Session B8: New Particles, Theoretical

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B 8 New Particles, Theoretical

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This report is the written summary of papers presented in Parallel Session B8 at the XIXth International Conference on High Energy Physics, August 24–30, 1978, Tokyo, Japan. The work of the minirapporteur was divided three ways: $Q\bar{Q}$ interaction and charmonium (Jackson), implications of the upsilon states (Rosner), and weak decays of heavy quarks (Quigg).

I. Introduction

"New particles" are defined for the purposes of Parallel Session B8 to be hadrons (almost totally mesons) connected with "charm" and other new "flavors," *i.e.*, $\psi, \psi', \ldots, D, D^*, \ldots F, F^*, \ldots, \Upsilon, \Upsilon', \ldots$. New leptons ($\tau$) and other new mesons (exotic, baryonium) are discussed elsewhere. The new particles are taken to be bound states of heavy quarks $Q(Q=c, b, t, \ldots)$ and light quarks q ($q=u, d, s$), states like ψ or Υ being $(Q\bar{Q})$ and states like D being $(Q\bar{q})$. In Part II the $Q\bar{Q}$ interaction and spectroscopy and dynamics of charmonium are discussed. The upsilon regime, including all that can be concluded from the just now emerging spectroscopy at 10 GeV, is treated in Part III. Weak decays of heavy quarks, not only the decays of D and F, but also the presumably sequential decays of the $(b\bar{q})$ mesons ($M\simeq 5$ GeV), are the subject of Part IV.

Only passing mention, if that, is made of many topics concerning the new particles and of many papers contributed to the Conference. For the latter neglect, the authors apologize. For the former, the reader is referred to other conference proceedings,[1,2] or to recent reviews, either largely theoretical[3,4] or more experimental.[5–7]

II. $Q\bar{Q}$ Interaction and Spectroscopy of Charmonium

A. Background

Since the discovery of the $\psi(3095)$ in November 1974, there has accumulated evidence for 8 relatively narrow states between 2.8 GeV and the charm threshold at 3.73 GeV, as well as the $\psi''(3772)$ and several less well resolved resonances up to 4.5 GeV. Many, but not all, of these states fit comfortably into an atomic energy level pattern analogous to positronium, as expected from a confined $c\bar{c}$ system based on QCD, with a potential,

$$v(r)=-\frac{4}{3}\frac{\alpha_s}{r}+V_c(r) \quad (1)$$

The first term is the QCD form of Coulomb's law, while the second is the confining potential,

Table I. Static and quasi-static interactions.

Type of interaction	Lorentz property: 4-vector $\gamma_\mu \otimes \gamma_\mu$	4-scalar $1 \otimes 1$	4-pseudoscalar $\gamma_5 \otimes \gamma_5$
Static potential	$V(r)$	$S(r)$	$P(r)$
Spin-orbit	$\frac{3}{2m^2}\frac{1}{r}\frac{dV}{dr}\boldsymbol{L}\cdot\boldsymbol{S}$	$-\frac{1}{2m^2}\frac{1}{r}\frac{dS}{dr}\boldsymbol{L}\cdot\boldsymbol{S}$	0
Tensor force	$\frac{S_{12}}{12m^2}\left[\frac{1}{r}\frac{dV}{dr}-\frac{d^2V}{dr^2}\right]$	0	$-\frac{S_{12}}{12m^2}\left[\frac{1}{r}\frac{dP}{dr}-\frac{d^2P}{dr^2}\right]$
Fermi hyperfine	$\frac{\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2}{6m^2}\nabla^2 V$	0	$\frac{\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2}{12m^2}\nabla^2 P$

traditionally assumed to vary linearly with distance at large separations. Early rough fits to the level scheme and values of the wave function at the origin were made with $\alpha_s \simeq 0.2$ (the value inferred from the naive estimate, $\Gamma_t(\psi)=\Gamma(\psi\to 3g)$), $m_c \simeq 1.4$ GeV/c^2, and $V_2(r)=ar$, $a\simeq 0.2$ GeV^{-2}.

The static limit yields no information on spin dependence. The relative amounts of central (1), spin–spin ($\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2$), tensor ($S_{12}$) and spin-orbit ($\boldsymbol{L}\cdot\boldsymbol{S}$) interactions are a priori unknown. Asymptotic freedom implies that at short enough distances and for heavy enough quarks the exchange of one massless vector gluon should dominate the interaction and lead to the familiar Breit-Fermi interaction (see Table I, $\gamma_\mu\otimes\gamma_\mu$). At larger separations, the situation may be different. The effective coupling of the "quanta" of the potential to the quarks may correspond to other Dirac operators. The proportions of 1, $\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2$, S_{12} and $\boldsymbol{L}\cdot\boldsymbol{S}$ can be different (see Table I for $1\otimes 1$ and $\gamma_5\otimes\gamma_5$).

Empirically, the spin splittings are large ($\sim$140 MeV for total ^3P_J interval, $\sim$250 MeV for $^3S_1-{}^1S_1$ intervals), and, in the P-states, very unlike the Landé interval rule. The Breit–Fermi interaction (effective γ_μ coupling to the quarks), based on either the whole $V(r)$ in eq. (1) or only the first term, gives too small a spin–spin splitting, too small a tensor force, and too large a spin-orbit splitting for the P-states. The early idea[8,9] of adding an anomalous Pauli ($\sigma_{\mu\nu}$) coupling did not solve all the problems. The spin-spin and tensor forces were enhanced, but so also was the already too large spin–orbit term. Another possibility is that the confining potential is an effective Lorentz scalar. This means a spin–orbit energy of opposite sign to the short-range QCD contribution. With a decrease in m_c and an increase in α_s, the QCD spin–spin and tensor forces can be made acceptably large (or nearly so), while the negative long-range spin–orbit energy combines with the large positive short-range contribution to give an acceptable overall effect,[10] at least for the ^3P states.

Incidentally, the view that the confining potential has an effective Lorentz scalar coupling to the quarks gains support from the Budapest version of the bag model.[11] The "scalar" spin-orbit energy of Table I is just the Thomas precessional energy,[12] produced at large separations in the bag model by rotation of the bag itself with quarks locked inside.

With potentials based on eq. (1), radiative transitions and wave functions at the origin (Γ_e) are in semi-quantitative agreement with data (modulo factors <2), except for the supposedly M1 transitions involving the $X(2.83)$ and $\chi(3455)$.[1,2,4] A contribution[13] points out again that naive estimates of rates for radially forbidden M1 transitions can be grossly wrong, but fails by a factor of 5 to remove the discrepancy for the $\psi\to\gamma X(2.83)$ transition.

B. Subsequent Developments in the Phenomenology

1. *Vector-scalar mixture*

Numerous authors[14–17] have fine-tuned the potential, using a short-range Coulombic potential, $-4\alpha_s/3r$, with $1/m^2$ spin-dependent corrections from $(\gamma_\mu+\kappa\sigma_{\mu\nu})\otimes(\gamma_\mu+\kappa\sigma_{\mu\nu})$, plus a long-range confining potential $V_c=ar$, with $1/m^2$ corrections from $(1-f)(1\otimes 1)+f(\gamma_\mu\otimes\gamma_\mu)$, where κ is the anomalous moment parameter and f is a small parameter measuring the amount of effective vector coupling in the confining potential. There are 5 parameters here (m_c, α_s, a, κ, f). Plausibly good fits are obtained with $\alpha_s\simeq 0.2$, $\kappa\simeq 4-5$, $f\lesssim 0.1$. If $\kappa\equiv 0$, then one needs a larger α_s, typically $\alpha_s\simeq$ 0.4–0.5 for charmonium. The various calculations are compared in detail in ref. 4.

2. Asymptotic freedom potentials and effective value of α_s

Fits to the spectroscopy indicate $\alpha_s\simeq 0.4$ in eq. (1), twice as large as the 3-gluon annihilation value. Should these be the same? Probably not. The smaller value, $\alpha_s\simeq 0.2$, is associated with short distances ($r\sim 1/m_c$) or equivalently $q^2\sim 2$ GeV^2. The effective value in the potential is associated with $q^2\simeq$ 0.3–0.4 GeV^2, the mean square momentum of the bound state. Asymptotic freedom implies that the coupling constant in the potential should be larger, although at such momenta it is difficult to be quantitative.

