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THE PHENOMENOLOGY OF THE NEXT LEFT-HANDED QUARKS

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The observation of $\Upsilon(9.5)$ suggests that the -onium of at least one new quark has been discovered. We discuss the production and decays of the lowest-lying vector states. Recent observations have no indications of right-handed currents in antineutrino-nucleon scattering. We discuss the properties of new states made of t (charge = $\frac{2}{3}$) or b (charge = $-\frac{1}{3}$) quarks in a model with just left-handed currents. Particular attention is paid to decay modes, production by neutrinos or antineutrinos, the analogues of $K_0 - \bar{K}_0$ mixing, and *CP* violation.

To our friend Benjamin W. Lee
who cannot share with us the
joys of new discoveries.

1. Introduction

There have recently been two fundamental advances in our knowledge about quarks beyond charm. On the one hand, a number [1,2] of recent deep inelastic ν and $\bar{\nu}$ scattering experiments see no evidence for right-handed currents coupling to new quarks. On the other hand, evidence has been reported [3] for a state or states Υ with mass $\sim 9\frac{1}{2}$ GeV, produced in hadron-hadron collisions and decaying into lepton pairs. It seems very likely that the -onium of one or more new quarks has been discovered. Since such new quarks have low enough masses to have been excited in ν or $\bar{\nu}$ collisions, we interpret the absence of gross right-handed current effects as indicating that the new quark or quarks have left-handed weak interactions. The simplest model which could incorporate such quarks is a six-quark gene-

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ralization of the standard Glashow-Iliopoulos-Maiani (GIM) [4] -Weinberg-Salam-Ward [5] $SU(2)_L \times U(1)$ model which was first discussed by Kobayashi and Maskawa (KM) [6]. There is even an argument, based on the embedding of this model in a unified SU(5) [7] theory including strong interactions, that the next charge $-\frac{1}{3}$ quark should [8] have a mass in the range 4 to 10 GeV. We are strongly aware that this model is in apparent conflict with present limits on parity violation in atomic physics [9], but adopt a "wait and see" attitude.

Previous phenomenological analyses [10–12] of new quarks have mainly been based on models with right-handed currents. Previous analyses [13,14] of the KM model [6] have concentrated on its implications for *CP* violation in the light (u, d, s, c) quark sector. In this paper we analyse more thoroughly the decays and production of bound states of the heavy bottom b (charge = $-\frac{1}{3}$) and top t (charge = $+\frac{2}{3}$) quarks, with a view to present and forthcoming experiments.

We start in sect. 2 with some estimates of the decay modes and production by hadron-hadron and photon-hadron collisions of the lightest orthobottom- and orthotoponia [15], which do not depend on the details of the weak interaction model. In contrast to charmonium, where we expect

$$\frac{\sigma(\mathrm{p}+\mathrm{p}\to\psi'+\mathrm{X})\,B(\psi'\to\mu^+\mu^-)}{\sigma(\mathrm{p}+\mathrm{p}\to\mathrm{J}/\psi+\mathrm{X})\,B(\mathrm{J}/\psi\to\mu^+\mu^-)}\approx(2\tfrac{1}{2}-5)\%\,, \tag{1.1}$$

not inconsistent with experiment [16], we find that for bottomonia $\mathrm{J_B}\equiv \mathrm{b\bar{b}}:1^3\mathrm{S}_1$, $\mathrm{J'_B}\equiv\mathrm{b\bar{b}}:2^3\mathrm{S}_1$, etc.:

$$\sigma(\mathrm{pp}\to\mathrm{J_B}+\mathrm{X})\,B(\mathrm{J_B}\to\mu^+\mu^-):\sigma(\mathrm{pp}\to\mathrm{J'_B}+\mathrm{X})\,B(\mathrm{J'_B}\to\mu^+\mu^-):$$
$$:\sigma(\mathrm{pp}\to\mathrm{J''_B}+\mathrm{X})\,B(\mathrm{J''_B}\to\mu^+\mu^-)\approx1:0.30:0.15\,, \tag{1.2}$$

whereas for toponia $\mathrm{J_T},\mathrm{J'_T},\ldots\equiv\mathrm{t\bar{t}}:1^3\mathrm{S}_1,2^3\mathrm{S}_1,\ldots$ we find

$$\sigma(\mathrm{pp}\to\mathrm{J_T}+\mathrm{X})\,B(\mathrm{J_T}\to\mu^+\mu^-):\sigma(\mathrm{pp}\to\mathrm{J'_T}+\mathrm{X})\,B(\mathrm{J_T}\to\mu^+\mu^-):$$
$$:\sigma(\mathrm{pp}\to\mathrm{J''_T}+\mathrm{X})\,B(\mathrm{J''_T}\to\mu^+\mu^-)\approx1:0.17:0.05\,. \tag{1.3}$$

Thus more than one bottomonium state may be visible in hadron-hadron collisions, but since their mass differences are expected to be O(400) MeV, distinguishing them would require a high resolution experiment. One of the reasons for the high rates (1.2, 1.3) for $\mathrm{J'_B}$ or $\mathrm{J'_T}$ production is that we estimate a very small branching ratio for $\mathrm{J'_B}\to\mathrm{J_B}\pi\pi$ or $\mathrm{J'_T}\to\mathrm{J_T}\pi\pi$. We also estimate the production cross section in photoproduction to be

$$\sigma(\gamma+\mathrm{p}\to\mathrm{J_B}+\mathrm{X})\,BB(\mathrm{J_B}\to\mu^+\mu^-)\approx(\tfrac{1}{10}\text{ to }1)\text{ pb}\,, \tag{1.4}$$

with $\sigma(\gamma+\mathrm{p}\to\mathrm{J_T}+\mathrm{X})\,BB(\mathrm{J_T}\to\mu^+\mu^-)=O(10)$ larger, for m_b or $m_\mathrm{t}\sim5$ GeV and

Table 1
Bottom and top production by $\bar{\nu}$ and ν ★

Bottoms

Fundamental subprocess	Production strength	Fraction of total cross section		Multimuon signals 2μ	3μ	4μ
$\bar{\nu} + u \to \mu^+ + b$, $b \to c$	$Vs_1^2 s_3^2 (1-y)^2$	$<3 \times 10^{-3}$	$<\frac{1}{2}\%$	$+-$ / $<1 \times 10^{-3}$	$+-+$ / $<1.5 \times 10^{-4}$	$+-+-$ / $<2 \times 10^{-5}$
$\bar{\nu} + c \to \mu^+ + b$ ($\bar{c}$ spectator), $b \to c$	$C(s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta) \times (1-y)^2$	$<2 \times 10^{-3}$		$++$ / $<7 \times 10^{-4}$	$+--$ / $<5 \times 10^{-5}$	
$\nu + \bar{u} \to \mu^- + \bar{b}$, $\bar{b} \to \bar{c}$	$Ss_1^2 s_3^2 (1-y)^2$	$<3 \times 10^{-5}$	$<7 \times 10^{-4}$	$-+$ / $<2 \times 10^{-4}$	$-+-$ / $<3 \times 10^{-5}$	$-+-+$ / $<3 \times 10^{-6}$
$\nu + \bar{c} \to \mu^- + \bar{b}$ (c spectator), $\bar{b} \to \bar{c}$	$C(s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta) \times (1-y)^2$	$<7 \times 10^{-4}$		$--$ / $<1 \times 10^{-4}$	$-++$ / $<1.5 \times 10^{-5}$	

Tops

Fundamental subprocess	Production strength	Fraction of total cross section		Multimuon signals 2μ	3μ	4μ
$\bar{\nu} + \bar{d} \to \mu^+ + \bar{t}$	$Ss_1^2 s_2^2$	$<1.5 \times 10^{-4}$	$<5 \times 10^{-3}$	$+-$ / $<7 \times 10^{-4}$		
$\bar{\nu} + \bar{s} \to \mu^+ + \bar{t}$	$S(s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta)$	$<5 \times 10^{-3}$		$++$ / $<2 \times 10^{-4}$		
$\nu + d \to \mu^- + t$	$Vs_1^2 s_2^2$	$<3 \times 10^{-3}$	$<8 \times 10^{-3}$	$-+$ / $<10^{-3}$		
$\nu + s \to \mu^- + t$	$S(s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta)$	$<5 \times 10^{-3}$		$--$ / $<3 \times 10^{-4}$		

★ In the upper (lower) half of the above table we assume the b (t) quark to be the lighter one. The like sign dimuon events due to t production require very substantial $T-\bar{T}^0$ mixing.

