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A MODEL WITH UNCONFINED ψ CONSTITUENTS*

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ABSTRACT

We propose a model for the ψ particles in which these objects are made of unconfined constituents. In our picture no analog of charmed particles exist and no Okubo-Zweig-Lizuka rule is needed.

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I. INTRODUCTION

ψ spectroscopy indicates that the new particles[1,2] are composed of spin $\frac{1}{2}$, charged and isosinglet fermions and anti-fermions. If these are charmed (c) quarks which together with the ordinary (q) quarks form the basic SU(4) multiplet then the existence of the new particles has been predicted ahead[3,4] and their finding is a major theoretical triumph.

However, to verify this we have to discover charmed particles ($c\bar{q}$ composites) and the appropriate SU(4) pattern, in particular for the weak GIM[5] currents, should emerge. No evidence for charmed bosons with mass $m \approx$ 1.6-2 GeV (the value suggested both by SU(4) and the threshold in $R = \sigma(e^+e^- \to \text{"hadrons"})/\sigma(e^+e^- \to \mu^+\mu^-)$) has been found to date.[1] Also no really satisfactory understanding of the almost exact Okubo-Zweig-Lizuka rule for ψ's has emerged.

We propose to consider the possibility that the ψ constituents do not belong to the quark family, do not participate in the standard colored gauge interactions and, in particular, are unconfined.[6] The L's must possess a new kind of strong interaction among themselves but conceivably have significantly weaker interactions with quarks and ordinary hadrons.[7]

In such a picture no analog of the charmed particles exists ($L\bar{q}$ states need not bind and being colored are even forbidden as a physical state.) No Okubo-Zweig-Lizuka rule need be invoked since the interactions responsible for ψ decay, ψ binding and hadron binding are all independent.

We note that the knowledge of the few low lying states in the ψ spectrum does not require an infinite confining potential and many other models can fit just this experimental information.[7]

We suggest, in particular, that the so-called newly discovered "heavy leptons"[8] are the unconfined constituents, and to avoid contradictory notation we call them L's.[9]

We find the suggestion that the "heavy lepton" threshold fortuitously opened up at W = 4 GeV within a few percent of the charm threshold and so helps to maintain the charm hypothesis[2] perfectly possible but contrived. A simpler, neater possibility is to abolish the charm hypothesis altogether and make the L's the sole moving force behind all the structures in e^+e^- cross section, narrow spikes and broad thresholds alike.

In our scheme we have to specify three new types of interactions:

(i) H_W — a "weak" Hamiltonian responsible for the L decays $L \to \nu_L + \ell + \nu_\ell$ and $L \to \nu_L$ + hadrons where ν_L is the associated "neutrino" and ℓ is either e or μ.

(ii) H_B — the interaction specific to the L world and which is responsible for the $\bar{L}L$ binding to form ψ, χ, η_c states and the rest of the broad bumps beyond $\bar{L}L$ thresholds.

(iii) H_H — an interaction responsible for the $\bar{L}L \to \bar{q}q$ transitions (e.g. $\psi \to$ hadrons and $\psi' \to \psi$ + hadrons).

This means in particular that we have to specify the nature of the "ν_L", to indicate whether we have one L or more, and to specify the number, spins, masses and couplings of the new bosons responsible for the H_W, H_H, and H_B interactions. As we will see, the available experimental data strongly restrict this richness.

II. THE WEAK INTERACTION OF L's

In the standard model for the L's which interprets them as heavy leptons

$$H_W = g\bar{L}(1-\gamma_5)\not{W}\,\nu_L \qquad g^2/m_W^2 \approx G_F \approx 10^{-5}\,\mathrm{GeV}^{-2} \tag{1}$$

where ν_L, the analog of ν_e, ν_μ, participates then only in weak interactions. (If ν_L participated in the strong L interactions, $\psi \to \nu_L \bar{\nu}_L$ would be a dominant decay mode.) A priori L could carry e- or μ-number and we would replace ν_L in Eq. (1) by ν_e or ν_μ. This alternative will be discussed later.

Equation (1) yields

$$\Gamma(L \to \nu_L + \ell + \bar{\nu}_\ell) \sim 10^{12}\,\mathrm{sec}^{-1} \tag{2}$$

which is consistent with the $e\mu$ events detection distance.

Also, using the W boson coupling to quarks, one can estimate the partial width of semileptonic decays (which is proportional to the R value at the relevant Q^2)[10]

$$L \to \nu_L + \text{hadrons} \qquad (m^2_{\text{hadronic}} = Q^2 \approx 1\text{-}2\ \mathrm{GeV}^2) \tag{3}$$

so that

$$\Gamma(L^\pm \to \nu_L \ell^\pm \nu_\ell)/\Gamma(\mathrm{tot}) \equiv r \simeq 0.17 \quad (\text{for each } \ell) \tag{4}$$

and also specific decay channels

$$\Gamma(L^\pm \to \rho^\pm \nu_L)/\Gamma(\mathrm{tot}) \simeq 0.3 \tag{5}$$

$$\Gamma(L^\pm \to \pi^\pm \nu_L)/\Gamma(\mathrm{tot}) \simeq 0.1 \quad . \tag{6}$$

We have considered in Eqs. (1)-(6) only the simplest possible interaction scheme.

A crucial test for our scheme is the consistency of $\sigma(e^+e^- \to e^\pm\mu^\mp + \text{neutrals})$ with the rise in R. The latter results from

$$e^+e^- \rightarrow \text{resonances} \rightarrow L\bar{L} \tag{7}$$

or

$$e^+e^- \rightarrow L\bar{L} + nB \tag{8}$$

if the mediators of H_B interactions have a small mass (m_B) and $W \geq 2m_L + nm_B$.

The following difficulties may arise:

(a) Our scheme predicts

$$z = \frac{\sigma(\text{observed; } e^+e^- \rightarrow e^{\mp}\mu^{\pm} + \text{neutrals})}{R - 2.5} = 2r^2 x \tag{9}$$

where $x \approx 0.1$ [8] is the acceptance for the $e\mu$ events and following Gilman[11] we take $R = 2.5$ to represent the "old physics" value of R.

Using r from Eq. (4), z is predicted to be $\approx 5 \cdot 10^{-3}$ whereas the experimental value varies and the worst discrepancy occurs at $W \simeq 4.2$ GeV where $R \approx 5$-6 and $z \approx 10^{-3}$. A small suppression by about a factor of two in the purely leptonic versus the semileptonic amplitude may correct this discrepancy.

(b) The $e\mu$ signal seems to peak later than R.[12] Let us assume that there are (at least) two lumotons L_1 and L_2 with $m_{L_1} \approx 1.7$-2 GeV, $m_{L_2} \approx 2.3$-2.6 GeV. The prominent structure in R at $W \approx 4.2$ GeV involves decays into $L_1\bar{L}_1$ only. If the pure leptonic branching ratio of L_2 exceeds considerably that of L_1 then the delayed $e\mu$ signal reflects the delayed onset of $L_2\bar{L}_2$ threshold. The assumption of two or more L's is also helpful in understanding the high asymptotic R value and in ψ-spectroscopy.