A rephrasing of these considerations leads to a variant of the QCD potential. The single-gluon exchange amplitude is proportional to $\alpha_s(q^2)/q^2$. If $\alpha_s(q^2)$ is approximated as a constant, the corresponding static potential is $1/r$. If, however, the famous logarithmic variation of α_s with q^2 is retained when taking the Fourier transform, the resulting potential is $1/r\sum_{n=0}^{\infty} a_n[\ln(r_n^2/r^2)]^{-(2n+1)}$. The inverse logarithms make the potential fall off with distance less rapidly than a pure Coulomb field. Calculations with such potentials, plus a Lorentz scalar confining potential, in coordinate space[18,19] or directly in momentum space with the Salpeter equation,[20] yield tolerable fits to both the charmonium spectroscopy and what little is known about the upsilon family, although the spin–spin and tensor forces are somewhat too small.

3. Inverted multiplets

The postulate of effective Lorentz scalar coupling to the quarks for the long-range confining potential is plausible and surely helps the interpretation of the spectroscopy. Schnitzer[21] has pointed out a probable consequence, namely "inverted multiplets" in $Q\bar{q}$ systems like the D family or F family. The point is that the QCD spin-orbit potential from gluon exchange has a coefficient that depends on the masses of the constituent quarks in such a way that it is reduced by a factor of 3 (relative to the term of opposite sign from the long-range Lorentz scalar potential) in the $Q\bar{q}$ system, compared with the $Q\bar{Q}$ system.[22] If, in charmonium ($Q\bar{Q}$), the proportions are such that the vector part dominates, but not overwhelmingly, the reduction by a factor of 3 *and* the larger size of the $Q\bar{q}$ system (because of smaller reduced mass) should cause the average spin-orbit interaction to have a negative coefficient of $\boldsymbol{L}\cdot\boldsymbol{S}$. The multiplets are inverted.

The argument is nonrelativistic. One concern is that relativistic effects may vitiate the argument, another is that tensor forces complicate it. Nevertheless, it seems probable that in the ^3P_J states of the D-family the $J=2^+$ state will lie lowest. Observation of such inverted multiplets would give strong support to the dominantly scalar Lorentz transformation property of the effective coupling of the confining potential. The experimental problems are not exactly trivial, but theorists can hope!

C. Instantons, Gluonic Excitations, Four-Quark States

1. Instanton contribution to the $Q\bar{Q}$ interaction

The existence of instantons in Euclidean QCD implies the existence of an instanton-generated interaction energy between Q and $\bar{Q}$, with origins in the difference in time evolution of the wave functions of each quark as they move through the "external" fluctuating instanton fields. The spin–spin part of this interaction has been considered as a potential solution to the uncomfortably large splitting between the $\psi(3095)$ and $X(2830)$.[23] Subsequent calculations[24,25] have examined the tensor force and spin–orbit interaction as well. It turns out that the effective coupling to the quarks is equivalent to the sum of scalar ($1\otimes 1$) and pseudoscalar ($\gamma_5\otimes\gamma_5$) exchanges with equal weight. From Table 1 it is evident that the spin-spin, tensor, and spin-orbit interactions thus have coefficients $+1/2$, -1, and $-1/3$, respectively, compared to vector coupling (for the same static potential).

Qualitatively, the instanton contribution to the potential is of the right sign to assist in fitting experiment for the spin–spin and spin–orbit parts (see Sec. A above), but not for the tensor force. Quantitatively, little can be said because the strength and radial dependence of the potential is a sensitive function of the distribution of sizes of the instantons, especially the maximum size.[130] Barring some peculiar radial dependence, it seems that, if the strength is such as to explain all of the 270 MeV splitting between ψ and X, the tensor and spin–orbit parts will be far too large for the phenomenology. Since instanton technology is in a state of rapid development at the moment, changes, clarification and/or improvements can be anticipated.

2. *Gluonic excitations*

A consequence of QCD, with its not so small coupling constant, is the possibility of excitation of gluonic degrees of freedom that give rise to additional states, beyond the Schrödinger spectrum. Various models are used to describe the combined system of quarks and gluons. One is the MIT/Budapest bag.[26,27] These models have extra states involving $Q\bar{Q}$ with gluonic excitations, and also purely gluonic excitations.[28] Another model[29] considers quantized vibrations of the color electric flux tube linking quarks. Still another[30] treats the transverse degrees of freedom of the gluon field on an equal footing with the quarks, all particles interacting *via* instantaneous static (gluonic) potentials. The models generally put the states with gluonic excitation 1 GeV or more above the lowest $Q\bar{Q}$ states. Thus in charmonium, these states lie above the charm threshold and will be difficult to identify. Perhaps for the upsilon sector, the lowest such state will lie below the flavor threshold! (Even so, it will not have $J^{PC}=1^{--}$, and thus will have to wait identification as an extra state among the analogs of the χ states.)

3. $Q\bar{Q}$ $q\bar{q}$ *states*

Problems with the simple $c\bar{c}$ picture (X(2830, χ(3454)—if pseudoscalars, too large splittings from ψ and ψ', absence of hadronic decays) lead to speculation[31,32] that the X(2830) and/or χ(3454) are not $^1S_0$ states of $c\bar{c}$, but are more complicated states, *e.g.*, $c\bar{c}q\bar{q}$, falling by chance among the charmonium levels.

For the X(2830) such an explanation is not implausible. The absence of a signal in the inclusive photon spectrum from the ψ(3095) and the observation of the state in $\psi\to\gamma\gamma\gamma$[33] permit the radiative width for $\psi\to\gamma X$ to be bounded, 10 eV$<\Gamma<$1300 eV, with the lower limit coming from the extreme assumption that $X\to\gamma\gamma$ dominates its decay. If X(2830) were the η_c, the radiative width should be of the order of 7–30 keV.[2,13] If, however, $X=c\bar{c}q\bar{q}$ (or even $c\bar{c}g$), a significantly smaller radiative width is expected because of the radically different dynamical nature of X and ψ. In this interpretation of X(2830), the pseudoscalar partner of the ψ is assumed to be fairly broad ($\sim$30 MeV) and close enough to the ψ to have so far escaped discovery.

The interpretation of χ(3454) as a beast different from $c\bar{c}$ appears necessary, but beset with some difficulty. Here the corresponding data[33] imply the bounds, 1.8 ± 0.9 keV$<\Gamma(\psi'\to\gamma\chi(3454))<6$ keV. Naively, the radiative width is estimated to be $\lesssim$17 keV. With a significantly different quark content for χ(3454) the photonic width is expected to be much smaller even than the lower bound from experiment. Furthermore, the assumption of χ (3454)$\to\gamma\psi$ dominating the χ decay, attendant to obtaining the lower bound, seems a priori unlikely. A possible state at 3.6 GeV (see D(4) below) adds to the confusion.

D. Problems and Limitations

Numerous contributions to the conference concerned the improvement of the description of the charmonium spectroscopy in terms of the nonrelativistic Schrödinger equation for the $c\bar{c}$ system. Some of these have been discussed in Section B. It is now time to issue a warning on the limitations of such approaches.

1. *Relativistic corrections*

Charmonium is far from a truly nonrelativistic system. In the ground state the mean kinetic energy is $\langle T\rangle\simeq$200–250 MeV. This means that each quark has $\beta^2\simeq\langle T\rangle/m_c\simeq0.16$ with respect to the center of mass. Relativistic corrections of the order of 10 to 30% can be anticipated for any dynamical quantity. The effect of relativistic corrections on spin–flip radiative transitions[13] has already been mentioned. The Salpeter equation provides a consistent description of relativistic effects, within the framework of a known kernel. Some work in this direction already exists.[10,20]

2. $\Psi(0)$ and Γ_e

Experimental values of Γ_e are invariably used to constrain phenomenological fits, with the simple proportionality of Γ_e to $|\Psi(0)|^2$ taken for granted. There are, however, corrections, some kinematic and some dynamic. The kinematic ones are relativistic. Just as the coupling of a bound state of angular momentum l to e^+e^- (or other exothemic channels) is proportional in lowest order to lth derivative of the radial wave function, divided by M^l,[35,3] for $l=0$ there are corrections involving $\Psi^{(2n)}(0)/M^{2n}$, $n=1, 2, \ldots$, in addition to the lowest order $\Psi(0)$.[35] These can be exhibited in closed form.[36] Dynamical corrections of order α_s arise from gluonic radiative corrections to the $Q\bar{Q}\gamma$ vertex.[37,38] Assuming that the QCD radiative corrections parallel exactly the QED ones, the formula for Γ_e has a factor $(1-16\alpha_s(m_c^2)/3\pi+\ldots)$ multiplying $|\Psi(0)|^2$. This can introduce factors of 2 in the values of $|\Psi(0)|^2$ used as constraints on potentials. Models with $\alpha_s\simeq0.4-0.5$, formerly rejected because of too large predictions for Γ_e, now become permitted, but uncertainty in the exact expressions for the corrections make too much fine-tuning of the potential unproductive.