$E_\gamma \gtrsim 100$ GeV. The photoproduction of dissociated pairs of naked top or bottom states for $E_\gamma \gtrsim 100$ GeV is expected to be O(1–10) nb.

In sect. 3 we proceed to estimates of the decay characteristics and neutrino production of naked top and bottom states. We start by bounding the various angles appearing in the KM weak coupling matrix [6,14]. We then find that if bottom particles are lighter than tops, charmed particles should predominate in their decays. As suggested by previous authors [11,12], both $q\bar{q}$ and $q\bar{q}q\bar{q}$ final states are expected to be significant in bottom and top decays. Because of small mixing angles to light quarks and a small enhancement of non-leptonic decays [17], the lifetimes of the lightest naked bottom or top states are expected to be $\gtrsim 10^{-13}$ sec, so that they may leave visible emulsion tracks. We discuss the production of bottom and top states by ν and $\bar{\nu}$. The production of bottoms by $\bar{\nu}$ is expected to be between $\frac{1}{2}$ % and 10^{-5} (!) of the total cross section. Thus no dramatic threshold is expected in the total cross section, but interesting multilepton signatures may exist ($\mu^+\mu^+$, $\mu^+\mu^-\mu^+$, $\mu^+\mu^-\mu^-$, $\mu^+\mu^-\mu^+\mu^-$). Our conclusions on b and t production by ν and $\bar{\nu}$ are summarized in table 1. We also discuss the mass matrices for new neutral mesons ($B^0 \equiv b\bar{d}$, $T^0 \equiv t\bar{u}$) analogous to the ($K^0 - \bar{K}^0$) and ($D^0 - \bar{D}^0$) systems. We find

$$\left(\frac{\Delta m}{\Gamma}\right)_{B^0} \approx \frac{m_t^2}{700\ \mathrm{GeV}^2}\,, \qquad \left(\frac{\Delta m}{\Gamma}\right)_{T^0} \approx \frac{m_b^2}{2001\ \mathrm{GeV}^2}\,. \tag{1.5}$$

$$\left(\frac{\Delta\Gamma}{2\Gamma}\right)_{B^0(m_t>m_b)} \approx \tfrac{1}{12}\,, \qquad \left(\frac{\Delta\Gamma}{2\Gamma}\right)_{T^0(m_b>m_t)} \approx \tfrac{1}{50}\,.$$

If $m_t > m_b$, $B^0 - \bar{B}^0$ mixing is expected to be larger than for $D^0 - \bar{D}^0$ ($> \frac{1}{2}$ %) and becomes sensitive to the mass of the t-quark if it is $\gtrsim 8$ GeV.

We also analyse *CP* violation in the $B^0 - \bar{B}^0$ and $T^0 - \bar{T}^0$ systems: in both cases:

$$\frac{\mathrm{Im}\, m}{\Delta m} \approx \tan 2\delta\,, \tag{1.6}$$

where δ is the KM [6] phase which can be much larger than the 10^{-3} characteristic of *CP* violation in the $K^0 - \bar{K}^0$ system.

Sect. 4 summarizes our conclusions, and assesses the prospects for discriminating between possible assignments of the $\Upsilon(9.5)$ GeV as bottomonium, or toponium, or a combination of the two. We have a very slight preference for believing it to be bottomonium, but it is absurdly premature to draw this inference.

2. Decay modes and production of hidden top and bottom states

In this section we summarize some guesses about the decay modes and production mechanisms of orthobottomonia and orthotoponia [11,12,15], with particular emphasis on the ratios of 1^3S_1, 2^3S_1, and 3^3S_1 signals in lepton-pair production by hadron-hadron collisions.

2.1. Decay modes

Eichten and Gottfried [15] have recently applied their successful charmonium model to heavier quarks, estimating mass differences between $^3S_1$ states. For a quark with mass ~5 GeV they get

$$m(2^3S_1) - m(1^3S_1) \approx 420 \text{ MeV}, \qquad m(3^3S_1) - m(2^3S_1) \approx 330 \text{ MeV} \tag{2.1}$$

and conclude that the 3^3S_1 is below the threshold for producing pairs of bound states of new and light quarks. It may therefore have a substantial branching ratio into lepton pairs, and three spikes might therefore be visible in $p + p \to \ell^+\ell^- + X$.

Their estimates of the decay modes of these three states are shown in table 2. The only additions to their results that we have made are to have divided by four their estimates of radiative decays to apply to the bottomonium case, and to have estimated the decay modes $J'_B \to J_B\pi\pi$, etc.

If a $V' \to V\pi\pi$ matrix element

$$g_{V'V}\epsilon'_\mu\epsilon^\mu \tag{2.2}$$

is introduced, and $z \equiv m^2_{\pi\pi}/(m' - m)^2$: $m_{V'} \equiv m'$, $m_V \equiv m$, then

$$\frac{d\Gamma}{dz} = \frac{1}{64\pi^2}\frac{g^2}{4\pi}\frac{(m'-m)^3(m'+m)}{m'^3}\sqrt{1 - \frac{4m_\pi^2}{(m'-m)^2 z}}\sqrt{1-z} \times \sqrt{1 - \left(\frac{m'-m}{m'+m}\right)^2 z}\,. \tag{2.3}$$

Taking

$$m = 9.5 \text{ GeV}, \qquad m' = 9.5 \text{ GeV} + 420 \text{ MeV}, \tag{2.4}$$

we get

$$\frac{\Gamma(V' \to V\pi\pi)}{\Gamma(\psi' \to J/\psi\pi\pi)} = \left(\frac{g^2_{V'V}}{g^2_{\psi' J/\psi}}\right) \times 2.5\%\,. \tag{2.5}$$

Assuming that $|g_{V'V}| \leqslant |g_{\psi' J/\psi}|$, as seems reasonable, and using the experimental value of $\Gamma(\psi' \to J/\psi\pi\pi)$, we find from eq. (2.5) that

$$\Gamma(V' \to V\pi\pi) < 3 \text{ keV}\,. \tag{2.6}$$

The estimate (2.6) is small by comparison with the direct hadron decays shown in table 2, but may actually be a gross overestimate. Various authors [18] have pointed out that chiral symmetry would imply an Adler zero in the $V'V\pi\pi$ coupling, and this is consistent with the dipion spectrum observed in $\psi' \to J/\psi + \pi\pi$ decay. If we

Table 2
Decay modes of bottom- and toponia

Mode	State					
	J_B	J'_B	J''_B	J_T	J'_T	J''_T
e^+e^-	0.7	0.4+	0.4−	2.7	1.7	1.5
$\mu^+\mu^-$	0.7	0.4+	0.4−	2.7	1.7	1.5
$\tau^+\tau^-$	0.7	0.4+	0.4−	2.7	1.7	1.5
$\gamma^* \to$ hadrons	2.8	1.7	1.5	10.8	6.8	6.0
direct hadrons	13.7	8.8	7.8	13.7	8.8	7.8
$\gamma\chi_{B,T}$		8	12.5		30	50
$J_{B,T}\pi\pi$		$\lesssim 0.3$	small ?		$\lesssim 0.3$	small ?
$J'_{B,T}\pi\pi$			$\lesssim 0.1$			$\lesssim 0.1$
Total	~19	~20	~23	~33	~51	~68
BR $\to \mu^+\mu^-$	3.5%	2.1%	1.6%	8.2%	3.3%	2.2%