(c) Inclusive spectra. Since an extra ν_L is emitted in all L decays our scheme helps to understand the qualitative features of small average $y \equiv E_{ch}/W$, the fraction of energy of charged tracks. Naive statistical arguments suggest equipartition, i.e. $y \approx 0.67$ whereas experimentally y appears to decrease from y = 0.6 for low W to y = 0.5 at $W \approx 8$ GeV.[1]

The non-observation of a sharp decrease in y at the new R threshold is a difficulty both in our model and, as we show in Appendix A, also in standard charm or multi-flavor models. If m_B is large then reaction (7) dominates (8). Since the L decay is energy (W) independent and leads to few, and typically one, charged particles we expect:

(i) A slower rise in $n_{charged}$ beyond $W \approx 4$ GeV which is consistent with the data.

(ii) Contributions due to the 2 body decay channels (5) and (6) to the inclusive charged particle distribution s $d\sigma/dx$ at $\frac{1}{2} \leq x \leq 1$. This is further discussed in Appendix B. These effects are not inconsistent with the present data.

The L's may manifest also in various anomalies in ν_μ induced reactions. Since these effects are much more model dependent, we restrict ourselves to a few comments. Experimentally the reactions $\nu_\mu \to \mu^+$ + anything are suppressed by a factor ~ 100-1000 relative to the "normal" process $\nu_\mu \to \mu^-$ + anything.[13] If we take L to be a lepton as suggested by various gauge models (so that e.g. μ^-, ν_μ, M^o, M^+ and e^-, ν_e, E^o, E^+ both form one multiplet) then $\nu_\mu \to M^+ \to \mu^+\nu\nu$ or $M^o \to \mu^-\mu^+\nu$ would be a possible source for such wrong sign μ's. This possibility seems to be ruled out experimentally.[13]

If L^- is a lepton with a μ number then $\nu_\mu \to L^- \to \mu^-\nu\nu$ will be a source of lower energy μ's and apparent violation of scaling.

For ν_μ induced reactions it is of little importance if the L neutrino is ν_e or a new neutrino ν_L. If however we adopt the present scheme in which the L's have moderately strong interactions with hadrons then one may attempt to account for $\mu^- \mu^+$ pair production via diagrams like those in Fig. 1.[14]

We have worried above only about phenomenological consequences of the weak L decay. If the ν_L has no strong H_B interaction then the weak vertex $L \rightarrow W_\mu + \nu_L$ does not conserve (for vector B) the L charge so that special care has to be taken to insure renormalizability of the model. Note however that in many schemes only the overall strong + weak Lagrangian is renormalizable (due to e.g. cancellation of anomalies in quark and lepton triangles).

Also in Appendix D we suggest a scheme where the ν_L do have strong interaction.

III. ψ SPECTROSCOPY

In this section we discuss the binding of $\bar{L}L$ pairs via the H_B interactions to form the ψ - like states.

Since we do not have to worry about the admixture of light quarks and the B particles (the mediators of the B interactions) will turn out to be massive ($m_B \geq 0.7$ GeV and in most schemes $m_B \geq 3$ GeV) the description in terms of just $\bar{L}L$ bound states may be adequate. Furthermore, if the potential generated via the exchange of the B's is smoothly varying over the ψ, ψ' spatial region then a simple non-relativistic Schrodinger equation approach may be attempted.

This is very similar to the simple charmonium model with the difference that we do not confine the constituents by infinitely rising potentials. The absolute binding $B(1S) = (2m_{L_1} - m_\psi)$ etc. should also be calculated and compared with the experimental values.

The charmonium model fits imply that in order to account for the level ordering and the slow decline of the wave functions at the origin for the radial excitations $\psi'(3.7)$, $\psi''(4.2)$, ... (as inferred from leptonic widths), potentials which deviate considerably from simple Yukawa's are required. In particular, an opposite (positive) curvature over the relevant r range seems to be required.

Non-confining simple potentials of the variety which lead to a reasonable explanation of the whole spectrum by excitations of one $\bar{L}L$ pair have indeed been found.[7,24]

In our version we have an extra degree of freedom, namely choosing the $L_1\bar{L}_1 - L_2\bar{L}_2$ mixtures of the states. We could thus start from separate pure

$L_1\bar{L}_1$ and $L_2\bar{L}_2$ levels shifted by $2(m_{L_2} - m_{L_1})$ and via mutual mixing and two level repulsion obtain the experimental structure as indicated (see Fig. 2). The choice that the $\psi'(3.7)$ is the ground state of $L_2\bar{L}_2$ essentially helps to explain its large leptonic width and could, if χ, P(3.5) are $L_1\bar{L}_1$ states, explain the small e.m. transitions $\psi' \to \chi\gamma$, etc.[24]

The interpretation of the potential in terms of the t-channel exchanges entailed some difficulties because of the following conflicting requirements:

(i) The range of the potential has to be large ($\geq \frac{1}{4}$ Fermi) to get the small $(E_{2P} - E_{1S}) \sim 0.35$ GeV splitting.

(ii) The masses of the exchanged B's have to be large to prevent fast decays of ψ particles.

Clearly no such difficulty arises in the orthodox approach where identical potentials for ψ and ordinary hadron binding are generated via the exchange of the same confined massless gluons.

To elaborate on (i) we note that for Abelian interactions one particle (vector or scalar) exchanges always yield attractive forces

$$V(r) = - \sum_i g_i^2 \frac{e^{-\mu_i r}}{r} \tag{10}$$

For such potentials (and also for the more general case $- \int \sigma(\mu)\, e^{-\mu r}/r \, d\mu)$[15] one can prove that if a P-wave bound state exists then $B_{1S} - B_{2P}$ is bigger than a minimum value obtained for $V = -g^2 e^{-\mu_0 r}/r$ with g^2 fitted to the P wave binding B_{2P}. For the particular choice $B_{2P} \approx 0$, i.e. $m_{2P} \approx 3.4$ GeV, $m_{L_1} \approx 1.7$ GeV the numerical estimate of min $(B_{1S} - B_{2P})$ is sufficiently simple (Appendix C) and $B_{1S} - B_{2P} \approx 0.3 - 0.4$ GeV implies $\mu_0 \leq 0.5$-1 GeV.

A vector B_V of such a mass (or any mass $\leq 2m_{L_1}$) would be produced in e^+e^- collisions

$$e^+e^- \to "\gamma" \to "\bar{L}L" \to B_V \tag{11}$$

and then decay slowly via the weak L quark interaction by $B_V \to "\bar{L}L" \to \bar{q}q$.[16] Since no significant narrow spikes in $\sigma(e^+e^- \to \text{hadrons})$ were observed below the ψ and between ψ, ψ' a light B_V is unlikely. (Also $\eta_c \to 2B_V$ should be avoided.)