3. Mixing and coupled channels

In the $c\bar{c}$ Schrödinger description mixing of triplet states with the same J, but different L values, occurs *via* the tensor force. In particular, $^3S_1$ mixes with $^3D_1$. The $\psi(3772)$ is interpreted[39] as mainly $^3D_1$, but with an admixture of $^3S_1$ from the nearly $\psi(3684)$. Quantitative estimates of the mixing are hindered by two effects, the presence of the dynamical and kinematic corrections just described and the presence of strongly coupled channels($\psi\leftrightarrow D\bar{D}$). Indeed, for $\psi(3772)$ and higher states inclusion of all open and nearby closed channels is necessary for a consistent description of the dynamics. Treatment of the $c\bar{c}$ sector in this energy range without other communicating channels is likely to be totally misleading, at least in detail.

Only the Cornell group has made a serious attempt at a coupled-channel calculation,[39,1,40,129] and even that is with a stylized model. The complexity apparent in the total cross section in e^+e^- annihilation from 3.7 to 4.5 GeV argues for more and better, fully relativistic calculations.

4. Peculiar states

The observation by the DESY–Heidelberg group[41] of a state at 3.6 GeV in the cascade decay $\psi'\to\gamma_1\chi$, $\chi\to\gamma_2\psi$, with a product of branching ratios $B_1B_2=(2.8\pm1.2)\times10^{-3}$, brings to three the number of peculiar states in the charmonium spectrum. The unusual properties of $X(2830)$ and $\chi(3454)$ have been widely discussed,[1–7] the main one being the absence of detected hadronic decay modes. The new state, if so it be, is no less peculiar. Its closeness below $\psi'(3684)$ indicates a $^1S_0(\eta_c')$ assignment, but the relative largeness of B_1B_2 is a problem. The first radiative transition (allowed M1) will have $\Gamma_1\lesssim1$ keV, or $B_1\lesssim4\times10^{-3}$. This means that $B_2>0.7$! Amusingly enough, the dominance of the radiative decay $\chi\to\gamma_2\psi$ is consistent with the absence of any peak at 3.6 GeV in the hadronic decays of the χ states.[42] But an η_c' with negligible hadronic decays is as difficult to stomach at 3.6 GeV as it was at 3.454 GeV.[43]

Other assignments ($^1D_2$ of $c\bar{c}$, $c\bar{c}g$) are equally implausible, given the value of B_1B_2. The radiative width for $\psi'\to\gamma_1{}^1D_2$ can be appreciable ($\lesssim0.5$ keV) because of the $^3D_1$ admixture in ψ'. But the second radiative transition should be drastically inhibited (spinflip, $\Delta L=2$, with negligible $^3D_1$ admixture in ψ). Since the hadronic width of a D state should be of the order of 25 keV or greater,[44] it is very difficult to understand $B_2\simeq1$. The assignment $c\bar{c}g$ (gluonic excitation) has a different enough quark/gluon configuration from a $c\bar{c}$ ψ' that it seems difficult to get a large enough matrix element for the first (or the second) radiative transition.

The assignment of $c\bar{c}q\bar{q}$ is somewhat more viable. De Rújula and Jaffe[31] predict that their lowest "molecular" charmonium states ($I=0$, 1; $J^{PC}=0^{++}$) lie at about 3.6 GeV. The radiative transitions would both be El. The very different wave function [*e.g.*, 0.915

$(DD+X\eta)+0.404(D^*D^*+\psi\omega)$ for $I=0]$ will presumably make B_1 much smaller than the scaled p-state value of $\sim 3\times 10^{-3}$. The apparently dominant decay $\chi(3.6)\to\gamma\psi$ might possibly come (*via* VMD) from the presence of $X\eta$ or $X\pi$ in the wave function, but the $X\eta$ or $X\pi$ content should make hadronic decays most important.

It is too early to be certain of anything about the 3.60 GeV region, but the reported state certainly warns us to be alert for phenomena outside the nonrelativistic $c\bar{c}$ model.

E. Theorems and Near Theorems for the Schrödinger Equation

Although Schrödinger equation considerations have serious limitations in the details of $Q\bar{Q}$ spectroscopy, the gross behavior is described well enough. As $m_Q\to\infty$, the dynamics should become more nonrelativistic and the Schrödinger equation a better guide to future spectroscopy, apart from "extra" states. We therefore summarize relevant theorems and near-theorems.

1. Level ordering

Let the energy levels be labelled $E(n_r+1, l)$ where n_r is the number of radial nodes and l is the orbital angular momentum. Provided the potential $V(r)$ satisfies some weak constraints[45] (conditions A and B of Grosse and Martin), the ordering of levels is $E(1\text{s})<E(1\text{p})<E(2\text{s})<E(1\text{d})<E(2\text{p})$.

2. l dependence

For any potential, $E(1, l)$ is a concave function of $l(l+1)$, as is also $\sum_{k=1}^{n} E(k, l)$. For example, $E(1\text{p})>(1/3)E(1\text{d})+(2/3)E(1\text{s})$. For charmonium, with $E(1\text{s})=0$, $E(2\text{s})\simeq 590$ MeV, $E(1\text{p})\simeq 430$ MeV (*e.g.*, ^3P_J), $E(1\text{d})\simeq 680$ MeV. Hence LHS=430 MeV and RHS=226 MeV, a rather loose inequality.

It is conjectured by Grosse and Martin[46] that, if V satisfies their A and B, then $E(1, l)$ is a concave function of l. Then $E(1\text{p})>1/2[E(1\text{d})+E(1\text{s})]>1/2[E(2\text{s})+E(1\text{s})]$. Now the the numbers are $430>340>295$ MeV, inequalitites that are becoming interestingly close. This conjecture has the status of a "near theorem," being valid in a number of special circumstances ($V=r^2+\lambda\delta V$, $\delta V\in A, B$; $V=r^\gamma$, $0>\gamma>2$; large l).

3. $|\Psi_n(0)|^2$ for different n, same $Q\bar{Q}$ system

If the potential is such that $V''(r)\gtrless 0$ for all r, then[47] $|\Psi_2(0)|^2\gtrless|\Psi_1(0)|^2$. It is conjectured that this inequality has the generalization, $|\Psi_{n+1}(0)|^2\gtrless|\Psi_n(0)|^2$, valid for all n. For large n it has been established for power law potentials[48] and more generally,[49] within the framework of the WKB approximation.

4. Level spacings for the same $V(r)$, but different reduced masses

Let the energy levels be written in an obvious notation as $E(M, n, l)$ and $E(m, n, l)$. Then, for any potential and any n value,[50] $E(M, n, 0)<E(m, n, 0)$ if $M>m$. Study of the equation for continuous l values leads to the results,[46] $E(M, 1, 1/2(\sqrt{M/m}-1))<E(m, 1, 0)$ for any potential provided $M>m$, and $E(M, 1, \sqrt{M/m}-1)<E(m, 1, 0)$ if the potential satisfies A and B. Application of the last inequality and the concavity conjecture in § 2 above, using the ψ, Υ, and Υ' masses, leads to the result, $M_b-m_c>3.29$ GeV,[46] 0.11 GeV better than the naive comparison of ψ and Υ masses, assuming equal binding.

5. Mass dependence of $|\Psi_n(0)|^2$ for same $V(r)$

For $n=1$ and $V''<0$ with $V'\geqq 0$, it can be proved that $d[|\Psi_n(0)|^2/m]/dm>0$, as the quark mass m is varied.[50] The inequality also holds for arbitrary n, provided the potential has power law form (and $V''<0$),[48] and also for any potential in the limit of large n where WKB arguments can be applied.[51] This inequality can be employed in a comparison of Γ_e of ψ with Υ and ψ' with Υ' to make statements about the magnitude of the charge of the upsilon's quark Q. [See Part III.]