Decay widths in keV

take a matrix element

$$\tilde{g}_{V'V}\epsilon'_\mu \epsilon^\mu [(p_{\pi_1} + p_{\pi_2})^2 - nm_\pi^2] , \tag{2.7}$$

then the value [18] $n \approx 4$ would imply that

$$\frac{\Gamma(V' \to V\pi\pi)}{\Gamma(\psi' \to J/\psi\pi\pi)} \approx \frac{\tilde{g}^2_{V'V}}{\tilde{g}^2_{\psi' J/\psi}} \times 3 \times 10^{-3} , \tag{2.8}$$

which is very small for $|\tilde{g}_{V'V}| \lesssim |\tilde{g}_{\psi' J/\psi}|$, as we expect. If we use expression (2.8), table 2 shows that we estimate

$$B(J_B \to \mu^+\mu^-) \approx 3\tfrac{1}{2}\ \% , \qquad B(J'_B \to \mu^+\mu^-) \approx 2\% ,$$

$$B(J''_B \to \mu^+\mu^-) \approx 1\tfrac{1}{2}\ \% , \tag{2.9a}$$

$$B(J_T \to \mu^+\mu^-) \approx 8\% , \qquad B(J'_T \to \mu^+\mu^-) \approx 3\% ,$$

$$B_T(J''_T \to \mu^+\mu^-) \approx 2\% . \tag{2.9b}$$

2.2. *Production in hadron-hadron collisions*

Before estimating J_B, J'_B, ..., J_T, J'_T, ... production, we first warm up on charmonium. The mechanism whereby J/ψ and ψ' are produced in hadron-hadron collisions is obscure. The absence of extra muons [19] produced in association with

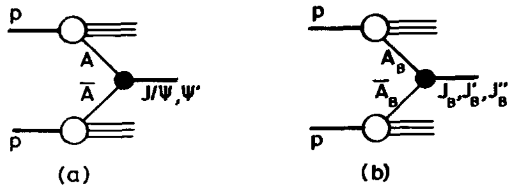

Fig. 1. Generalized Drell-Yan mechanisms for the production of charmonium (a) and bottomonium (b) states.

the J/ψ suggests that they are not made by $c\bar{c}$ quark fusion. The inequality of $\sigma(p + p \to J/\psi + X)$ and $\sigma(\bar{p} + p \to J/\psi + X)$ [20] suggests that J/ψ is not made by gluon amalgamation [21] alone. The small ratio $\sigma(\bar{p} + p \to J + X)/\sigma(p + p \to J + X)$ suggests that J/ψ is not made by non-strange $q\bar{q}$ annihilation [22] alone. But let us assume [23] (i) that the J/ψ and ψ' are made by a similar Drell-Yan [24] production mechanism, $A\bar{A} \to J/\psi, \psi'$ for some constituents A illustrated in fig. 1a, and (ii) that the decays $J/\psi \to A\bar{A}$, $\psi' \to A\bar{A}$ are responsible for fractions F, F' of the total J/ψ, ψ' decay widths, respectively. Then

$$\frac{\sigma(pp \to \psi' + X)\, B(\psi' \to \mu^+\mu^-)}{\sigma(pp \to J/\psi + X)\, B(J/\psi \to \mu^+\mu^-)} = \left(\frac{m_{J/\psi}}{m_{\psi'}}\right)^3 \frac{\mathcal{L}(m_{\psi'}/\sqrt{s})}{\mathcal{L}(m_{J/\psi}/\sqrt{s})}$$

$$\times \frac{F'\Gamma(\psi' \to \mu^+\mu^-)}{F\Gamma(\psi/J \to \mu^+\mu^-)}\,, \tag{2.10}$$

where $\mathcal{L}(m/\sqrt{s})$ is the luminosity for $A\bar{A}$ collisions. If we identify $\Gamma(J/\psi, \psi' \to A\bar{A})$ with decay widths into conventional hadrons, then

$$F \simeq 0.86\,, \qquad F' \approx 0.1 \text{ to } 0.2\,. \tag{2.11}$$

Examination of the excitation function for J/ψ production suggests that

$$\mathcal{L}\left(\frac{m_{\psi'} = 3.7}{\sqrt{s} = 27}\right) \Big/ \mathcal{L}\left(\frac{m_{J/\psi} = 3.1}{\sqrt{s} = 27}\right) \approx \tfrac{3}{4}\,. \tag{2.12}$$

Inserting (2.11) and (2.12) into (2.10) gives [25]

$$\frac{\sigma(pp \to \psi' + X)\, B(\psi' \to \mu^+\mu^-)}{\sigma(pp \to J/\psi + X)\, B(J/\psi \to \mu^+\mu^-)} \approx \left(\frac{3.1}{3.7}\right)^3 \times \frac{3}{4} \times \frac{(0.1 \text{ to } 0.2) \times 2.1}{0.86 \times 4.8}$$

$$= (2.5 \text{ to } 5)\%\,. \tag{2.13}$$

The ratio (2.13) is not too dissimilar from experiment, and indicates *no need to invoke the predominance of intermediate P_c/χ production* to account for the observed small rate of ψ' production relative to J/ψ production. This finding contrasts with the models of ref. [21], which get most of their J/ψ production from gluon amalgamation to P_c/χ states. Armed with this conclusion, we now estimate

$$\frac{\sigma(pp \to J'_B + X)\, B(J'_B \to \mu^+\mu^-)}{\sigma(pp \to J_B + X)\, B(J_B \to \mu^+\mu^-)} \approx \left(\frac{m_{J_B}}{m_{J_{B'}}}\right)^3$$

$$\times \frac{\mathcal{L}_B(m_{J_{B'}}/\sqrt{s})\, F'_B \Gamma(J'_B \to \mu^+\mu^-)}{\mathcal{L}_B(m_{J_B}/\sqrt{s})\, F_B \Gamma(J_B \to \mu^+\mu^-)} \,, \tag{2.14}$$

where the J_B and J'_B are supposed to be made by a similar production mechanism $A_B\bar{A}_B \to J_B, J'_B$ illustrated in fig. 1b, with a corresponding luminosity $\mathcal{L}_B(m/\sqrt{s})$, and F_B, F'_B are defined analogously to the F, F' defined above. From table 2 [15] we have

$$F_B \approx 0.89\,, \qquad F_{B'} \approx 0.53\,, \qquad \frac{\Gamma(J_{B'} \to \mu^+\mu^-)}{\Gamma(J_B \to \mu^+\mu^-)} \approx 0.6\,.$$

If we further assume that

$$\frac{\mathcal{L}_B(m_{J_{B'}}/\sqrt{s})}{\mathcal{L}_B(m_{J_B}/\sqrt{s})} \approx \frac{\mathcal{L}(m_{J_{B'}}/\sqrt{s})}{\mathcal{L}(m_{J_B}/\sqrt{s})} \,, \tag{2.16}$$

so that the threshold behaviour for J_B production is proportional to that for J/ψ production (we do not assume at this stage that $\mathcal{L}_B = \mathcal{L}$), then we find

$$\frac{\sigma(pp \to J'_B + X)\, B(J'_B \to \mu^+\mu^-)}{\sigma(pp \to J_B + X)\, B(J_B \to \mu^+\mu^-)} \approx 0.8 \times \tfrac{0.53}{0.89} \times 0.6 \approx 30\%\,. \tag{2.17}$$

Similar analyses can be performed for J''_B, and for the J_T, J'_T, J''_T sequence using other figures in table 2: they give the results quoted in eqs. (1.2) and (1.3). One caveat should be mentioned: we have not included in these figures cascade decays: $J'_B \to \chi_B + \gamma$, $\chi_B \to J_B + \gamma$, etc. In the case of b (charge $= -\frac{1}{3}$), we do not expect these to make big corrections to the ratio (2.17) (cf. the relatively small $J'_B \to \chi_B + \gamma$ branching ratio in table 2), but these feedthroughs might be significant for J_T, J'_T, and J''_T, depressing further the rates for J'_T and J''_T relative to Υ. We note that Herb et al. [3] find that a good fit to the data is obtained with two peaks with relative production rates 1 : 0.4. The assumption (2.16) would yield an increase in the Υ by (10–30) at the ISR as compared with a 400 GeV proton beam at FNAL or SPS. This is compatible with present limits [3].