A scalar B[17] should be heavier than $m_{\psi'} - m_\psi = 0.6$ GeV so as not to show up as a narrow resonance in the $\pi\pi$ invariant mass in the $\psi' \to \psi\pi\pi$ decay. If B_s is an exact isoscalar then ψ, ψ', and ψ'' cannot decay strongly via H_B into B's due to charge conjugation nor would 1^+ (P wave) and 0^- ("η_c") states because of the Bose symmetry requirements. A model with light B's has some attractive features outside the realm of just ψ spectroscopy itself. For $W \geq 2m_L + m_B$ we have also $e^+e^- \to \bar{L}L + B_s$, $B_s \to \pi\pi$ and not only $e^+e^- \to \bar{L}L$ which might be helpful in understanding $s\,d\sigma/dx$ at small x. Also multi-B production may account for the quick rise in $\sigma_{tot}(\psi\,p)$ as inferred from ψ photo-production. Finally if $m_B \approx 0.7$ GeV it could help to explain the strong peaking of $d\sigma/dm_{\pi\pi}$ for the $\psi' \to \psi\pi\pi$ decay towards $m_{\pi\pi} = m_{\psi'} - m_\psi$.

The following difficulties do occur however in the light B_s models.

(i) The natural parity χ state observed in $\psi \to \gamma + \chi$, $\chi \to \pi\pi$ could decay $\chi \to nB_s$ inducing a large width.

(ii) $\psi \to \gamma + nB_s$ may be too strong.

(iii) An $L_1\bar{L}_1 - L_2\bar{L}_2$ mixing has to proceed via (virtual) annihilation into B's and thus cannot occur in the 1^{--}, 0^{-+} and 1^{++} states. If ψ' is mainly

$L_2\bar{L}_2$ and the P-wave states $L_1\bar{L}_1$ then mixing is required to account for ψ' photonic transitions. In general, ruling out mixing decreases our ability to obtain a good level ordering.

The difficulty with potentials like (10) can be avoided without introducing low mass B's if we allow for non-Abelian exchanges.

As an example consider the ideal case of exactly degenerate L_1 and L_2 states forming an "L-spin" doublet. The difference between the shapes of the R-2.5 and $e\mu$ signals if indeed a real effect would then require (see discussion of Section II) a third L. The bound states will now be eigenstates of L-spin with L = 0/L = 1: $L_1\bar{L}_1 \pm L_2\bar{L}_2$. Let us assume that also in the t channel we have exchanges of L-spin = 0 and L-spin = 1, i.e. Abelian and non-Abelian interactions. We will then have the following potentials in $L_1\bar{L}_1 + L_2\bar{L}_2$ and $L_1\bar{L}_1 - L_2\bar{L}_2$ states, respectively:

$$\begin{aligned} V^{(+)} &= V^{(0)} + 3V^{(1)} \\ V^{(-)} &= V^{(0)} - V^{(1)} \end{aligned} \tag{12}$$

where $V^{(i)}$, the potentials due to the exchange with L = 0 and L = 1, respectively, are each of the purely attractive Yukawa form.

Thus the ψ spectroscopy splits into two pieces, the more strongly bound L-spin = 0 systems and the more loosely bound L-spin = 1 systems.

To make a concrete choice let us assume that L_1 and L_2 have opposite charges, i.e. $Q = 2L_3$ and the photon is L-spin isovector. All states produced via one photon, ψ, ψ', $\psi''(4.2)$, $\psi'''(4.4)$, etc., are therefore L = 1 states and $\chi(3.4)$, P(3.5), "η_c" (2.8) and in general any state obtained from the ψ family via a photonic transition has L = 0. This is so because three objects with L = 1 and $L_3 = 0$, e.g. ψ', $\chi_{L=1}$, γ, cannot couple.

Such a model has several nice points:

(a) Assuming ordinary hadrons have L = 0, the $\psi \rightarrow$ hadrons decay violates L-spin and hence, if analogy with I-spin is any clue, may be weaker by order α than the "strong" H_B pure $\bar{L}L$ interactions. The hadronic decay of the L = 0 (χ, P, η_c) states could still be somewhat dynamically suppressed because $H_B > H_H$ but may have widths of the order of MeV's.

(b) Since the L-spin = 0 states are much more strongly bound, the fact that P-wave (L-spin = 0) states occur below the 2S ψ' (L-spin = 1) state is qualitatively explained. Also the large ($\sim$ 300 MeV) $\psi - \eta_c$ splitting which is very difficult to obtain in ordinary charm models need no longer be a concern since the η_c is L-spin = 0 and more strongly bound. In fact the $\eta_c(2.8)$ could even be a radially excited state.

(c) Since we are not restricted to small mass exchanges we can attempt to realize a bootstrap scheme where the ψ, ψ' etc. exchanges will generate the binding potentials. Note that the more attractive $V^{(+)}$ potential will cause Abelian interaction of longer range than the corresponding non-Abelian interaction.

(d) In the present model, the ψ', ψ'' etc. are all excitations of the same ground state so that the problem of similar values of the wave function at the origin is relevant and could not be resolved if the potential binding these states is a pure attractive Yukawa type, Eq. (10). Note however that precisely for this case of L-spin = 1 the binding potential is a superposition of attractive and repulsive potentials with the attractive ranges being systematically longer. (see comment (c) above.) It is relatively straightforward to construct by such superposition Wood-Saxon type potentials which lead to reasonable predictions for the wave function at the origin.

(e) A severe difficulty in many of the conventional schemes, namely the small electric and magnetic dipole transitions $\psi' \to \gamma P$ and $\psi \to \gamma + \eta_c$,[4] respectively, may be avoided. Even though the energies of these levels happen to be relatively close, these wave functions are in qualitatively different potentials and in particular the ψ, ψ' are likely to be much more spread out in r than the corresponding L-spin = 0 states. The overlap matrix elements are therefore likely to be small.

Finally, in the framework of this model there is room for much stronger bound, lower lying states of the L-spin = 0 family, whereas the discovery of such lighter narrow states would be detrimental to the ordinary charmonium and heavier quark models. Note also that the L-spin doublet model implies doubly charged ψ-like states.

We would like, however, to emphasize the independence of this specific L-spin model and the much more general concept of unifying the $e\mu$-anomaly, the rise in R and ψ physics, which is the main theme of our paper.

The motivation for presenting this model was to show that the difficulties with the 2P-1S small splittings and large exchanged masses are perhaps not insurmountable.

Another non-Abelian model, somewhat less appealing is presented in Appendix D.