6. Inverse scattering problem

For the conventional central field problem with an energy continuum in addition to bound states, knowledge of the positions of the bound states, and the phase shift as a function of energy for all energies is necessary and sufficient for construction of the unique potential for that partial wave. For a confining potential, with no continuum and a spectrum of discrete states extending to infinity, it has recently been shown[52] that the necessary and sufficient data for $l=0$ states are the positions

Table II. Comparison of ψ and Υ families.

	ψ Family			Υ Family		
3S	ψ'' (4.04)			Υ'' (10.38)		
2S	ψ' (3.686)		0.945 $\pm$0.010	Υ' (10.016)		0.92 $\pm$0.04
		0.591 $\pm$0.003			0.556 $\pm$0.003	
1S	J/ψ (3.095)			Υ (9.46)		

of the levels and the values of $\Psi_n(0)$. When one recalls the relation between the phase shift $\delta(E)$ and the wave function $\Psi_E(0)$ *via* the Jost function, one sees that the conventional theorem and the new one are basically the same.

Educational calculations have been made by Thacker, Quigg and Rosner,[53] showing for reflectionless potentials (simulating confinement) how much (or little) knowledge can be gained about the potential from the location of a few bound states.

III. Implications of the Υ States

A. Introduction

The Υ family discovered at Fermilab[54–56] and confirmed at ISR[57] and DORIS[58] implies the existence of a fifth quark "*b*."[59] It has charge $e_Q=-1/3$. It is very likely a color triplet, just like the first four quarks (u, d, s, c). It may have a heavier $e_Q=2/3$ partner "t", or a charge $-1/3$ relative "h."

Alternative interpretations of the Υ family are discussed (and found unlikely) in § B. As a corollary, properties of systems containing quarks of other masses, charges, or color representations than b are noted. Some implications of the new quark for heavy particle spectroscopy are mentioned in § C. The possibilities for still heavier quarks, and for searches for other new particles, are greatly enhanced by the discovery of the b (§ D).

B. The Υ as a $b\bar{b}$ State

The signal for the Υ as a $\mu^+\mu^-$ resonance in hadronic interactions was very similar to that of J/ψ at a lower mass: a sharp peak above a rapidly falling continuum. The peak is narrow[58] and has at least two higher-mass partners.[54,56] All of these properties are similar to the charmonium system (J/ψ, ψ', . . .) and, indeed, the mass splittings in the two families are remarkably similar. A comparison is shown in Table II (ψ masses: ref. 7; Υ splittings: refs. 56, 58).

The ψ family is a bound system of a charmed quark c and antiquark.[1–4] This idea was generalized to heavier quarks well before the discovery of the Υ. Thus, the existence of three narrow levels[60] and the value[60,61] $\Gamma(\Upsilon\to e^+e^-)\sim 1$ keV (for $\Upsilon=b\bar{b}$, $e_b=-1/3$) were anticipated. The remarkable coincidence of mass splittings (Table II) was somewhat more of a surprise.[1,60] The large Υ'–Υ splitting is not a problem for a customary $b\bar{b}$ interpretation[62] of the Υ levels. It has, however, led to some interesting alternative proposals.[63–69]

The only nonrelativistic potential for which the level structure is independent of quark mass is $V(r)=C\ln(r/r_0)$.[69] (This potential was first suggested for charmonium[70] because it gives an orderly decrease of leptonic widths of $n\ {}^3S_1$ states in accord with experiment.) "Duality" schemes also give equal $2S$–$1S$ splittings for all vector meson states.[71] Now, equal $2S$–$1S$ splittings for two different families arise from a wide variety of potentials. In the Coulomb+linear example, which has some theoretical underpinnings,[1–4] $M_{\Upsilon'}-M_{\Upsilon}\simeq M_{\psi'}-M_{\psi}$ when one doubles the strength of the Coulomb force[16,18,69,72,129] with respect to the value used in ref. 60. The nonrelativistic prediction[73] for leptonic widths then increases, since the larger Coulomb interaction pulls the wave function toward the origin. Since relativistic corrections tend to reduce leptonic widths,[37,38,129] this is probably acceptable, as already mentioned in Part II, § D(2).

Figure 1 compares level splittings in two

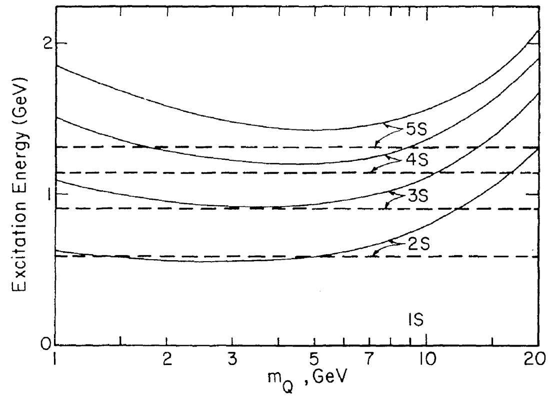

Fig. 1. Comparison of level splittings in two quarkonium potentials as functions of quark mass m_Q. Solid lines: $V(r)=-0.56/r+0.163r$. Dashed lines: $V(r)=0.733\ \ln r$. (Units are in GeV or GeV^{-1}.) [Note added: the experimental $\psi'-\psi$ and $\Upsilon'-\Upsilon$ splittings can be reproduced with $V(r)=-0.507/r+0.17\ r$, *i.e.*, with $\alpha_s=0.38$.]

extreme examples with equal $\Upsilon'-\Upsilon$ and $\psi'-\psi$ splittings: the logarithmic potential, and a Coulomb+linear potential. The effects of the Coulomb potential are clearly enhanced as the quark mass increases and hence as the shorter-distance part of the potential is probed. Note the similarity of the $3S$ levels in the two potentials for the Υ family. A sixth quark (§ D), especially if it gives rise to a new vector meson "ζ"[69] heavier than Υ, will distinguish between the two.[74] If there really is a short-distance Coulomb interaction between quarks, the Υ family is telling us that it will be easier to see this interaction (that is, lower quark masses will suffice) than originally anticipated.[75]

Let us discuss some of the evidence that Υ really *is* a $b\bar{b}$ family. Several points are summarized in Table 3.

The narrowness of the Υ may be ascribed to the Okubo–Zweig–Iizuka, *et al.* (OZI)[76] rule. The Υ then must be below some threshold, indicating the need for new heavy objects. The Υ and presumably the Υ', Υ'', ... are then viewed as bound states of these objects.

We have assumed Υ' and Υ are related. They may not be.[67] The decay $\Upsilon'\to\Upsilon+\text{hadrons}$, estimated using scaling arguments[1,77] (§ C) to be $\approx 40\%$ of all Υ' decays, would provide evidence that Υ and Υ' are members of the same family. Radiative decays $\Upsilon'\to\Upsilon\gamma\gamma$ would be still more conclusive, though rarer.[78]

If the Υ is made of spinless bosons,[63] it is not the ground state. It decays rapidly to the ground state and a photon of energy several hundred MeV (also to hadrons), leaving only a fractional-percent branching ratio to lepton pairs. Preliminary indications[56,79] are that $B(\Upsilon\to\mu^+\mu^-)$ exceeds a percent, as expected on the "standard" model.[62] We shall thus assume the constituents of Υ are fermions. In this manner the Υ can be a $^3S_1$ state. Its decay to any lower $^3S_0$ state is presumably at least as rare as that of the ψ, probably occuring with a rate well below a percent.

Could the fermions in the Υ have spin 3/2? Then $R\equiv\sigma(e^+e^-\to\text{hadrons})/\sigma(e^+e^-\to\mu^+\mu^-)$ above flavor threshold ($\approx 10\frac{1}{2}$ GeV; see § C) should grow rapidly. We shall assume spin 1/2 quarks.

