confining potential has a universal character, independent of quark mass, the answer has been given by Eichten and Gottfried[41] and displayed in Fig. 7. The greater mass causes the level spacings to decrease, as $m_Q^{-1/3}$ for a purely linear potential and empirically as $m_Q^{-0.26}$ in Fig. 7a. At the same time it is expected that the gap between the ground state and the Zweig-allowed continuum increases, as shown in Fig. 7a. At m_Q = 5 GeV there are three bound s-states and probably two bound d-states. For such a hypothetical new quark, in e^+e^- annihilation between 10 and 11 GeV there would thus be seen three ψ-like narrow states, plus two states akin to the $\psi''(3772)$ with small Γ_e, but narrow because below the "beauty" threshold. Since two p-states would be bound as well, there would be a rich spectroscopy of radiative and hadronic transitions among the levels. Figure 7b indicates the hadronic transitions expected. Radiative rates will scale crudely as $m_Q^{-5/3}\, e_Q^2$ for transitions between corresponding states. For example, if $e_Q = 1/3$, the $\psi' \to \gamma\, \chi_J$ transitions will have analogues in the beautiful spectra with partial widths summing to 1.5-2.0 keV, compared to the 50 keV in charmonium. Since fine structure decreases rapidly with mass (as $m_Q^{-5/3}$ for a linear potential), multiplet splittings will be much smaller than in charmonium and will be difficult to resolve. Hadronic (annihilation) rates for the s-states decrease somewhat because of a decrease in α_s with mass, not compensated completely by an increase in $|\psi(0)|^2$.



(a) (b)

Fig. 7

Speculations of this sort may be more than an amusing exercise in quantum mechanics. Lederman, these Proceedings, has reported a beautiful new peak at 9.6 GeV $\mu^+\mu^-$ invariant mass in $p + (Z, A) \rightarrow \mu^+\mu^- X$. The peak is rather lop-sided with an observed width of $\sim$ 1.3 GeV and experimental mass resolution of $\sim$ 0.5 GeV. Perhaps we are on the threshold of a whole beautiful new "new particle spectroscopy"[42]!

FOOTNOTES AND REFERENCES

1) For detailed discussion of the experiments and deductions, see G.J. Feldman, ψ spectroscopy, Proc. 1976 Summer Institute on Particle Physics, Stanford Linear Accelerator Center (ed. M.C. Zipf), SLAC Report No. 198 (November 1976), p. 81; B.W. Wiik and G. Wolf, Electron-positron interactions, DESY 77/01 (1977), (Les Houches Summer School, 1976); G.J. Feldman and M. Perl, Physics Reports C (to be published).

2) P.A. Rapidis et al., Phys. Rev. Letters 39, XXX (1977).
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3) B.J. Bjorken and S.L. Glashow, Phys. Letters 11, 225 (1964).

4) Because there are no asymptotic single-quark states in a theory with confinement, the mass of a quark is an ambiguous concept. We take it as a parameter to be determined empirically.

5) Even though the real world is "colourless", the colour degree of freedom manifests itself in various ways, several of them quantitative, where factors of three or nine appear, giving agreement with experiment. For a simple survey, see J.D. Jackson, Lectures on the new particles, Proc. 1976 Summer Institute on Particle Physics, Stanford Linear Accelerator Center (ed. M.C. Zipf), SLAC Report No. 198 (November 1976), p. 147.

6) S. Okubo, Phys. Letters 5, 165 (1963).
G. Zweig, CERN TH 401, 412 (1964).
J. Iizuka, Progr. Theor. Phys. Suppl. 37-38, 21 (1966).

7) The inhibition hinges on the quarkless intermediate state (Fig. 1b) consisting of a small number of "hard" gluons (with consequent weak couplings). Why not a large number of "soft" gluons (each with strong coupling)? One appeals to short distances for the annihilation, but the arguments are far from foolproof.

8) T. Appelquist and H.D. Politzer, Phys. Rev. Letters 34, 43 (1975); Phys. Rev. D 12, 1404 (1975).

9) A. Martin, Phys. Letters 67B, 330 (1977).
H. Grosse, Phys. Letters 68B, 343 (1977).

10) A. Martin, CERN preprint TH 2338 (June 1977).

11) K. Gottfried, *in* Proc. Canadian Inst. Particle Phys. Internat. Summer School, McGill University, 21-26 June 1976 (eds. R. Henzi and B. Margolis) (McGill University, Montreal, 1977).

12) The ninth, as yet unseen, state is $^1D_2$ ($J^{PC} = 2^{-+}$), part of the 1d multiplets. It may be hard to find because its neighbour, the new 1^{--} state at 3.772 GeV (presumably $^3D_1 + \varepsilon\ ^3S_1$, at least in the $c\bar{c}$ sector) decays strongly into $D\bar{D}$, as will higher lying states that would otherwise feed the $^1D_2$ radiatively. But see Section 5.1.

13) J. Pumplin, W. Repko and A. Sato, Phys. Rev. Letters 35, 1538 (1975). H.J. Schnitzer, Phys. Rev. Letters 35, 1540 (1975).