IV. ψ - HADRON INTERACTIONS

There are two types of ψ - hadron interactions: (i) $\psi(\bar{L}L)$ annihilations into hadrons, e.g. $\psi \to 3\pi$, and (ii) ψ bilinear processes, e.g. $\psi' \to \psi\pi\pi$ and $\psi p \to \psi p$. The calculation of both types of interactions in a simple, reasonable way is a great challenge in all approaches. Experimentally the total widths of the lowest ψ states are $\sim 10^{-3}$ smaller than typical hadronic widths. Exclusively we have, for example, $\Gamma(\psi \to \rho\pi)/\Gamma(\rho \to \pi\pi) \sim 10^{-5}$. The bilinear decay $\psi' \to \psi\pi\pi$ is much faster and $\sigma_{tot}(\psi p) \approx 1$ mb at high energies ($s \approx 200\ \mathrm{GeV}^2$) is the value derived from coherent ψ photoproduction by using VDM. The last value of $\sigma_{tot}(\psi p)$ is an order of magnitude smaller than ordinary meson proton cross sections. In the following we take the point of view that this suppression indicates that the quark- L interactions are smaller than ordinary quark-quark interactions and the L's do not belong to the quark family. We do bear in mind, however, the possibility that this last suppression may be explained by some dynamical or kinematical mechanism related to the large charmed quark mass. Thus $\sigma_{tot}(\psi p) \approx 1$ mb may in fact be the best single piece of evidence for the hadronic nature of the ψ's.

To explain the ψ decays we introduce a new vector (1^{--}) boson H which couples to both quarks and L's and for simplicity will be taken to be a singlet of both quark and L symmetry groups. Since we want to keep vector couplings universal we do not identify H with B, the vector particle responsible for $L\bar{L}$ binding. The relevant interaction Hamiltonian is

$$H_H = g\bar{L}\not{H}L + g\bar{q}\not{H}q \tag{13}$$

and the hadronic ψ decay proceeds then via the diagram of Fig. (3.a). This can be compared to the leptonic annihilation diagram Fig. (3.b) which involves the same $\psi(0)$ (wave function at the origin) factor and hence

$$\frac{e^4/m_\psi^4}{g^4/m_H^4} = \frac{\Gamma_{\psi \to e\bar{e}}}{\Gamma_{\psi \to \mathrm{had}}} \approx \frac{1}{10} \tag{14}$$

and

$$\frac{g^2}{m_H^2} \approx (2\text{-}3) \times 10^{-3}\,\mathrm{GeV}^{-2} \; ; \tag{15}$$

we assumed above $m_H >> m_\psi$. In principle a small m_H need not be ruled out since the small coupling to $\bar{L}L$ will reduce its effect on $\sigma_{tot}(e^+e^- \to$ "hadrons") and on ψ particle decays (e.g. $\psi' \to H + \eta_c$). However, unless $2m_H \geq m_{\eta_c}$, $\eta_c \to 2H_V$ could be the dominant decay of η_c. For the other extreme case $m_H^2 << m_\psi^2$, Eq. (15) is replaced by

$$\frac{e^2/m_\psi^2}{g^2/m_\psi^2} \approx \tfrac{1}{4} \; ; \quad \text{i.e. } g^2 \approx 4e^2 \sim \tfrac{1}{2} \tag{16}$$

not a very small coupling.

Note that Eq. (15) or (16) is our substitute for the Okubo-Zweig-Lizuka rule. It is simply built in by choice of sufficiently small coupling constant[18] (and/or large m_H).

The ratio $\Gamma_{\mathrm{hadronic}}/\Gamma_{\mathrm{leptonic}}$ is expected to be the same for ψ and ψ'. This yields $\Gamma_{\psi' \to \mathrm{hadrons}} \approx 20\text{–}30$ keV in accord with the lack of strong branching ratio to identified final hadronic states in ψ' decay.

Charge conjugation forces ψ bilinear processes to occur in second order in H_H of Eq. (13) with very small widths or cross sections in contradiction with the data.[19] We are therefore led to couple also a scalar meson H_s to both quarks and L's

$$H_{H_s} = g_s^L \bar{L} L H_s + g_s^q \bar{q} q H_s \quad . \tag{17}$$

For the scalar coupling there is no universality and $g_s^L >> g_s^q$ is possible.[20] The scalar H_s could then be identified with the light scalar B_s (of mass $m \approx 0.7$-1 GeV) introduced in our discussion of ψ spectroscopy.

A large effective coupling $g^2_{\psi' \psi (\pi\pi)}/4\pi \approx 0.4$ is required for ψ decay and hence $g_s^q g_s^L \approx 1$ also. Crossing over to the $\psi\pi \to \psi\pi$ or $\psi p \to \psi p$ scattering region, we find that $\sigma^{tot}(\psi p \to \psi p) \approx 30\ \mu b$ at $s = 20\ GeV^2$ — the value indicated by VDM and ψ photoproduction.

The single scalar exchange diagram cannot, however, explain the sharp rise of $\sigma_{tot}(\psi p)$ to ≈ 1 mb at $s \approx 200\ GeV^2$. One possible way is to associate this rise with inelastic final states of the type

$$\psi p \to \begin{cases} \bar{L}L + p \\ \bar{L}L + nB_s + p \end{cases} \quad .$$

A multiperipheral structure analogous to that of hadronic collisions could arise particularly if we have strong trilinear B_s coupling. No new $L\bar{q}$ (analogs of charmed $c\bar{q}$ composite) states need be invoked to explain the rising $\sigma_{tot}(\psi p)$.

The prediction of a significant $e\mu$ signal in γp collisions is a common feature of any L model. This prediction is independent of the details of the final state in

$$\gamma p \to L\bar{L}p \tag{18}$$

or

$$\gamma p \rightarrow L\bar{L} + nB_s + p \ . \qquad (19)$$

The Bethe-Heitler cross section for reaction (18) has been calculated to be $\sigma_{B.H.} \approx 0.3$ nb at $E_\gamma = 100$ GeV for $m_L = 2$ GeV.[21]

Since ψp scattering is likely to be largely inelastic we expect in our scheme

$$\sigma_{\gamma p \rightarrow \bar{L}L + X} > \sigma_{\gamma p \rightarrow \psi p}$$

and the latter is $\approx 0.5\ \mu$b at $E_\gamma = 100$ GeV.

V. COMPARISON TO CHARM

The detailed predictions of the charmonium scheme which cover extensively all facets of ψ spectroscopy (ψ ψ' χ P_c η_c ... levels and transitions $\psi' \to x+\gamma, \ldots$) can be regained in nonconfining models of the type discussed here.[7.24]

We believe that the sucess of the charmonium picutre rests mainly on the following key features:

(i) "Large" $|\psi_n(0)|^2$ wave functions at the origin.

(ii) The correct level spacing and in particular $E(2p) < E(2s)$.

Both of these can remarkably well be obtained with linear potentials.