The quarks in the Υ probably have charge $-1/3$. This was expected on the basis of production estimates,[62] and is much more likely as a result of the measurement (average of values in ref. 58),

$$\Gamma(\Upsilon\to e^+e^-)=1.26\pm0.21\ \text{keV}. \quad (2)$$

This value is much more compatible with $e_Q=-1/3$ than with $e_Q=2/3$ in various specific

Table III. Evidence that $\Upsilon=bb$. (b=charge $-1/3$ color triplet quark)

Hypothesis	Υ is a bound state of new objects	New objects are fermions (quarks)	Quarks have $e_Q=-1/3$	Quarks are color triplets
Motivation	Narrow width; (ref. 76) Similarity of ψ, Υ families (Table II)	Large (few %) leptonic branching ratio (refs. 56, 79)	$\Gamma(\Upsilon\to e^+e^-)$ (refs. 18, 19, 38, 50, 62, 69) Universal $\Gamma/e_Q{}^2$ (refs. 71, 80, 81, Fig. 2)	Large (few %) leptonic branching ratio (refs. 56, 78)
Further tests	$\Upsilon'\to\Upsilon+\dots$ implies Υ and Υ' are related (Rate est., refs. 76, 77)	$\Gamma(\Upsilon\to 0^+\gamma)$ big; $B(\Upsilon\to e^+e^-)$ small if Υ made of bosons (ref. 63)	If $\Gamma(\Upsilon'\to e^+e^-)$ <0.6 keV, e_Q must be $-1/3$ (ref. 50) [fulfilled (ref. 58)]	$B(\Upsilon\to e^+e^-)$ $<\frac{1}{2}\%$ for color sextet, $e_Q=-1/3$ (ref. 64)

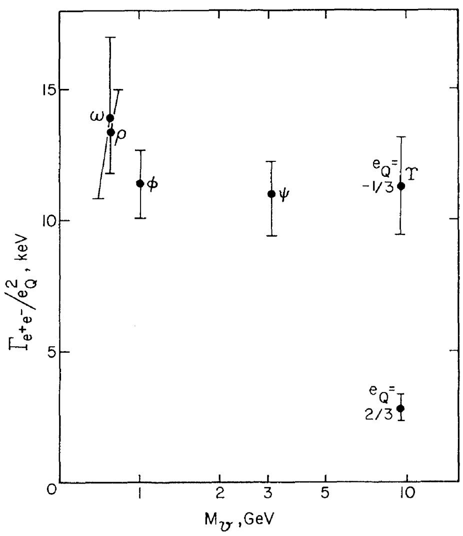

Fig. 2. Leptonic widths $\Gamma_{e^+e^-}$ (refs. 7, 81) normalized by squares of quark charges $e_Q{}^2$, as functions of vector meson mass $M_{\mathcal{V}}$.

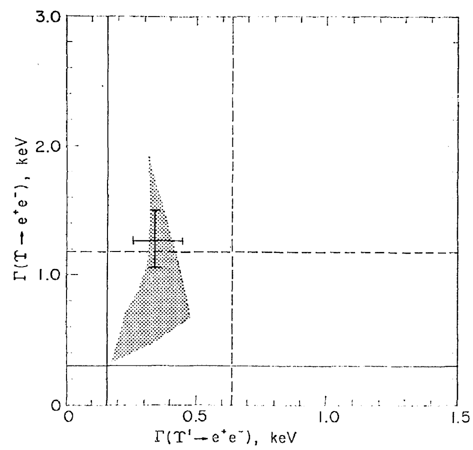

Fig. 3. Lower bounds for leptonic widths of Υ and Υ' (ref. 50), together with data presented at this Conference (ref. 58). The shaded area represents the range of predictions of twenty potentials reproducing the ψ and ψ' masses and leptonic widths, for $e_Q=-1/3$. Solid and dashed lines correspond to lower bounds for $e_Q=-1/3$ and 2/3, respectively. Equation (2) and $\Gamma(\Upsilon''\to e^+e^-)=0.36\pm0.09$ keV are used.

potential models.[18,19,38,50,62,69,72] The leptonic widths of ρ, ω, ϕ, and ψ obey a nearly universal law[18,71,80]

$$\Gamma(\mathcal{V}\to e^+e^-)/e_Q^2=11.9\pm0.8 \text{ keV}, \qquad (3)$$

as shown in Fig. 2.[5,81] The Υ is consistent with this behavior for $e_Q=-1/3$, but not for $e_Q=2/3$.

Since leptonic widths are proportional to the square of the wave function at the origin,[73]

$$\Gamma(\mathcal{V}\to e^+e^-)=\frac{16\pi\alpha^2}{3}Ne_Q^2|\Psi(0)|^2/M_{\mathcal{V}}^2 \qquad (4)$$

(N=dimension of quarks' color representation) and since

$$|\Psi(0)|^2=\frac{m_Q}{4\pi}\left\langle\frac{dV}{dr}\right\rangle, \qquad (5)$$

one can relate leptonic widths in the Υ family to those in the ψ family if one knows how $\langle dV/dr\rangle$ changes with m_Q. This has been done for a restricted class of potentials[50] (see Part II, § E (5)). There results a set of lower bounds

$$\Gamma(\Upsilon\to e^+e^-)\geq(0.3,\ 1.2) \text{ keV}$$
$$\Gamma(\Upsilon'\to e^+e^-)\geq(0.17,\ 0.63) \text{ keV} \qquad (6)$$

for $e_Q=(-1/3,\ 2/3)$, respectively. These are conservative, based on $m_Q/m_c\geq2.6$. Most potential models have m_Q/m_c lying between 3 and 4, and Grosse and Martin[46] have established $m_Q-m_c\geq3.29$ GeV for $m_Q/m_c\geq3$.

While the experimental result (1) does not permit a distinction between $e_Q=-1/3$ and 2/3, the measurement of $\Gamma(\Upsilon'\to e^+e^-)$[58] is very helpful. This is because $\Upsilon\to e^+e^-$ probes a *terra incognita* (the deepest part, in fact, yet seen) of the $Q\bar{Q}$ potential, while the physics of the higher-lying Υ' level is restricted to a much greater degree by information from charmonium, and thus is a better indicator of e_Q.[53] Since the measured value for $\Gamma(\Upsilon'\to e^+e^-)$ lies below 0.63 keV, e_Q must be $-1/3$ (see Fig. 3).

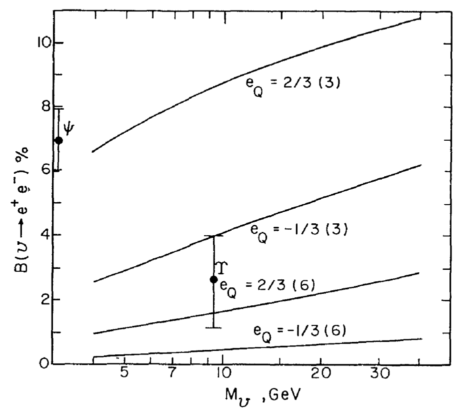

Fig. 4. Predicted leptonic branching ratios for quarks of various charges (−1/3, 2/3) and colors (3, 6): $B=[\Gamma_h/\Gamma_l+7]^{-1}$, with α_s extrapolated from ψ using asymptotic freedom (ref. 83).

Color sextet quarks[82,64–66] raise predicted leptonic widths by 2 (eq. (3)), but hadronic widths[64] by 49/2. This is because sextet quarks couple copiously to gluons. The predicted branching ratio for $\Upsilon\to e^+e^-$ is far lower for color sextet quarks, as may be seen in Fig. 4. Here we have used[3,4]

$$\Gamma(\Upsilon\to 3\text{ gluons})=\frac{160}{81}(\pi^2-9)\alpha_s^3\frac{|\Psi(0)|^2}{M_\Upsilon^2}\times\begin{Bmatrix}1 \text{ for color 3 quarks}\\ 42/2 \text{ for color 6 quarks}\end{Bmatrix} \quad (8)$$

and extrapolated α_s from the ψ using asymptotic freedom.[83] We have also taken account of $\Upsilon\to\mu^+\mu^-$, $\Upsilon\to\tau^+\tau^-$, and $\Upsilon\to\gamma\to$hadrons, assuming $R=4$[58] for the last process. With $e_Q=-1/3$, $B(\Upsilon\to e^+e^-)$ is about 0.4% for color sextets, and nearly ten times that value for color triplets. Color sextet quarks are more strongly bound in QCD; this is one reason they were suggested for the Υ states. The stronger binding spreads apart the $1S$ and $2S$ levels.[64–66] It also packs more narrow $^3S_1$ levels below flavor threshold; for color triplet quarks one estimates[84] three or possibly four levels (§ C) while one specific sextet model[64] predicts five.

If the jump in R above flavor threshold can be measured precisely enough, and if no other quark or lepton thresholds lie in the same region, both e_Q and N follow:

$$\Delta R=e_Q^2N=\begin{Bmatrix}1/3\\2/3\\4/3\\8/3\end{Bmatrix}\text{ for }(e_Q,N)=\begin{Bmatrix}(-1/3,3)\\(-1/3,6)\\(2/3,3)\\(2/3,6)\end{Bmatrix}.$$

If the Υ were composed of color sextet quarks, these could not be stable when incorporated singly into hadrons. Two experiments[85] indicate the cross section for production of particles of mass $M\simeq M_\Upsilon/2\simeq 5$ GeV/c^2 with lifetimes more than 5×10^{-8} sec is less than 1/10 that of the Υ at 400 GeV/c. To enable sextet quarks to decay, one would have to introduce a new vector boson carrying both color and flavor.