14) W.H. Furry, Phys. Rev. 50, 784 (L) (1936).

15) For $\chi(3552)$ a spin value of $J = 4$ or higher seems quite implausible. With $J = 0$ excluded, $J = 2$ seems inevitable. Two $J = 2$ states near in mass in the $c\bar{c}$ system seems rather unlikely. Thus $J = 0$ for $\chi(3415)$ appears inevitable, too. But see Ref. 35.

16) The spins can be established unambiguously by observation of angular correlations in the decay chain, $\psi' \to \gamma_1\chi \to \gamma_1\gamma_2\psi \to \gamma_1\gamma_2(\mu^+\mu^-)$. See G. Karl, S. Meshkov and J.L. Rosner, Phys. Rev. D 13, 1203 (1976); P.K. Kabir and A.J.G. Hey, Phys. Rev. D 13, 3161 (1976); H.B. Thacker and P. Hoyer, Nuclear Phys. B106, 147 (1976).

17) Perhaps the most detailed calculations, both above and below the charm threshold, are those of the Cornell group: E. Eichten et al., Phys. Rev. Letters 34, 369 (1975); Phys. Rev. Letters 36, 500 (1976); K.D. Lane and E. Eichten, Phys. Rev. Letters 37, 477 (1976); E. Eichten et al. (in preparation).

18) The E1 matrix elements are in rough accord with experiment, as are the values of Γ_e, only if $e_Q = 2/3$. This can be taken as (weak) evidence for the charges of the c quarks, independently of the step in R at the charm threshold.

19) H.J. Schnitzer, Phys. Rev. D 13, 74 (1976).

20) R. Barbieri, R. Kögerler, Z. Kunszt and R. Gatto, Nuclear Phys. B105, 125 (1976).

21) A.B. Henriques, B.H. Kellet and R.G. Moorhouse, Phys. Letters 64B, 85 (1976).

22) G. Feinberg and J. Sucher, Phys. Rev. Letters 35, 1740 (1975).

23) C.J. Biddick et al., Phys. Rev. Letters 38, 1324 (1977).

24) B.J. Edwards and A.N. Kamal, Phys. Rev. Letters 39, 66 (1977).

25) Equation (4) should be viewed as symbolic. At the level of M^{-2} in the amplitude there are relativistic corrections to the wave function that are not displayed.

26) Second and third papers of Ref. 17.

27) For some mysterious reason, everyone finds about 3.77 GeV for the mass -- our spinless potential of Section 4.2 puts the 1d state 84 MeV above the 2s; with relativistic corrections, Ref. 20 gets 3.78 GeV; the spinless, but coupled, channel calculation of the Cornell group gives 3.77 GeV; the relativistic calculation of Ref. 21 gives 3.77 GeV for two different sets of parameters! No great significance should be attached to any of these coincidences.

28) A. De Rújula, H. Georgi and S.L. Glashow, Phys. Rev. Letters 37, 398 and 785 (1976).

29) Lane and Eichten, Ref. 17.

30) L.B. Okun' and M. Voloshin, Zh. Eksper. Teor. Fiz. 23, 369 (1976); A. De Rújula et al., Phys. Rev. Letters 38, 317 (1977).

31) See, for example, Fig. 16 of R.F. Schwitters, Proc. 18th Internat. Conf. on High Energy Physics, Tbilisi, July 1976 (JINR, Dubna, 1977), Vol. 2, p. B34.

32) T. Appelquist et al., Phys. Rev. Letters 34, 365 (1975).

33) W. Braunschweig et al., Phys. Letters 67B, 243 (1977).

34) J.S. Whitaker et al., Phys. Rev. Letters 37, 1596 (1976). Recent confirmation of $\chi(3454)$ has come from the Pluto group at DESY (U. Timm, these Proceedings).

35) H. Harari, Phys. Letters 64B, 469 (1976).

36) J.D. Jackson, CERN preprint TH 2305 (April 1977, to appear in Proc. 12th Rencontre de Moriond).

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38) M.S. Chanowitz and F.J. Gilman, Phys. Letters 63B, 178 (1976).

39) J.D. Jackson, Phys. Rev. Letters 37, 1107 (1976).

40) Sum rules of a different sort (relativistic dispersion relations for $\gamma \to$ hadrons, $\gamma\gamma \to$ hadrons, etc.) are employed by V.A. Novikov et al. [Phys. Rev. Letters 38, 626, 791 (E) (1977); Phys. Letters 67B, 409 (1977); A.I. Vainshtein, these Proceedings] to obtain relations among and values of various radiative and hadronic widths. This approach complements the more detailed potential models and yields closely similar results in areas of overlap.

41) E. Eichten and K. Gottfried, Phys. Letters 66B, 286 (1977).

42) Lest the reader thinks I am overly satisfied, complacent, and confident about the basic truth of asymptotically-free QCD, I remind us all that we do not really understand Zweig's rule and/or the use of lowest-order perturbation theory for annihilation rates, or quark confinement. Our apparently successful spectroscopy rests on shaky foundations indeed.