Once (i) and (ii) are achieved then the order of magnitude of radiative transitions are essentially fixed by e^2 and μ, the charge and mass of the fermions via the various dipole sum rules.

Also large $|\psi_n(0)|^2$ ensures large $\psi \to e^+e^-$ and $\psi' \to e^+e^-$ and may be of some help in increasing spin (η_c-ψ) splitting.

Clearly by postulating a potential that will approximate a linear potential effect, as was done to a certain extent in the work of Feinberg and Lee[7] (this was kindly pointed out to us by E. Eichten), these essential features can indeed be achieved. This model also incorporates the sharp rise in R around W = 4 GeV and some of the resonances above. It is furthermore clear that one will be even more successful trying two lepton types with different potentials, a case that is discussed in detail by Hagiwara and Sanda.[24]

The advantage of the charmonium scheme lies therefore in our view mainly in the natural way in which linear potentials arise (hopefully already at distances of .2 Fermis), due to the asymptotic freedom and infrared slavery of the theory as contrasted with the difficulty of obtaining potentials of this variety in the case of non-confining potentials.

VI. SUMMARY AND CONCLUSIONS

We have introduced above a scheme with unconfined ψ constituents. We find that the available data do not force on us the conventional confined charmed quark assumption though it is not easy to find as simple and elegant a scheme as the charm-SU(4) and color gauge theory.

The main motivation behind our suggestion is the desire to trace back to a single source all the spikes, resonances and $e\mu$ anomalies in the e^+e^- colliding beams.

The present scheme is much more flexible than the charm scheme. As we have seen various difficulties can be averted by including non-Abelian gluon exchange forces, low mass scalar exchanges, neutral strongly interacting L's, etc.

What are the decisive experiments which could disprove our hypothesis? It is clear that if new charged stable mesons with nonleptonic decays are found to be significantly pair produced at SPEAR or DORIS at the peaks observed around $W \simeq 4$ GeV then the charm hypothesis (or some similar model with several heavy quarks) is presumably true. The unique prediction of our scheme is that at $W \geq 2m_L \simeq 4$ GeV no resonant structure in any set of completely constrained events should be found as we sweep through the new resonances. The resonances decay predominantly into $L\bar{L}$ and each L emits an unobserved ν_L. More evidence for the charm scheme would be the production of a baryon with a sharply defined mass which decays semileptonically (the apparent violation of the $\Delta S = \Delta Q$ rule is in the GIM case a very efficient indicator of these events). The point is that while our scheme may in some version (similar to that of the non-Abelian model presented in Section III) contain low lying $L^+\bar{L}_o$, i.e. "mesonic" bound states, it is not expected to have baryon-L bound states of any kind.

On the other hand the medium strong interactions of the "heavy leptons" can be confirmed, e.g. by finding strong excess of $e\mu$ events in photoproduction (and also in meson-proton and proton-proton collisions) as compared to the estimated QED or ordinary hadronic background.

If repeated attempts to find charmed particles will continue to fail, alternatives will have to be looked for. If this indeed is the case and if the particular scheme we suggest here is confirmed then a new world of non-quark strongly interacting particles exists. It would mean that we are much further from a complete understanding of elementary particles but it is certainly a very exciting possibility.

ACKNOWLEDGEMENT

One of us (R.R.) wishes to thank Sidney Drell for warm hospitality extended to him at SLAC. We are grateful to several members of the SLAC theory group for discussions and, in particular, to J. Bjorken, F. Gilman and M. Perl. R.R. also thanks the Herman Rosenberg Foundation for financial support.

APPENDIX A

We would like to show that the experimental data on $y = \langle E_{ch}/W \rangle$ are inconsistent with the following set of assumptions often made in the "standard approach" to the new threshold in R :

(a) We have two components in the e^+e^- annihilation, the "old physics" component involving ordinary quarks only and which would scale at its $W \simeq 3.5$ GeV value to give $R \simeq 2.5$ and also a scaling $s\, d\sigma/dx$ inclusive distribution , and a "new" component .

(b) The new hadronic component with a threshold around $W \simeq 4$ GeV and which accounts for the rise in R is due to excitations involving new heavy isoscalar quarks c, c', etc.

$$e^+e^- \to c\bar{c} \to D\bar{D} \quad + \text{mesons} \tag{A.1}$$

Now the following "equipartition" theorem and identity of inclusive cross sections

$$\frac{d\sigma}{dP_{\pi^+}} = \frac{d\sigma}{dP_{\pi^o}} \; ; \; \frac{d\sigma}{dP_{K^+}} = \frac{d\sigma}{dP_{K^o}} \tag{A.2}$$

can be proved for the pions and kaons produced directly, i.e. not in the weak D decay. This results from I-spin invariance and is most easily illustrated by using Mueller's theorem to relate the inclusive cross section to the imaginary part of a fictitious scattering amplitude involving an outgoing π and a virtual "$\gamma_{c\bar{c}}$" target.

Let us write y as a weighted average of the various contributions

$$y = \frac{R_{old}\, y_{old} + R_{charm}\, (f_1\, y_1 + f_2\, y_2)}{R} \tag{A.3}$$

where $f_1 \approx (W - 2m_D)/W$ and $f_2 \approx 2m_D/W$ are roughly the fractions of final state energy taken by the directly produced mesons and by $D\bar{D}$, respectively.

(It is unlikely from a naive phase space consideration that the $D\bar{D}$ system will carry energy much in excess of the heavy $D\bar{D}$ rest mass.)

From scaling y_{old} should stay at its value ($\simeq .57$) at $W \approx 3.5$ GeV before the new threshold and $y_1 \approx .67$ due to equipartition.

For W = 8 GeV $f_1 \approx f_2 \approx 0.5$, $R_{old} \approx R_{new} \approx R/2$ and $y = 0.5$. Solving (A-3) for y_2 we find $y_2 = 0.2-0.3$. This small value should then show up as a break in the curve of y versus W at W = 4.2 GeV where the effect of D decays is maximal. These conclusions change somewhat to $y_2 = 0.3-0.4$ if a single heavy (ordinary) lepton is introduced.

In our scheme $y_2 = y_{\text{heavy lepton}} \approx 0.3-0.4$ arises very naturally, and again a break of the y curve at W = 4.2 GeV should occur.

Thus the systematic trend of decreasing y down to 0.5 at W = 8 GeV is very difficult to understand in either scheme.

One other mechanism which should be considered is that the jet structure recently observed at SLAC [1] is more pronounced for charged particles (which more closely follow the direction of the charged c parton as is the case at large Q^2 production of positive pions on protons). Also the jet angular distribution is somewhat forward-backward peaked. Thus the corrections due to loss of forward-backward moving fast charged particles is more severe than that of the neutrals.