Finally, we test the hypothesis that the constituents (gluons? ordinary light $q\bar{q}$?) which collide to form J/ψ and ψ' are the same as those that make the $\Upsilon(9.5)$. So we now assume

$$\mathcal{L}(m/\sqrt{s}) = \mathcal{L}_B(m/\sqrt{s})\,, \tag{2.18}$$

in which case we have

$$\frac{\sigma(\mathrm{pp} \to \mathrm{J_B} + \mathrm{X})\, B(\mathrm{J_B} \to \mu^+\mu^-)}{\sigma(\mathrm{pp} \to \mathrm{J}/\psi + \mathrm{X})\, B(\mathrm{J}/\psi \to \mu^+\mu^-)} \approx \left(\frac{m_{\mathrm{J}/\psi}}{m_{\mathrm{J_B}}}\right)^3 \frac{\mathcal{L}(m_{\mathrm{J_B}}/\sqrt{s})}{\mathcal{L}(m_{\mathrm{J}/\psi}/\sqrt{s})}$$

$$\times \frac{\Gamma(\mathrm{J_B} \to \mu^+\mu^-)}{\Gamma(\mathrm{J}/\psi \to \mu^+\mu^-)} \tag{2.19}$$

$$\approx \mathrm{O}(10^{-4})\,, \tag{2.20}$$

if we use table 2 for $\Gamma(\mathrm{J_B} \to \mu^+\mu^-)$ and eyeball the J/ψ excitation function to estimate the ratio of luminosities. The ratio corresponding to (2.19) for $\mathrm{J_T}$ production would be a factor of 4 higher. Experimentally,

$$\frac{\sigma(\Upsilon)\, B(\Upsilon \to \mu^+\mu^-)}{\sigma(\mathrm{J}/\psi)\, B(\mathrm{J}/\psi \to \mu^+\mu^-)} \approx 3 \times 10^{-5}\,, \tag{2.21}$$

so that the hypothesis (2.18) does not grossly overestimate the cross section for Υ production – the agreement would be worse if the Υ were to be identified with $\mathrm{J_T}$. At this point we may perhaps express a slight preference for the hypothesis that the Υ is bottomonium $\mathrm{J_B}$ rather than toponium $\mathrm{J_T}$, based on the results (1.2), and the smallness of the cross section. However, any such inference is grossly premature. In particular, it is quite possible that the top and bottom might be so close that their -onia would overlap.

2.3. Photoproduction cross sections

Various formulae and empirical rules exist in the literature which enable one to estimate the photoproduction of Υ if one adopts the bottomonium or toponium hypothesis, as well as the cross section for producing pairs of naked top or bottom particles. Empirically, one has

$$\sigma(\rho\mathrm{N}) : \sigma(\phi\mathrm{N}) : \sigma(\mathrm{J}/\psi\mathrm{N}) \approx \frac{1}{m_\rho^2} : \frac{1}{m_\phi^2} : \frac{1}{m_{\mathrm{J}/\psi}^2}\,, \tag{2.22}$$

and we may extend this to Υ:

$$\frac{\sigma(\Upsilon\mathrm{N})}{\sigma(\mathrm{J}/\psi\mathrm{N})} \approx \frac{m_{\mathrm{J}/\psi}^2}{m_\Upsilon^2}\,. \tag{2.23}$$

One may then assume [26] that the intermediate states in ΥN scattering almost always contain naked b or t particles, in which case

$$\frac{\sigma(\gamma \to \mathrm{b\bar{b}})}{\sigma(\gamma \to \mathrm{c\bar{c}})} \sim \left(\frac{m_{\mathrm{J}/\psi}^2}{m_{\mathrm{J_B}}^2}\right)\left(\frac{g_{\mathrm{J}/\psi}^2}{g_{\mathrm{J_B}}^2}\right)\,, \tag{2.24}$$

where the g's are defined by $\langle\gamma|V\rangle = em_V^2/g_V$, so that the observed $J/\psi \to \mu^+\mu^-$ decay width and the $J_B \to \mu^+\mu^-$ decay width estimated in table 2 give $g_{J/\psi}^2/g_{J_B}^2 = (m/\Gamma_{ee})_{J/\psi}/(m/\Gamma_{ee})_{J_B} = 4.5 \times 10^{-2}$, implying (if $\Upsilon \equiv J_B$)

$$\frac{\sigma(\gamma \to b\bar{b})}{\sigma(\gamma \to c\bar{c})} \approx 4 \times 10^{-3} \ . \tag{2.25}$$

If $\sigma(\gamma \to c\bar{c}) \approx 1\ \mu b$, as is often believed, then

$$\sigma(\gamma \to b\bar{b}) \sim 4 \text{ nb} \ . \tag{2.26}$$

The cross section for $\gamma \to t\bar{t}$ would be a factor four larger if $\Upsilon \equiv J_T$. An alternative method of estimating $\sigma(\gamma \to b\bar{b})$ comes from the sum rule of Shifman et al. [27]:

$$\int_{\nu_0}^{\infty} \frac{d\nu}{\nu^2} \sigma(\gamma \to b\bar{b}) = \frac{22\pi}{405} \rho \frac{\alpha}{4} \frac{\alpha_s(m_b^2)}{m_b^4} \ , \tag{2.27}$$

where $\rho \approx \frac{1}{2}$ is the fraction of the proton momentum carried by gluons, and $\alpha_s(m_b^2)$ is the appropriate QCD coupling strength. The kinematic threshold for $b\bar{b}$ production is

$$\nu_0 = 2m_b(m_b + m_p) \ . \tag{2.28}$$

If we approximate the threshold behaviour by a step function well above the kinematic limit,

$$\sigma(\gamma \to b\bar{b}) \equiv \sigma_\gamma^b \theta(\nu - 4m_b^2) \ , \tag{2.29}$$

then we estimate

$$\sigma_\gamma^b \approx \frac{11\pi}{405} \alpha \frac{\alpha_s(m_b^2)}{m_b^2} \approx 2 \text{ nb} \ . \tag{2.30}$$

The agreement between (2.26) and (2.30) is encouraging, and leads us to expect

$$\sigma(\gamma \to b\bar{b}) = O(1 \text{ to } 10) \text{ nb} \qquad \text{for } E_\gamma \gtrsim 100 \text{ GeV} \ , \tag{2.31}$$

if $\Upsilon \equiv J_B$, with $\sigma(\gamma \to t\bar{t}) \sim$ four times higher if $\Upsilon \equiv J_T$.

Assuming that the photoproduction of Υ is the diffractive shadow of $b\bar{b}$ or $t\bar{t}$ production [26], we estimate

$$\frac{\sigma(\gamma + p \to J_B + p)}{\sigma(\gamma \to b\bar{b})} \approx \left(\frac{e^2}{g_B^2}\right)^{-1} \frac{1}{16\pi b} \sigma(\gamma \to b\bar{b}) \ , \tag{2.32}$$

where b is the slope of the diffractive peak. Putting in the estimates (2.26) or (2.30) for $\sigma(\gamma \to b\bar{b})$, and taking $b = 2$ GeV^{-2}, yields

$$\sigma(\gamma + p \to \Upsilon + X) \approx \sigma(\gamma + p \to \Upsilon + p) = O(1 \text{ to } 10) \text{ pb} \ , \tag{2.33}$$

if $\Upsilon \equiv J_B$, with $\sigma(\gamma + p \to \Upsilon + X)$ a factor 4 higher if $\Upsilon \equiv J_T$. The ratio of J'_B to J_B photoproduction may be expected to be $O(\frac{1}{2})$ on the basis of (2.32).