The ratio of Υ' to Υ leptonic widths has been quoted as[58]

$$\frac{\Gamma(\Upsilon'\to e^+e^-)}{\Gamma(\Upsilon\to e^+e^-)}=\begin{Bmatrix}3.4\pm0.9 \text{ (DESY-Heidelberg)}\\ \approx 3 \text{ (DASP II)}\end{Bmatrix}.$$

This ratio can be used to extract $|\Psi_{2S}(0)|^2/|\Psi_{1S}(0)|^2$ with the help of (3). Figure 5 shows the corresponding ratios for ρ and ρ',[86] J/ψ and ψ',[7] and Υ and Υ',[58] along with the predictions for various potentials. A trend toward Coulomb-like behavior is clearly visible as m_Q increases (hence as the quark Compton wavelength decreases, probing shorter distances).

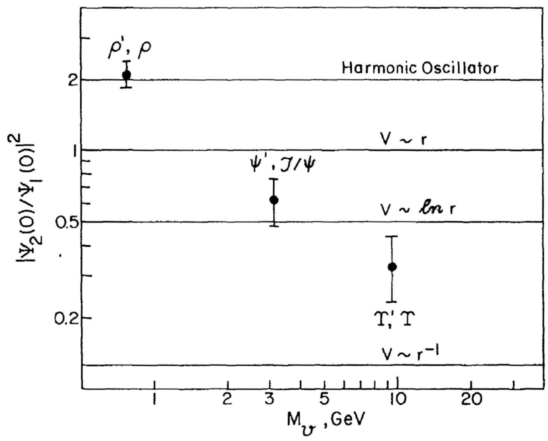

Fig. 5. Ratios of $2S$ and $1S$ squares of wave functions at the origin for various potentials: 2 for oscillator, 1 for linear, ~0.5 for logarithmic, and 1/8 for Coulomb.

C. Heavy Particle Spectroscopy

How good is the nonrelativistic approximation for $Q\bar{Q}$ systems? To illustrate, consider the logarithmic potential that gives the $\psi'-\psi$ and $\Upsilon'-\Upsilon$ splitting.[69] The constant internal kinetic energy is $T=0.37$ GeV. The β^2 for each quark is $\beta^2\simeq0.21$ for $m_Q\simeq1.5$ GeV and $\beta^2\simeq0.07$ for $m_Q\simeq5.0$ GeV, going as $\beta^2\simeq T/m_Q$ for large m_Q. Relativistic corrections are still appreciable at the ψ, but die away rapidly above the Υ. Heavy quarks thus could be a boon to nonrelativistic spectroscopy. In particular, the Υ system should allow reliable reconstruction of a $Q\bar{Q}$ potential via inverse-method.[53]

The ψ' was difficult to observe in hadronic interactions, but the Υ' appeared almost directly with the Υ. Production ratios at 400 GeV/c are[56]

$$B_{\mu\mu}\left.\frac{d\sigma}{dy}\right|_{y=0}(\Upsilon:\Upsilon':\Upsilon'')=1:(0.30\pm0.03):(0.155\pm0.016). \quad (9)$$

(These agree with estimates of ref. 71 and Ellis, *et al.*, ref. 62.) Figures 4 and 6, the latter incorporating some scaling arguments,[48,69,77,78] show why the Υ' was relatively more prominent than the ψ'. The Υ leptonic branching ratio is expected to be about half that of the ψ; the leptonic branching ratio of the Υ' could approach nearly double that of the ψ'. Moreover, the production ratios of the two states could be more similar for Υ' and Υ than for ψ' and ψ.

The successful description[62,71] of the ratio

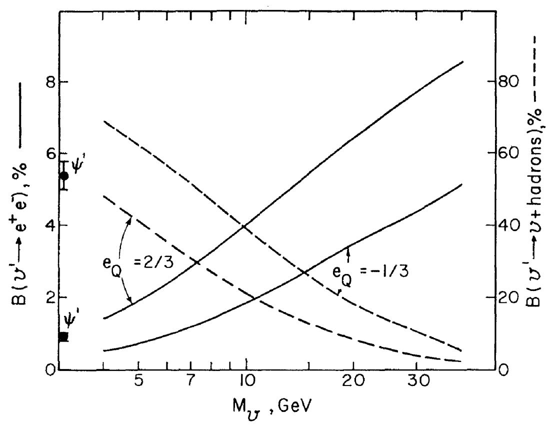

Fig. 6. Branching ratios for $\Upsilon' \to e^+e^-$ (solid curves, left-hand scale, lower ψ' point) and $\Upsilon' \to \Upsilon +$ hadrons (dashed curves, right-hand scale, upper ψ' point) as functions of vector meson mass. Color triplet quarks and $\Gamma(\Upsilon' \to e^+e^-) = 5e_Q{}^2$ keV are assumed. [*cf.* universality in Fig. 2; the coefficient is chosen to reproduce $\Gamma(\psi' \to e^+e^-)$.] All Υ' widths proportional to $|\Psi(0)|^2$ scaled accordingly. Radiative Υ' widths scaled via $e_Q{}^2(M_\Upsilon/M_\psi)^{-1}$ from assumed ψ' value of 60 keV. (see refs. 78, 69, 1.) Hadronic widths scaled from $\Gamma(\psi' \to \psi + \text{hadrons}) =$ 124 keV *via* $(M_\Upsilon/M_\psi)^{-2}$. (See ref. 77.)

(9) removes one of the major reasons for suggesting that the Υ' and Υ are made of different quarks.[67] In fact, it appears *difficult* to obtain the Υ'/Υ ratio (9) if both states are ${}^3S_1$ ground states of two different $Q\bar{Q}$ pairs.

Decays of the Υ are reviewed elsewhere.[58,87] Expectations for radiative and hadronic decays of excited $b\bar{b}$ systems have been set forth in ref. 78. These should be richer than in charmonium because of the higher threshold. One can prove[84] semiclassically for an arbitrary potential that the number n_{Th} of narrow ${}^3S_1$ $Q\bar{Q}$ levels below flavor threshold is

$$n_{Th} = a(m_Q/m_c)^{1/2}; \tag{10}$$

with $a \simeq 2$ since charm threshold lies just above the second ${}^3S_1$ ($c\bar{c}$) level. For m_b/m_c between 3 and 4, $n_{Th} = 3$ or 4, corresponding to $E_{Th} \simeq 10\frac{1}{2}$ GeV (Fig. 1).

New quark flavors may be produced by photons with somewhat greater ease than in hadronic reactions. Estimates still are somewhat model-dependent,[88] but encouraging nonetheless.

D. Expectations for New Objects

Even before charm had been confirmed, it was becoming apparent that more than four leptons and the four corresponding quarks had to exist.

1. There was an indication of a new heavy lepton with $M \simeq 1.8$ GeV. The cross section $\sigma(e^+e^- \to$ "hadrons") was too large above the supposed "charm" threshold to be due to charm alone, and evidence specifically in favor of the lepton came from production of $\mu^\pm e^\mp$ pairs at SPEAR.[89]

2. The new lepton unbalanced the quark–lepton analogy that had been one of the motivations of charm. A popular means of dealing with this situation[59,90] was to introduce a new quark doublet $\begin{pmatrix} t \\ b \end{pmatrix}_L$ to go with $\begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L$ Here $e_t = 2/3$.

3. Models attempting to retain *triplets* of quarks: (u, d, s); (c, b, h) were proposed.[91] These had extra $e_Q = -1/3$ quarks. Their main justification (in retrospect) seems to have been aesthetic: some of them were based on *exceptional groups*, which had the property of limited rank and hence limited representation size.

The central question in such models seems not to be *whether* there is a sixth quark, but *what its charge and mass are*. A property of *both* models[90,91] is their tendency to introduce a *new charged lepton for every charge* $-1/3$ *quark*. Hence if a *fourth* charge $-1/3$ quark is discovered, the temptation will be great to look for a fourth lepton, and vice versa, *regardless* of the specific model.