A final comment (due to K. Lane) related to the question of observed energy is that at the ψ and ψ' $y = 2/3$ (apart from minor correction due to the decays via a single photon, which amount to $\approx 20\%$ and have anyway $R = 0.57$, and the η_c cascades, in the ψ' case also the $\psi' \to \gamma + \chi$ contribution), and hence is significantly different from y at the neighboring background points.

APPENDIX B

The $\gamma \rightarrow L\bar{L}$ production - if it is indeed pointlike as in the standard models - has implications for the inclusive spectrum s $d\sigma/dx$ at "large" ($x\approx 1$) and medium ($x\approx 0.5$) x values. The reason is that L decays into a few particles, e.g. $L \rightarrow \pi^+ \nu_L$, $\rho^+ \nu_L$, $A_1^+ \nu_L$ with fixed energy independent branching ratios. At sufficiently high W (and $m_{\nu_L} = 0$) π^+, ρ^+, A_1^+ have x distributions which are flat between zero and one (the upper limit extends to $x = 0.93$ at $W = 8$ GeV and to 0.99 at $W = 16$ GeV for $L \rightarrow \pi^+ \nu$). Since the experimentally observed s $d\sigma/dx$ is smaller by two orders of magnitude near $x = 1$ than its maximal value it is clear that the $x = 1$ region particularly at higher colliding beam energies could be a very direct test of the point-like heavy lepton production. Competing processes like $e^+e^- \rightarrow \pi^+\pi^-$ are expected to be extremely small because of form factors ($F_\pi(t)^2 \approx 10^{-4}$ at $t = 64$ GeV^2).

What happens when strong $L\bar{L}$ interactions are introduced ? A priori we might expect to have deviation from a point-like production because of the existence of a form factor, e.g. due to diagrams like in Fig. 4. However, models with only heavy B's ($m_B \geq 4$ GeV) responsible for the BL interaction we may have still effectively point-like cross sections for most of the region explored so far, i.e. up to $W = 8$ GeV, since roughly speaking the $L\bar{L}$ once formed has no strongly communicating channels into which it can break. Thus all our comments above will be relevant in this case and even more so since a larger fraction of R is attributed to L. There is, in this case, no obvious mechanism for systematically enhancing s $d\sigma/dx$ at small x as W increases further and further beyond $2m_L$.

The situation is quite different for the second class of models with $m_B \leq$ 1 GeV when $L\bar{L}$ production will in general be associated with one or more B's.

$$e^+e^- \rightarrow L\bar{L} + B + \ldots$$

where the B eventually decays into hadrons populating the small x region. Also the L's take less energy than $W/2$. Clearly the last process, if indeed important, will also effect for W significantly larger than $2m_L$ the $e\mu$ signal, its magnitude and also $e\mu$ energy distribution.

APPENDIX C

The partial wave Schrodinger equation for zero energy is an eigenvalue condition for the coupling g^2

$$-\frac{d^2}{dr^2} + \frac{\ell(\ell+1)}{r^2}\; U_\ell(r) \;=\; V(r)\,U_\ell(r) \quad . \tag{C.1}$$

It can be cast into an integral equation

$$U_\ell(r) \;=\; g^2 \int_0^\infty G_\ell(r,r')\,V(r')\,A_\ell(r')\,dr' \tag{C.2}$$

with

$$G_\ell(r,r') \;=\; \begin{cases} \dfrac{1}{2\ell+1}\, r^{\ell+1}/r'^{\ell} & r < r' \\[2ex] \dfrac{1}{2\ell+1}\, r'^{\ell+1}/r^{\ell} & r' < r \quad . \end{cases} \tag{C.3}$$

For V(r) a Yukawa potential we can symmetrize the kernel and change variables to

$$\ell^{-\mu_0 r/2} = X , \qquad \ell^{-\mu_0 r'/2} = Y .$$

This gives us

$$V_\ell(X) \;=\; g^2 \int_0^1 \left(\frac{2}{\mu_0}\right)^2 \frac{1}{2\ell+1}\, H_\ell(X,Y)\,V_\ell(Y)\,dY$$

$$H_\ell(X,Y) \;=\; \left(\frac{\ell n\, X_>}{\ell n\, Y_<}\right)^{\ell+\frac{1}{2}} \quad . \tag{C.4}$$

The smallest coupling g^2 for which (C.2) or (C.4) is satisfied is

$$g^2_{\min} \;=\; (2\ell+1)\,\mu_0^2/\lambda_{\max} \, .$$

λ_{max}, the largest eigenvalue of H, can be easily approximated by writing H as a finite matrix. We find for $\ell = 1$

$$g^{2\,\ell=1}_{min} = 2.2\,\mu_0^2$$

and for $\ell = 0$

$$g^{2\,\ell=0}_{min} = 0.36\,\mu_0^2$$

i.e. the coupling required for P-wave binding is roughly six times larger than the corresponding minimal coupling required for S-wave binding.

The estimate of B_{1s}, the S-wave binding, for $g^2 = g^{2\,\ell=1}_{min}$ is readily done by standard variational techniques. From Flugge's book[23] we find (page 190) $K \approx 12$, $p \approx 16$ and $B_{1s} = 0.8\ V_0$ where

$$V_0 = \frac{2.2\,\mu_0^2}{2m} = \mu_0 \text{ (in GeV)} \quad (m = \frac{m_L}{2} \simeq 1 \text{ GeV}) .$$

So if we insist on $B(1s) \simeq 0.4$ then $\mu_0 \leq 0.7$ GeV.

APPENDIX D

As an alternative non-Abelian scheme where instead of having just two L's we introduce two $SU(3)_L$ triplets

$$L_1 L_1^o, \quad \nu_{L_1} ; L_2 L_2^o, \nu_{L_2}$$

where ν_{L_i} are the associated "neutrinos" and weak interactions transform $L_i \rightarrow \nu_{L_i}$. We assume $m_{L_1} \simeq m_{L_1^o}$ and $SU(2)_L$ in the $L_1 L_1^o$ sector almost exact though $SU(3)_L$ is badly broken $(m(\nu_{L_1}) \approx m(\nu_{L_2}) \approx 0)$. The photon now carries L-spin and couples to L-spin = 0 and L-spin = 1 states. The ν_{L_i} and ordinary hadrons are L-spin singlets. The L-spin = 0 $L\bar{L}$ states are very broad since they always decay into $\nu_L \bar{\nu}_L$. Thus the ψ family and the P-wave states are to be associated with L-spin = 1.

For such states the exchange of L-spin singlet (triplet, i.e. non-Abelian) mesons gives rise to attraction (repulsion), respectively, as discussed in Section III and "nice" potentials can be likewise generated. Also the mass of the $B_{L\text{-spin}=0}$ particles can be very low since the ψ family would not be able to decay via emission of these ($\psi' \rightarrow \psi + B^V_{L=0}$ is forbidden for vector B^V) which could give long range potentials and further help with ψ spectroscopy.