3. The production and decays of new quarks with left-handed weak currents

After the discussion of the previous section, which made no reference to the weak interactions of the new quarks, we now specialize to the model with six left-handed quarks which generalizes the GIM [4] -Weinberg-Salam-Ward [5] model, and was first written down by Kobayashi and Maskawa [6]. A clairvoyant gauge theorist of the 1940's could have used the discovery of the muon to predict the existence of some extra weak couplings beyond

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L , \qquad \begin{pmatrix} u \\ d \end{pmatrix}_L \quad \text{and} \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L . \tag{3.1}$$

They were needed to cancel [28] the triangle anomalies which would otherwise have destroyed [29] the renormalizability of the weak and electromagnetic gauge theory. Nature chose to supplement (3.1) by a doublet [4]

$$\begin{pmatrix} c \\ s \end{pmatrix}_L , \tag{3.2}$$

which had the bonus of naturally suppressing $\Delta S = 1$ neutral currents to $O(G_F^2)$. Similarly, it has been commonplace to use the discovery of the heavy lepton τ, which probably has a weak coupling

$$\begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L \tag{3.3}$$

to predict the existence of a third left-handed quark doublet

$$\begin{pmatrix} t \\ b \end{pmatrix}_L , \tag{3.4}$$

which would retain the natural suppression of $\Delta S = 1$ and $\Delta C = 1$ effects to $O(G_F^2)$, as observed. A model combining (3.1) to (3.4) would preserve the phenomenological successes of the $SU(2)_L \times U(1)$ GIM-Weinberg-Salam-Ward model in ν and $\bar{\nu}$ scattering off hadrons (both by charged and neutral currents) and in purely leptonic interactions. The only phenomenological problem is the apparently small violation of parity in atomic physics [9]. The six-quark model has the bonus of giving *CP* violation in a natural way, as was first pointed out by Kobayashi and Maskawa [6].

3.1. The weak coupling matrix

The quark doublets in (3.1), (3.2), and (3.4) mix in general, and the charged weak current in the six left-handed quark model can be written as a unitary matrix with four parameters: three Euler angles generalizing the conventional Cabibbo angle θ_c, and a *CP*-violating phase [6]. The current can be written in the form:

$$J_\mu^- = (\bar{u}, \bar{c}, \bar{t})\, \gamma_\mu L \begin{pmatrix} c_1 & -s_1 c_3 & -s_1 s_3 \\ s_1 c_2 & c_1 c_2 c_3 - s_2 s_3 e^{i\delta} & c_1 c_2 s_3 + s_2 c_3 e^{i\delta} \\ s_1 s_2 & c_1 s_2 c_3 + c_2 s_3 e^{i\delta} & c_1 s_2 s_3 - c_2 c_3 e^{i\delta} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix}, \quad (3.5)$$

where $c_i(s_i) \equiv \cos\theta_i\,(\sin\theta_i)$, $i = 1, 2, 3$. In the limit $\theta_1 \to \theta_c$, $\sin\theta_2$, $\sin\theta_3 \to 0$ the KM current (3.5) reduces to the GIM current. In a previous analysis [14], it was shown that the experimental validity of universality between μ decay and the sum of u couplings to d and s quarks restricts the "leakage" of the u coupling to the b-quark: $s_1^2 s_3^2 < 0.003$. Since $s_1^2 \simeq 0.05$, this means that

$$s_3^2 < 0.06\,. \quad (3.6)$$

What other restrictions exist on the angles in (3.5)?

Gaillard and Lee [30] estimated

$$\langle \bar{K}^0 | -\mathcal{L}_{\text{eff}} | K^0 \rangle \approx \frac{G_F}{\sqrt{2}} f_K^2 m_K^2 \frac{\alpha}{4\pi} \left(\frac{m_c}{38\ \text{GeV}}\right)^2 \sin^2\theta_c \cos^2\theta_c \quad (3.7)$$

and the observed $K_L - K_S$ mass difference led them to estimate $m_c \sim 2$ GeV. The experimental success of this estimate leads us to require its validity in the KM model, where eq. (3.7) gets modified:

$$m_c^2 \sin^2\theta_c \cos^2\theta_c \to s_1^2 c_1^2 c_3^2 \left[c_2^4 m_c^2 + s_2^4 m_t^2 + \frac{2 s_2^2 c_2^2 m_t^2 m_c^2}{m_t^2 - m_c^2} \ln \frac{m_t^2}{m_c^2} \right]. \quad (3.8)$$

We therefore have a restriction on s_2^2 in terms of $\eta \equiv m_c^2/m_t^2$:

$$s_2^2 < \eta \ln \eta + \sqrt{(\eta \ln \eta)^2 + \eta}\,. \quad (3.9)$$

Putting in extreme values for m_t yields the bounds

$$\begin{aligned} m_t &= 5\ \text{GeV} && \to s_2^2 < 0.2\,, \\ m_t &\approx m_W \approx 65\ \text{GeV} && \to s_2^2 < 0.03\,. \end{aligned} \quad (3.10)$$

On the other hand, if the KM phase δ is to be responsible for the CP violation in the $K^0 - \bar{K}^0$ system, eq. (2.11) of ref. [14] implies that for $s_i^2 << c_i^2$ ($i = 1, 2, 3$), and $m_c^2 << m_t^2$:

$$\frac{1}{2} \left| \frac{\operatorname{Im} M_{12}^K}{\Delta m^K} \right| \approx 10^{-3} \approx s_2 s_3 \sin\delta \left[-\ln\eta - 1 + \frac{s_2^2}{\eta} \right]. \quad (3.11)$$

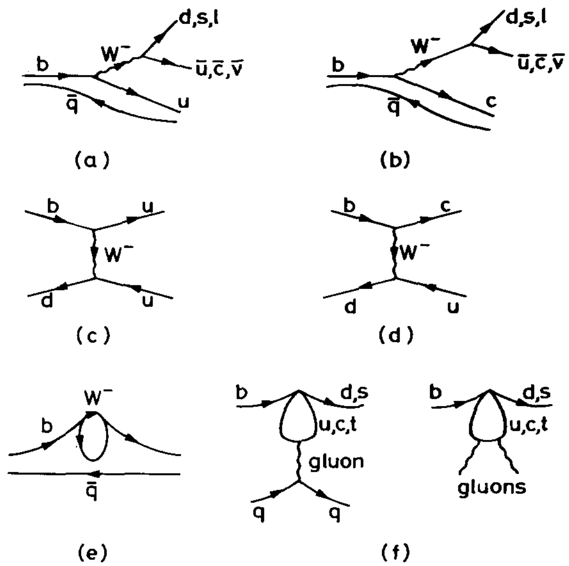

Fig. 2. Quark diagrams contributing to B decay in the free quark model [(a)–(d)], and in the presence of strong interactions [(e) and (f)].

The bounds (3.6) and (3.10) imply that the KM phase δ cannot be arbitrarily small:

$$m_t = 5\ \text{GeV} \rightarrow |\sin\delta| > 3 \times 10^{-3}\ ,$$

$$m_t = 65\ \text{GeV} \rightarrow |\sin\delta| > 6 \times 10^{-4}\ . \tag{3.12}$$

We still [14] leave the derivation of the small numbers (3.6), (3.10), and (3.12) as an exercise for ourselves and our readers.

3.2. Decays of bottom and top particles

We now turn to the weak decays of top ($t\bar{q}$: q = u, d) and bottom ($b\bar{q}$) mesons. Because of the symmetry of the KM matrix under the substitutions

$$t \leftrightarrow b\ , \qquad c \leftrightarrow s\ , \qquad u \leftrightarrow d\ , \qquad \theta_2 \leftrightarrow \theta_3\ , \tag{3.13}$$

we just consider the decays of bottom particles in the case $m_b < m_t$: the properties of top particles in the case $m_t < m_b$ can be obtained by using (3.13).