One prediction of the mass of the sixth quark,[92] based on an eight-quark model, fills in a $t\bar{t}$ state just below 30 GeV. Within the confines of six quarks, m_t cannot be predicted, though an estimate of m_b has been made[93] using a highly appealing and economical SU(5) model.[94]

The relative strengths of weak decays of the b quark to u and c can be measured;[95] these constrain models, but don't immediately distinguish between the "quark-doublet" and "quark-triplet" alternatives. The $b \to c$ and $b \to u$ decays could provide a massive weak current, whose importance for production of new particles (such as heavy leptons) has been stressed previously.[96]

We conclude by noting that b-quarks and their likely heavier relatives can be copious sources of the long-sought Higgs bosons, both neutral[97,98] and charged.[99,100] For quark masses at the highest PETRA and PEP ranges, the prospects are encouraging if the Higgs bosons are light enough (and if they exist at all!).

IV. Weak Interactions of Heavy Quarks

A. Status of Charm

The prediction of charmed particles[101,102] has been spectacularly confirmed in the recent past. Here we review the extent to which the detailed properties expected of charm have been verified by experiment.

1. Spectrum of charmed particles

To the familiar nonets of SU(3) mesons the charm hypothesis adds a $(c\bar{c})$ particle, and an SU(3) triplet with charm $C=-1$, and an SU(3) antitriplet with $C=+1$. The pseudoscalar $D^+(c\bar{d})$ and $D^0(c\bar{u})$ mesons and their antiparticles with $m_+=1868.3\pm0.9$ MeV and $m_0=1863.3\pm0.9$ MeV are firmly established[33] and some of their properties are known. Some evidence for the existence of $F^+(c\bar{s})$ with mass of 2030 MeV, based on observations of $F^+\to\eta+$anything and $F^+\to\eta\pi^+$ has been presented by the DASP collaboration.[103] Similarly, the vector particles D^{*+} and D^{*0} with $m_+^*=2008.6\pm1.0$ MeV and $m_0^*=2006\pm1.5$ MeV are established, while the $F^*\to F\gamma$ ($m^*\simeq2140$ MeV) has been indicated by the DASP data.

For charmed baryons the experimental situation is more indefinite. The Brookhaven neutrino event,[104] $\nu p\to\mu^-\Lambda\pi^+\pi^+\pi^+\pi^-$ can be interpreted as the production of the $J^P=(1/2)^+C_1^{++}$ (cuu, with mass 2426 MeV) which decays strongly into $\pi^+C_0^+$ (cud, with mass 2250 MeV), whereupon the weak decay $C_0^+\to\Lambda\pi^+\pi^+\pi^-$ ensues. Additional evidence for the photoproduction of the charmed antibaryons $\bar{C}_0^-$, $\bar{C}_1^{--}$, $\bar{C}_1^0$, extending the earlier observations of the Columbia–Fermilab–Illinois collaboration[105] was presented by W. Lee in Parallel Session B4. No compelling observations of the $J^P=(3/2)^+$ states have been made.

2. Form of the charm-changing weak current

We expect the $\Delta C=1$ charged current to take the form

$$J_\mu(\Delta C=1)=\bar{c}\gamma_\mu(1+\gamma_5)s\cos\theta_C-\bar{c}\gamma_\mu(1+\gamma_5)d\sin\theta_C,$$

i.e., to mediate transitions within the left-handed weak isospin doublet $\begin{pmatrix} c \\ s_\theta \end{pmatrix}_L$. This form is consistent with all experimental information. To what extent is it implied by experiment?

a. V–A structure

Much indirect evidence from neutral current information, "naturalness" requirements, and the like supports a V–A space-time structure.[106] The only direct test now available comes from the interpretation of antineutrino-induced dimuon events as

$$\bar{\nu}N\to\mu^+\bar{C}+\text{anything}$$
$$\to\mu^-+\nu_\mu+\text{hadrons},$$

which proceeds *via* the elementary interaction $\bar{\nu}\bar{s}_\theta\to\mu^+\bar{c}$. The standard V–A assignment gives an excellent description of this process.[107] For infinite neutrino energy, and in the absence of experimental cuts, the possibilities $\begin{pmatrix} c \\ s_\theta \end{pmatrix}_L$ and $\begin{pmatrix} c \\ s_\theta \end{pmatrix}_R$ are easily distinguished: for left-handed $\bar{c}s_\theta$ transitions, the lepton-energy-loss distribution $d\sigma/dy$ is flat, while for right-handed transitions, $d\sigma/dy\propto(1-y)^2$. At finite energies these idealized distributions are distorted by threshold effects and experimental apparatus. Computations[108] appropriate to the experimental conditions of the CDHS experiment[109] are compared with published data in Fig. 7. It is difficult to express a preference for $V-A$ over $V+A$, and impossible to set a limit on a right-handed coupling from these data. The many dilepton events now under analysis should permit a quantitative statement to be made in the near future.

Several other tests for the chirality of the $\bar{c}s_\theta$ transition have been proposed, but not executed. Semileptonic decays of charmed baryons give rise to characteristic energy and invariant mass spectra, and hyperon polarizations, which are sensitive to VA interference.[110] For charmed mesons, the energy spectra and angular correlations in the decays $D\to K^*e\nu$ and $D\to(K+\text{pions})e\nu$ probe the spacetime structure of the hadronic current.[111,112]

b. Does the current have the GIM form?

The dominance of the $c\bar{s}$ transition is established by the observations of charmed meson

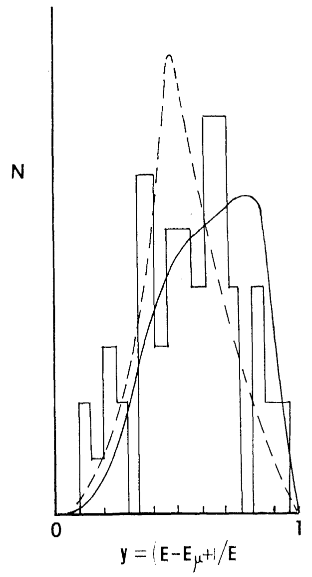

Fig. 7. Distribution of fractional leptonic energy loss $y=(E-E_{\mu^+})/E$ for di-muon events, $\bar{\nu}N\to\mu^+\mu^-X$. Data are from the CDHS experiment.[109] The curves[108] are calculated assuming the di-muon events come from charm particle decays and that the basic weak interaction is $\bar{\nu}\bar{s}_\theta\to\mu^+\bar{c}$. The curves include apparatus acceptance and threshold effects. Solid curve, $\binom{c}{s_\theta}_L$, dashed curve, $\binom{c}{s_\theta}_R$.

decays, although Cabibbo universality is not seriously tested. However, no Cabibbo-suppressed $\bar{c}d$ decays have been convincingly observed. Indeed, the only *experimental* evidence for the $c\to d+W^+$ transition comes from the observation[113] of a valence component of

$$\nu N\to\mu^- l^+ + \text{anything},$$

interpreted as

$$\nu d\to\mu^- c$$
$$\to l^+ + \text{anything},$$

at close to the expected rate.[107]

It is important to compare rates for the decays such as

$$\frac{\Gamma(D^0\to K^- l^+\nu)}{\Gamma(D^0\to\pi^- l^+\mu)}=\cot^2\theta_C,$$

$$\frac{\Gamma(D^0K^-\pi^+)}{\Gamma D^0\to\pi^+\pi^- \text{ or } K^+K^-)}=\cot^2\theta_C,$$

$$\frac{\Gamma(D^+\to\bar{K}^0 l^+\nu)}{\Gamma(D^+\to\pi^0 l^+\nu)}=\cot^2\theta_C,$$

$$\frac{\Gamma(D^+\to\bar{K}^0\pi^+)}{2\Gamma(D^+\to\pi^+\pi^0)}=\cot^2\theta_C,$$

where phase space differences have been neglected.