REFERENCES AND FOOTNOTES

1. S. Ting, G. Feldman, R. Schwitters, G. Abrams and B. Wiik, Proceedings of the 1975 Lepton Photon Symposium, Stanford University (1975).
2. H. Harari, loc. cit.
3. M. Gaillard, B. Lee and J. Rosner, Rev. Mod. Phys. 47, 277 (1975); T. Appelquist and H. Politzer, Phys. Rev. Letters 34, 43 (1975).
4. For the most extensive attempt to explain all the ψ spectroscopy data with only c quarks, see E. Eichten, K. Gottfried, T. Kinoshita, K. Lane, and T. M. Yan, to be published.
5. S. L. Glashow, J. Iliopoulos and L. Maiani, Phys. Rev. D 2, 1285 (1970).
6. We will not consider the possibility that these are ordinary hadrons as various $\bar{B}B$ models for the ψ suggested.
7. A similar model is presently being investigated by G. Feinberg and T. D. Lee, Columbia Univ. preprint CO-2271-74. A more general discussion of ψ-models is given by R. P. Feynman, talk at the Conference on Quarks and the New Particles, Irvine (1975).
8. Martin Perl, Stanford Linear Accelerator Center preprint SLAC-PUB-1664 (1975).
9. In Finnish the L could stand for "lumoton" which means uncharmed.
10. Y. S. Tsai, Phys. Rev. D 4, 2821 (1971); J. D. Bjorken and C. Llewellyn Smith, Phys. Rev. D 4, 887 (1973).
11. F. Gilman, loc. cit.
12. It is still very puzzling why the two thresholds coincide within few hundred MeV's.
13. B. Barish et al., Phys. Rev. Letters 32, 1387 (1974).

14. In our scheme it is difficult to understand why μ-pairs associated with ν_μ and with $\bar{\nu}_\mu$ have different distributions of kinematic variables. For a phenomenological discussion, see V. Barger, R. Phillips and T. Weiler, Stanford Linear Accelerator Center preprint SLAC-PUB-1688 (1975).

15. There does not seem to be a clear cut generalization of this argument to the field theoretic Bethe Salpeter formulation. The positivity of $\sigma(\mu)$ simply reflects the fact that the Born kernel has, when viewed from the crossed (t) channel a form of the sum of perfect squares. This may not be true in higher orders since this property is destroyed by excluding diagrams which are direct channel two particle reducible.

16. One way to decrease this effect is to have L_1 and L_2 with opposite charges or opposite B-charges so that the contribution of $L_1\bar{L}_1$ and $L_2\bar{L}_2$ loops to the γ-B transition will cancel to a certain extent.

17. A dominantly scalar B interaction leads to LL binding as well. We do not worry about this at the present time.

18. This suppression operates equally strongly in $\psi \to \omega\pi\pi$ and $\psi \to \phi\pi\pi$ decays. Some further suppression due to the $\bar{\lambda}\lambda$ Okubo-Zweig-Lizuka rule is still expected for the $\psi \to \phi\pi\pi$ decay.

19. If we adopt Eq. (15) and choose a very large mass m_H^2 we can obtain a linear rise versus E_ψ^{lab} of $\sigma_{tot}(Lp)$ and $\sigma_{tot}(\psi p)$. The magnitude of $\sigma_{tot}(Lp)$ at $E_L = 100$ GeV is only a factor $\left[(g^2/m_H^2)/G_F\right]^2 \approx 10^5$ above the corresponding νp cross section, which is far too small.

20. $H_s = B_s$ could couple to $\theta_\mu^\mu \approx \sum_i m_i \bar{\psi}_i \psi_i$ so that $g_s^q/g_s^L \approx m_q/m_L \approx 10^{-1}$.

21. Y. S. Tsai, Rev. Mod. Phys. 46, 815 (1974).

22. We have to bear in mind the possibility that the eμ events may be due to D (charmed meson) decays which could provide alternative explanation for such

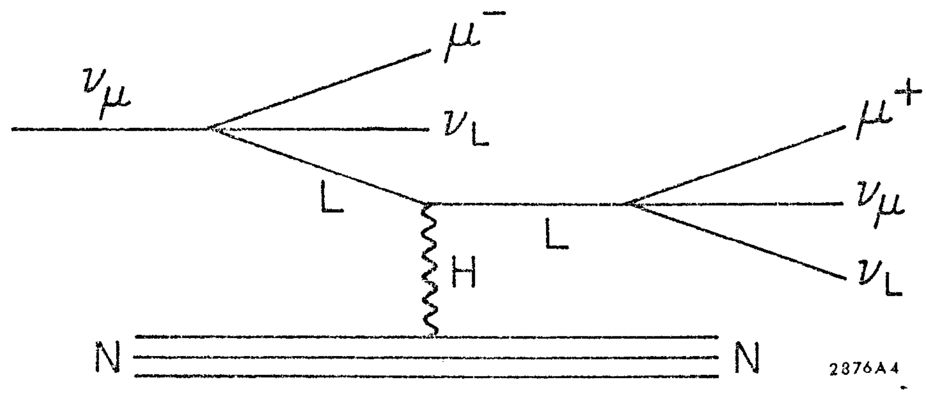
effects. All of the present work is clearly based on the premise that the identification of the eμ events as $\bar{L}L$ decays and the exclusion of the D decay hypothesis were indeed correct.

23. S. Flugge, Practical Quantum Mechanics (Springer-Verlag, Berlin, Germany, 1974).

24. B. Arbuzov, G. Segrè and J. Weyers, CERN preprint TH-2122-CERN; T. Hagiwara and A. Sanda, Rockefeller Univ. preprint COO-2232B-94.