There are three classes of diagrams which may contribute to bottom decays, which are shown in fig. 2. We first discuss figs. 2a and 3b which contain four-quark operators and $q\bar{q}q\bar{q}$ final states. The operators should be separated into symmetric and antisymmetric pieces, which have different anomalous dimensions, so that the antisymmetric piece is enhanced [31]. The short-distance enhancement factor for t-decays is

$$A_t \approx \left[\frac{\ln(m_W^2/\Lambda_6^2)}{\ln(m_t^2/\Lambda_6^2)}\right]^{12/21}, \tag{3.14}$$

while for lower mass quarks:

$$A_s = \left[\frac{\ln(m_c^2/\Lambda_3^2)}{\ln(\mu^2/\Lambda_3^2)}\right]^{12/27} A_c\,, \qquad A_c = \left[\frac{\ln(m_b^2/\Lambda_4^2)}{\ln(m_c^2/\Lambda_4^2)}\right]^{12/25} A_b\,,$$

$$A_b = \left[\frac{\ln(m_t^2/\Lambda_5^2)}{\ln(m_b^2/\Lambda_5^2)}\right]^{12/23} A_t\,. \tag{3.15}$$

The Λ_i^2 in (3.14) and (3.15) are chosen such that $\alpha_s(Q^2)$ is continuous at each quark threshold, with

$$\alpha_s(Q^2) = \frac{12\pi}{b_n \ln(Q^2/\Lambda_n^2)} \tag{3.16}$$

for $m_n^2 \leqslant Q^2 \leqslant m_{n+1}^2$:

$$\Lambda_{n+1}^2 = (m_{n+1}^2)^{(b_{n+1}-b_n)/b_{n+1}} \, (\Lambda_n^2)^{b_n/b_{n+1}}\,, \qquad b_n = 33 - 2n\,. \tag{3.17}$$

Taking $\mu = (0.5 - 1)$ GeV and $\Lambda_3^2 = 0.12\ \mathrm{GeV}^2$ to be consistent with scaling violations in deep inelastic electroproduction, we find

$$A_s \approx 2.1 - 3.4\,, \qquad A_c \approx 1.7\,, \qquad A_b \approx 1.4\,, \qquad \text{for } 5\ \mathrm{GeV} \leqslant m_t \leqslant m_W\,. \tag{3.18}$$

So short-distance enhancement effects are smaller [17] for bottom quarks than for strange or charmed quarks. Also, we believe [17] that long distance operator matrix element enhancements are much smaller for heavier quarks. Indeed, the nonleptonic decays of charmed mesons do not seem [1,32] to be greatly enhanced relative to their semileptonic decays. We therefore neglect henceforth nonleptonic enhancement factors for bottom decays.

Calculating figs. 2a and 2b in a free quark model, we have

$$\Gamma_{(a,b)} \approx \left\{\begin{matrix} s_1^2 s_3^2 \\ \\ (s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta)\, F\left(\dfrac{m_c^2}{m_b^2}\right) \end{matrix}\right\} \times 3 \times \frac{G_F^2 m_b^5}{192\pi^3} \tag{3.19}$$

where $F(x) = 1 - 8x + 8x^3 - x^4 - 12x^2 \ln x \approx 1$ for u, d, s quarks, and ≈ 0.3 for $q \equiv c$. We therefore have

$$\frac{\Gamma_{(b)}}{\Gamma_{(a)}} \approx \left(\frac{F(m_c^2/m_b^2)}{1}\right) \frac{s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta}{s_1^2 s_3^2} \approx 6\, \frac{s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta}{s_3^2}\,. \tag{3.20}$$

Since we expect on the basis of (3.6) and (3.10) that $s_2^2 + s_3^2 + 2s_2s_3 \cos\delta = O(s_3^2)$ we expect from (3.20) that $\Gamma_b >> \Gamma_a$.

Turning now to figs. 2c and 2d, we use a free quark model for $B^0 \equiv b\bar{d}$ decay (here and henceforth we identify $m_B \approx m_b$):

$$\Gamma_{(c,d)} \approx \left\{ \begin{matrix} s_1^2 s_3^2 m_u^2 \\ (s_2^2 + s_3^2 + 2s_2s_3 \cos\delta)\left(1 - \frac{m_c^2}{m_b^2}\right)^2 m_c^2 \end{matrix} \right\} \frac{1}{3}\, \frac{G_F^2}{8\pi} f_B^2 m_b\,, \tag{3.21}$$

where

$$\langle 0 | A_\mu^{b\bar{d}}(0) | B_0(k)\rangle \equiv i k_\mu f_B\,. \tag{3.22}$$

We therefore have

$$\frac{\Gamma_{(d)}}{\Gamma_{(c)}} \approx \frac{(1 - m_c^2/m_b^2)^2}{1}\, \frac{m_c^2}{m_u^2}\, \frac{s_2^2 + s_3^2 + 2s_2s_3\cos\delta}{s_1^2 s_3^2} \simeq 600\, \frac{s_2^2 + s_3^2 + 2s_2s_3\cos\delta}{s_3^2} \tag{3.23}$$

if we take $m_u \approx 300$ MeV. Again the charmed decay modes dominate over the non-charmed final states.

We now turn to the "penguin" diagrams of figs. 2e and 2f. In the free-field approximation the penguin diagrams do not contribute because they reduce to a non-diagonal mass renormalization. Furthermore, if $m_W^2 >> m_q^2$ for all quarks, they do not contribute in the leading log approximation for strong interaction corrections because of the generalized GIM mechanism. However, we believe [17] that they play an important role in the matrix element enhancement for strange particle decay, where soft gluon exchange should be understood in the generic penguin diagram of fig. 2e. For charm decay, there is no contribution to the dominant $\Delta S = \Delta Q$ transitions because the relevant operator is exotic in flavour, but *a priori* there may be a contribution to bottom decay; however, we expect a suppression of order $\alpha(m_b^2)/\alpha(\mu^2)$ relative to strange particle decay. The lowest order contributions are those of fig. 2f.

Shifman et al. [33] have developed a quantitative procedure for evaluating the contribution of the effective four-quark operator of fig. 2f which gives a reasonable description of strange particle decays. They separate the momentum integration into the region $\mu^2 \leqslant Q^2 \leqslant m_c^2$, n = number of flavours = 3, where the GIM mechanism is ineffective and $m_c^2 \leqslant Q^2 \leqslant m_W^2$, $n = 4$, where cancellation occurs. Then the penguin contribution is governed by $\ln(m_c^2/\mu^2)$. Adapting their analysis to b-decay, the relevant regions are $m_b^2 \leqslant Q^2 \leqslant m_t^2$, $n = 5$ and $m_t^2 \leqslant Q^2 \leqslant m_W^2$, $n = 6$. Since GIM is fully effective in the second region, there is no contribution for $m_b^2 = m_t^2$ and we obtain an upper bound by letting $m_t^2 \to m_W^2$. Then we find an effective coupling which is at most 13% of the usual Fermi coupling, and generally negligible. However, for "annihilation" processes of the type in fig. 2c this contribution can be relatively more important because there is no helicity suppression; the gluon in fig. 2f couples to both left- and right-handed quarks. For B^0 decay via annihilation, we find

$$\frac{\Gamma(B^0 \to \bar{d}s)}{\Gamma(B^0 \to c\bar{u})} = \frac{\Gamma(B^0 \to \bar{d}d)}{\Gamma(B^0 \to \bar{u}u)} \lesssim 1 . \tag{3.34}$$

(In the case of B^- decay via two quarks there is an additional suppression of $\frac{1}{9}$ because the colour counting is more favourable for the graphs analogous to figs. 2c and 2d.)

According to the general analysis of R.K. Ellis [34], the only independent operators of lowest dimension (5 and 6) are the four fermion operators. However, the momentum cut-off in the integral which determines the effective operator (dimension 7) for $\bar{q}q' \to 2$ (colour symmetric) gluons is much lower than the W mass, so dimensional arguments may be irrelevant. Transcribing the calculation [30] of $\bar{s}d \to 2\gamma$ to the present case we find

$$\frac{\Gamma(b \to s + GG)}{\Gamma(b \to c + \bar{u}d)} = \frac{\Gamma(b \to d + GG)}{\Gamma(b \to u + \bar{u}d)} \lesssim \frac{3\alpha_s^2(m_b^2)}{8\pi^2} \approx 3 \times 10^{-3} , \tag{3.35}$$

$$\frac{\Gamma(B^0 \to GG)}{\Gamma(B^0 \to c\bar{u})} \lesssim \frac{m_b^2}{m_c^2} \frac{6\alpha_s^2(m_b^2)}{\pi^2} \frac{s_1^2 s_3^2}{s_2^2 + s_3^2 + 2s_3 s_2 \cos\delta}$$

$$\approx 0.02 \frac{s_3^2}{s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta} . \tag{3.36}$$

The gluon decay channels therefore seem unimportant.