3. Is there a nonleptonic enhancement?

Counting quark diagrams gives $\Gamma(c\to \text{hadrons}+e\nu)=1$; $\Gamma(c\to\text{hadrons}+\mu\nu)=1$; $\Gamma(c\to\text{hadrons})=3$ (a color factor). Consequently we would expect $\Gamma(c\to\text{hadrons}+e\nu)/\text{all}=1/(1+1+3)=20\%$. The data[114] [DASPI $(8\pm2)\%$; Pb glass wall $(8.2\pm1.8)\%$; DELCO $(11\pm2)\%$] suggest instead that $\Gamma(D\to\text{hadrons}+e\nu)/\text{all}\approx10\%$. This implies that nonleptonic decays are enhanced by 8/3 in rate, which is a much smaller factor than the 20-fold enhancement of $\Delta C=0$ nonleptonic decays.[115]

4. SU(4) structure of the nonleptonic Hamiltonian

In a current-current picture, the nonleptonic Hamiltonian transforms as $\mathscr{H}_{NL}\cong$**[8]**$\oplus$**[27]** in the SU(3) Cabibbo theory, or as $\mathscr{H}_{NL}\cong$**20**$\oplus$**84** in the SU(4) GIM theory. In the SU(3) case the octet component is enhanced while the **[27]** contribution is suppressed. It is natural to suppose that the appropriate generalization to SU(4) is to suppress the **84**.[116] In the absence of an **84** piece of the Hamiltonian, the decay $D^+\to\bar{K}^0\pi^+$ and all Cabibbo-favored two-body decays of D^+ are forbidden.[117] The *rate* $\Gamma(D^+\to\bar{K}^0\pi^+)$ therefore measures the strength of the **84**-initiated transitions. The following branching ratios for nonleptonic decays of charmed mesons have been reported[118]

$D^+\to\bar{K}^0\pi^+$	$(1.5\pm0.6)\%$
$\to K^-\pi^+\pi^+$	$(3.9\pm1.0)\%$
$D^0\to K^-\pi^+$	$(2.2\pm0.6)\%$
$\to\bar{K}^0\pi^+\pi^-$	$(4.0\pm1.3)\%$
$\to K^-\pi^+\pi^0$	$(12\pm6)\%$
$\to K^-\pi^+\pi^-\pi^+$	$(3.2\pm1.1)\%$

A comparison of the relative rates for the decays $D^+\to\bar{K}^0\pi^+$ and $D^0\to K^-\pi^+$ may be had by measuring the relative lifetimes[119]

$$\frac{\tau_+}{\tau_0}=\frac{\Gamma(D^+\to\text{hadrons}+l^+\nu)/\text{all}}{\Gamma(D^0\to\text{hadrons}+l^+\nu)/\text{all}}.$$

5. Charmed-particle lifetimes

The lifetime of charmed particles is[120] $\tau=\pi^{\pm2}\times5\times10^{-13}$s. If the lifetime would exceed 10^{-11}s or be less than 10^{-14}s, our concept of charm would require dramatic revision.

B. Weak Interactions of the b-Quark

The starting point for our analysis is:

i) The b-quark exists, with a mass of about 5 GeV/c.[2]

ii) Its charge is $-1/3$.

iii) It is a member of a color triplet.

iv) It decays into u- or c-quarks.

v) It might or might not have a heavier partner t, with $e_Q=+2/3$. Specific gauge theories are discussed by Altarelli,[106] Fritzsch,[121] and Weinberg.[122]

1. *b-particle spectroscopy (and beyond)*

The table below shows the proliferation of meson and baryon states expected when a fifth (or sixth!) quark flavor is added to those already known. An argument[84] for the position of the new-flavor threshold suggests that the mass of the lightest meson $(b\bar{u})^-$ or $(b\bar{d})^0$ will be close to 5.3 GeV/c^2.

Table IV. Multiplicity of $Q\bar{Q}$ and QQQ states.

Symmetry group Quarks	SU(3) u, d, s	SU(4) $+c$	SU(5) $+b$	SU(6) $+t$
Mesons	$1\oplus 8$	$1\oplus 15$	$1\oplus 24$	$1\oplus 35$
$(1/2)^+$ Baryons	8	20	40	70
$(3/2)^+$ Baryons	10	20	35	56

2. *The b-quark lifetime*[123]

We assume that the relevant charged-current decays are $b\to u+W^-$, $c+W^-$. The possible transition $b\to t+W^-$ is shown to be energetically forbidden by the dimuon spectrum of the Columbia–Fermilab–Stony Brook experiment[54] and by the measurement of $R=\sigma(e^+e^-\to \text{hadrons})/\sigma(e^+e^-\to\mu^+\mu^-)$ at $\sqrt{s}=10.1$ GeV at DORIS.[41,58] A full-strength coupling to the u-quark leads to the decay $b\to ue\nu$, $u\mu\nu$, $u\tau\nu$, $u\bar{u}d_\theta$, $u\bar{c}s_\theta$. A free-quark model for the lifetime then leads to[124]

$$\tau_0\simeq 1.3\times 10^{-15}\text{s}$$

in the absence of nonleptonic enhancement. Because little enhancement is anticipated,[125] we expect that the b-quark lifetime $\tau_b\gtrsim\tau_0$.

It is possible that b might be uncoupled from u and c, in which case it would be absolutely stable.[126] Two searches at Fermilab[85] have found no stable charged particles in the 5 GeV/c^2 mass range at a sensitivity of $(0.1-0.2)\times$ the Υ production cross section. This implies that the lifetime of the b-quark is less than 5×10^{-8}s.

In a specific gauge model[90] with three left-handed doublets

$$\begin{pmatrix} u \\ d' \end{pmatrix}_L, \begin{pmatrix} c \\ s' \end{pmatrix}_L, \begin{pmatrix} t \\ b' \end{pmatrix}_L,$$

universality and the K_L–K_S mass difference suggest $\tau_b\gtrsim 10^{-14}$s.[125,127]

3. *Relative rates of $b\to(u, c)$ transitions*

Possible sources of b-particles are

i) $$\left.\begin{matrix}\bar{\nu}+u\to\mu^++b \\ \bar{\nu}+c\to\mu^++b\end{matrix}\right\}\lesssim 10^{-3}\times\sigma_{\text{tot}},$$

a discouraging rate, and likely to be even lower in νN collisions.

ii) Hadronic production. While the expected rates are not infinitesimal ($>10\times$ the Υ production cross section), hadronic experiments have not yet contributed to *charm* spectroscopy. This may change, and we should be especially alert for cascades of short tracks in emulsions or other high-resolution devices as signatures of $b\to c\to s$.

iii) $e^+e^-\to(b\bar{b})_{\text{unbound}}$. A $^3S_1$ upsilon level just above new-flavor threshold is expected to be produced at an appreciable rate.

Nonleptonic decays to specific final states in principle provide a measure of the relative rates for the transitions $b\to u+W^-$ and $b\to c+W^-$. In practice, small branching ratios probably cripple this approach. A more promising method would seem to be the analysis of $e^+e^-\to(b\bar{b})\to$multileptons.[124] Decays of the b-quark which lead to final-state electrons are enumerated in Table 5 below.

Table V. b-quark decay chains leading to electrons. ($b\to qW^-$, $W^-\to xy$, followed by q and xy decay)

q \ $W^-\to$	$e\nu$	$\mu\nu$	$\tau\nu$ no e^-	$\tau\nu$ e^-	$d\bar{u}$	$s\bar{c}$ no e^-	$s\bar{c}$ e^-
u	e^-	—	—	e^-	—	—	e^-
$c\to$no e^+	e^+	—	—	e^+	—	—	e^-
$c\to e^+$	e^+e^-	e^+	e^+	e^+e^-	e^+	e^+	e^+e^-

Unbound $(b\bar{b})$ therefore lead to final states containing (1) no $e^\pm$ (σ^0), (2) one $e^\pm$ ($\sigma_1=\sigma_{+}+\sigma_{-}$), (3) e^+e^- (σ_{+-}), (4) e^+e^- or e^-e^- ($\sigma_{ss}=\sigma_{++}+\sigma_{--}$), (5) $e^\pm e^+e^-$ ($\sigma_3=\sigma_{++-}+\sigma_{+--}$), or (6) $e^+e^+e^-e^-$ (σ_4). With little ambiguity, measurements of the relative cross sections determine $\Gamma(b\to u+W^-)/\Gamma(b\to c+W^-)$ and hence the relative weak couplings. The analysis has been shown to go through in the presence of neutral-particle mixing.

4. CP violation

It has been suggested by several authors[125,128] that neutral particle mixing might manifest observable CP violating effects in mesons composed of b-quarks. The magnitude of CP violation is measured by the charge asymmetries

$$(\sigma_{+}-\sigma_{-})/\sigma_{1},$$
$$(\sigma_{++}-\sigma_{--})/\sigma_{ss},$$
$$(\sigma_{++-}-\sigma_{+--})/\sigma_{3}.$$

An alternative approach to the study of neutral particle mixing and CP violation, which rests on a momentum cut to select the "primary" electron from the decay $b\rightarrow qe\nu$ has also been advocated.[128]

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