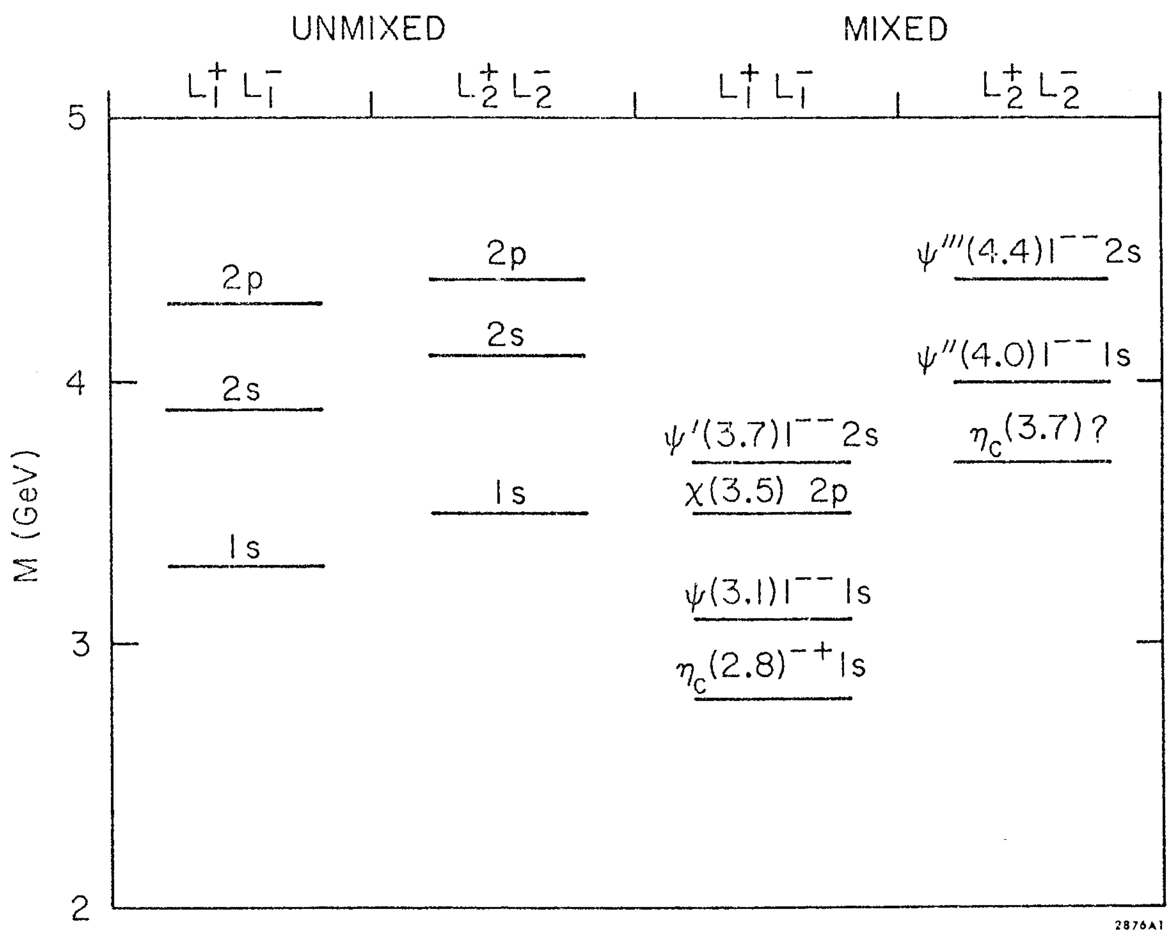
FIGURE CAPTIONS

1. A possible diagram for μ pair production.
2. A possible assignment of levels.
3. Some of the decays of ψ particles.
4. A diagram contributing to L form factor.



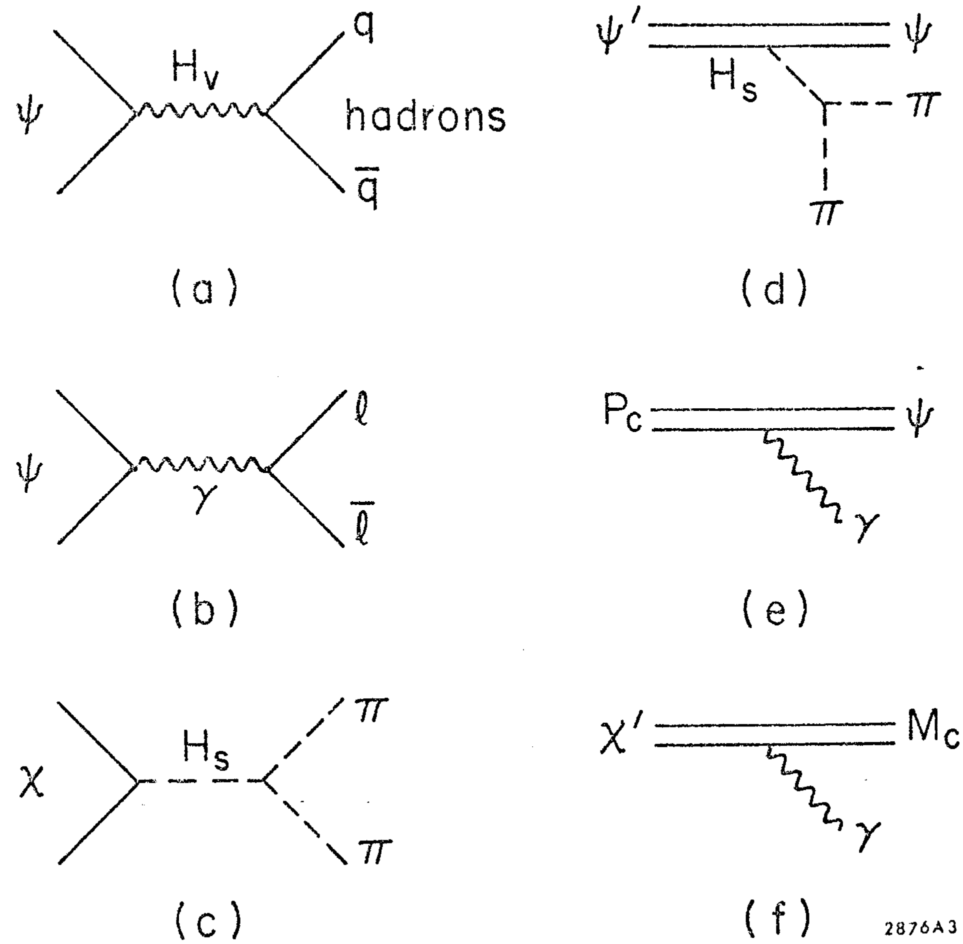
ν_μ
μ^-
ν_L
L
H
L
μ^+
ν_μ
ν_L
N
N
2876A4

Fig. 1



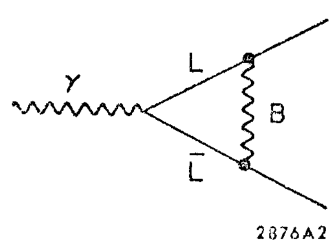
UNMIXED
MIXED
$L_1^+ L_1^-$
$L_2^+ L_2^-$
$L_1^+ L_1^-$
$L_2^+ L_2^-$
M (GeV)
5
4
3
2
2p
2s
1s
2p
2s
1s
$\psi'(3.7)1^{--}$ 2s
$\chi(3.5)$ 2p
$\psi(3.1)1^{--}$ 1s
$\eta_c(2.8)^{-+}$ 1s
$\psi'''(4.4)1^{--}$ 2s
$\psi''(4.0)1^{--}$ 1s
$\eta_c(3.7)$?
2876A1

Fig. 2



ψ
Hv
q
hadrons
q̄
(a)
ψ′
ψ
Hs
π
π
(d)
ψ
ℓ
γ
ℓ̄
(b)
Pc
ψ
γ
(e)
χ
Hs
π
π
(c)
χ′
Mc
γ
(f)
2876A3

Fig. 3



γ
L
B
L̄
2876A2

Fig. 4