It remains to estimate the relative strength of the "annihilation" processes with those of fig. 2a. For this we need an estimate of f_B; we have borrowed three different approaches from the literature, all of which give the estimate

$$|f_B| \approx 500 \text{ MeV} . \tag{3.37}$$

(a) Gerstein and Khlopov [35] assume SU(4) symmetry for the P-$\bar{q}_1 q_2$ couplings where P($\bar{q}_1 q_2$) is a pseudoscalar bound state and calculate the diagram of fig. 3 with a cut-off adjusted to fit the pion decay constant. Taking $m_u = m_d = 350$ MeV, $m_b = 5$ GeV, their result gives

$$|f_B| = 540 \text{ MeV} .$$

(b) Using a sum rule based on dispersion relations and asymptotic freedom, applicable for heavy quark bound states, Novikov et al. [36] obtain the inequality

$$|f_P| \lesssim m_P/\sqrt{32\pi} = 500 \text{ MeV}$$

for $m_P = 5$ GeV.

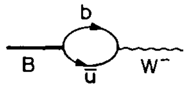

Fig. 3. Diagram for estimating [35] f_B.

(c) Using a non-relativistic quark model to estimate the B wave function at the origin

$$|\psi(0)|^2 \approx 0.1\ \mathrm{GeV}^3$$

Cahn and S.D. Ellis [11] obtain

$$f_B = 2\sqrt{\frac{3}{m_B}}\ |\psi(0)| \approx 500\ \mathrm{MeV}\ .$$

Thus we conclude that (3.37) is a safe estimate and we shall use it in the following.
Then from (3.19) and (3.21) we obtain, for example,

$$\frac{\Gamma_{(d)}}{\Gamma_{(b)}} = \frac{8\pi^2}{3}\frac{(1-m_c^2/m_b^2)^2}{F(m_c^2/m_b^2)} f_B^2 \frac{m_c^2}{m_b^4} \approx \tfrac{1}{10}\ . \tag{3.38}$$

The analogous ratio for B^- decay is O(1). Since the dominant penguin graph gives a contribution at most equal to Γ_d [eq. (3.34)], we conclude that, barring a fortituous cancellation between Cabibbo angles, charmed final states will dominate. For B^- decay one half of the hadronic decays may have a simple two-jet structure [11], but these should be much less in B^0 decays. Further we expect

$$\Gamma(B \to \ell\nu + \text{hadrons}) = \tfrac{1}{3}\,\Gamma_{(b)}\ , \qquad \ell = \mu, e\ ,$$

where again charmed final states should dominate, and

$$\frac{\Gamma(B^- \to \tau^-\nu_\tau)}{\Gamma_{(b)}} = \frac{8\pi^2 f_B^2 m_\tau^2 (1-m_c^2/m_b^2)^2 s_1^2 s_3^2}{F(m_c^2/m_b^2)\, m_b^4 (s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta)} \approx (1\text{–}2)\ \%\ , \tag{3.39}$$

whereas the semileptonic τ^- channel is suppressed by phase space and/or Cabibbo angles. Adding up the principal decay modes, we find

$$\Gamma_B \approx 2\Gamma_{(b)} \approx 2\,\frac{G_F^2 m_b^5}{192\pi^3}\,(s_2^2 + s_3^2 + 2s_2 s_3 \cos\delta)\ , \tag{3.40}$$

$$\tau_B \approx 0.4 \times 10^{-14}(s_2^2 + s_3^2 + 2s_2 s_3\cos\delta)^{-1} \gtrsim 10^{-13}\ , \tag{3.41}$$

where we have assumed that the relevant combination of mixing angles is unlikely to be greater than the Cabibbo angle. The semileptonic branching ratios should be close [17] to their free quark value of 20%.

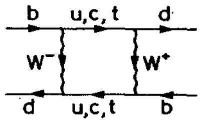

Fig. 4. Diagram which determines the $B^0 - \overline{B}^0$ mass matrix.

We expect strange final states to dominate T-decays even more strongly than charm for B-decays, because there is no phase-space suppression:

$$\Gamma_T \simeq 6\Gamma_{(b)} \gtrsim 3 \times 10^{-13} , \qquad \text{for } m_b > m_t , \tag{3.42}$$

$$\frac{\Gamma_T(\Delta S = 0)}{\Gamma_T(\Delta S = -1)} \approx s_1^2 \approx \tfrac{1}{20} \tag{3.43}$$

and two-jet final states should be negligible. Thus T-decays are essentially indistinguishable from charm decays except for the energy release.

3.3. Mass matrices and CP violation in the neutral bottom and top meson systems

We now consider the amount of mixing and *CP* violation in the $B^0 - \overline{B}^0$ and $T^0 - \overline{T}^0$ meson systems. As in sect. 3.2, we consider only $B^0 - \overline{B}^0$: results of the $T^0 - \overline{T}^0$ system can be obtained by the substitutions (3.13). Straightforward calculations of the diagrams in fig. 4, along the lines of refs. [30] and [14] give

$$\langle \overline{B}^0 | -\mathcal{L}_{\text{eff}} | B_0 \rangle \approx \frac{G_F}{\sqrt{2}} f_B^2 m_B^2 \frac{\alpha}{4\pi} \left[\frac{m_t}{38\ \text{GeV}} \right]^2 s_1^2 s_2^2 c_2^2 c_3^2 \cos 2\delta , \tag{3.44}$$

$$\left| \frac{\operatorname{Im} M_{12}^B}{\Delta m^B} \right| \approx \tan 2\delta , \tag{3.45}$$

if $m_t^2 >> m_c^2$, s_1^2, s_2^2, $s_3^2 << 1$ as we expect. The estimate (3.44) entails inserting the vacuum state, as in ref. [30]; this is equivalent to assuming a valence quark wave function for the B^0. We do not know any better, and hope the order of magnitude is correct. In any case, this uncertainty cancels out in the *CP* violating ratio (3.45). Comparing with (3.11) and recalling the bounds (3.6) and (3.10), we conclude that *CP* violation may be much stronger in the $B^0 - \overline{B}^0$ mass matrix than in the $K^0 - \overline{K}^0$ mass matrix.

To compare (3.44) with the magnitude of $\Gamma(B^0)$ in eq. (3.40), we use $f_B \sim 0.5$ GeV as before and find

$$\frac{\Delta m_{B^0}}{\Gamma_{B^0}} = \frac{8}{\sqrt{2}} \frac{\pi^2 \alpha f_B^2}{F(m_c^2/m_b^2)\, G_F m_b^4} \left(\frac{m_t}{38\ \text{GeV}} \right)^2 s_1^2 \left(\frac{s_3^2}{s_2^2 + s_3^2 + 2 s_2 s_3 \cos\delta} \right)$$

$$\approx m_t^2 / 700\ \text{GeV}^2 , \tag{3.46}$$

if we assume that the angular factor in parenthesis is order unity. The difference in lifetimes for the two B^0 states is determined by the relative importance of decay final states common to B^0 and $\overline{B}^0$, i.e. those with no net leptonic or flavour quantum numbers. These come essentially from fig. 2a, so from (3.20) and (3.40) we obtain

$$\frac{\Delta\Gamma}{2\Gamma} \approx \frac{1}{6}\,\frac{\Gamma_{(b)}}{\Gamma_B} \approx \tfrac{1}{12}\,. \tag{3.47}$$

The number of "wrong sign" dimuons from an initial $\overset{(-)}{B}{}^0$ state is given by [37]

$$\overset{(-)}{r} = \frac{\Gamma(\text{“}\overset{(-)}{B}{}^0\text{”} \to f\ell^+ \overset{(-)}{\bar{\nu}_\ell})}{\Gamma(\text{“}\overset{(-)}{B}{}^0\text{”} \to f\ell^{\overset{(+)}{-}} \overset{(-)}{\nu_\ell})}\,, \qquad \sqrt{r\bar{r}} = \frac{(\Delta m/\Gamma)_B^2 + (\Delta\Gamma/2\Gamma)_B^2}{2 + (\Delta m/\Gamma)_B^2 - (\Delta\Gamma/2\Gamma)_B^2}\,. \tag{3.48}$$

For $m_t \lesssim 8$ GeV, the contribution (3.47) is dominant and

$$\sqrt{r\bar{r}} \approx (\tfrac{1}{2} - 1)\,\%\,.$$

For $m_t > 8$ GeV,

$$\sqrt{r\bar{r}} \simeq \frac{1}{2}\left(\frac{m_t^2}{700\ \text{GeV}^2}\right)^2$$

becomes a measure of the top quark mass. Because the dominant $|\Delta S| \neq 0$ channel for top decay is not suppressed by phase space, the quantities analogous to (3.46) and (3.47) are suppressed by about $\frac{1}{3}$, and mixing effects will be smaller.

For *CP* violation the relevant parameter is

$$\epsilon_B \equiv \frac{\frac{1}{2}\,\mathrm{Im}\,\Gamma_{12}^B + i\,\mathrm{Im}\,M_{12}^B}{\frac{i}{2}\,\Delta\Gamma_B - \Delta m_B}\,, \qquad \sqrt{\frac{r}{\bar{r}}} = \left|\frac{1-\epsilon_B}{1+\epsilon_B}\right|^2\,, \tag{3.49}$$

where $\mathrm{Im}\,\Gamma_{12}$ depends on the *CP* violation in the decay modes common to B^0 and $\overline{B}^0$. Since these arise mainly from the graph of fig. 2a, which is *CP* conserving, we expect $\mathrm{Im}\,\Gamma_{12}$ to be negligible as in the kaon system. However, since $\tan\delta$ in eq. (3.45) is *a priori* arbitrary, and for $m_t \gtrsim 5$ GeV,

$$2|\Delta m/\Delta\Gamma| \gtrsim 0.4\,,$$

the value of ϵ can be quite large for the neutral B systems.

3.4. Neutrinoproduction of top and bottom particles

It is clear that the ideal way to study the properties of naked top and bottom particles will be in e^+e^- annihilation at PETRA, CESR or PEP. However, these experiments are for the future, and the best way to look for manifestations of top or bottom particles right now may be in neutrino experiments. It is trivial to calculate the characteristics of their production by ν and $\bar{\nu}$ using the standard quark-parton model: the results are in table 1, on which we make a few explanatory comments.

The symbols V, S, and C refer to the fraction of the momentum of the *weakly interacting* constituents of the nucleon (i.e. excluding gluons) which is carried by valence, non-charmed sea and charmed sea quarks, respectively. We estimated

$$C \approx 0.01 , \qquad S \approx 0.05 , \qquad V \approx 1 , \tag{3.50}$$

and have completely neglected interactions with the $b\bar{b}$ or $t\bar{t}$ parts of the sea. The upper limits (3.6) and (3.9) on s_3^2 and s_2^2 give stringent upper limits on the production of b and t quarks which are listed in column 3 of table 1. Unfortunately, lower limits on t or b production are derisory: using the bounds of sect. 3.1 we find a lower limit of 10^{-5} for b production by $\bar{\nu}$, and a similar limit for t production by ν. Since left-handed b quarks like to decay into c quarks as discussed in sect. 3.2, b quark production by $\bar{\nu}$ can yield *charmed* mesons in the final state, whereas the GIM model would only yield *anticharmed* mesons in $\bar{\nu}$-nucleon scattering. Therefore emulsion exposures to $\bar{\nu}$ could, in principle, yield events with two sequential finite length tracks, corresponding to the bottom and charmed particles propagating each with a lifetime $O(10^{-13})$ sec. The semileptonic decays of bottom and charmed particles would also yield distinctive multimeson final states, as categorized in column 4 of table 1. The upper limits on rates are calculated assuming semileptonic branching ratios O(15%), as observed for charmed mesons [1,32] and expected [17] for bottom and top particles because of the relative non-enhancement of non-leptonic decays. Present neutrino experiments [38] are now getting down to the production fractions of multilepton events which column 4 of table 1 suggests are interesting. Finding novel multimuon events at a much higher rate in $\bar{\nu}$-nucleon scattering than in ν-nucleon scattering would be very consistent with the assignment of the Υ as bottomonium, with the associated bottom particle masses starting at ~5 GeV. The threshold suppression factor [10] for heavy quark production probably means that the bounds in column 4 can only be closely approached for $E_{\nu,\bar{\nu}} >$ 100 GeV. Notice that top production is much less distinctive, as it generally mimics charm production, but is much less copious. The only exception is if there is substantial $T^0 - \bar{T}^0$ mixing, in which case like sign dimuons might occur. However, there are no substantial tri- or quadrilepton signals for t-production, unlike the b-production case.

4. Conclusions and discussion

Our main conclusions are the following: for the hadronic production of hidden bottom or top states we find:

Substantial (~30%) production of $\mu^+\mu^-$ *via* $J'_B \equiv 2^3S_1$ as well as $J_B = 1^3S_1$, with a possible contribution of $J''_B \equiv 3^3S_1$ also. This results from assuming similar production mechanisms for J_B, J'_B and J''_B, and calculations of low branching ratios for $J'_B \to J_B\pi\pi$ and other cascade decays.

For the decay properties of top and bottom particles in the six left-handed quark model [6] we find:

Charmed particles should appear in most decays of bottom particles, if the latter are lighter than tops.

Lifetimes $\gtrsim 10^{-13}$ sec for bottom or top particles with masses O(5) GeV.

The possibility of substantial B^0 ($\equiv b\bar{d}$) – $\bar{B}^0$ ($\equiv \bar{b}d$) meson mixing if $m_t > m_b$.

CP violating effects in the $B^0 - \bar{B}^0$ and $T^0 - \bar{T}^0$ systems which are considerably larger than in the $K^0 - \bar{K}^0$ system.

Distinctive multilepton signatures for b production by $\bar{\nu}$, but fewer corresponding signatures for t production if $m_t < m_b$. The absolute rates are not exactly predictable, but must be low.

What are the prospects for determining soon whether $\Upsilon(9.5)$ is bottomonium, or toponium, or a superposition of the two? The line structure of $\Upsilon(9.5)$ appears complicated with $\geqslant 2$ states. This meshes better with the bottomonium or superposition hypotheses. The low cross section and signal to background ratio may favour the bottomonium hypothesis. The situation could perhaps be resolved if one saw a higher mass state, so that one could compare its characteristics with those of $\Upsilon(9.5)$. The observation of anomalous multimuon events at a higher rate in $\bar{\nu}$ scattering than in ν scattering would favour the b-quark hypothesis. However, we see no clear chance of resolving the ambiguities before experiments with the next generation – PETRA, CESR, PEP – of e^+e^- colliding ring machines.

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Note added in proof

Since the publication of our paper more data have been publised [39] on the 9.5 GeV region in $pp \to \mu^+\mu^- + X$. The separation into at least two peaks $\Upsilon(9.4)$ and $\Upsilon'(10.0)$ has been established, and the possibility exists of a third peak $\Upsilon''(10.4)$. The mass separations are larger than anticipated by Eichten and Gottfried [19], but we do not believe they affect our qualitative conclusions about the ratios of J_B, J'_B and J''_B (J_T, J'_T and J''_T) production. Most critical is the decay rate for $V' \to V\pi\pi$. Using eq. (2.7) and $m' - m = 600$ MeV, we find

$$\frac{\Gamma(V' \to V\pi\pi)}{\Gamma(\psi' \to J/\psi\pi\pi)} \simeq \frac{\tilde{g}^2_{V'V}}{\tilde{g}^2_{\psi' J/\psi}} \times 0.15.$$

In order for our predictions of $\sigma B|_V : \sigma B|_{V'}$ to remain valid, we need $\tilde{g}^2_{V'V}/\tilde{g}^2_{\psi' J/\psi}$ somewhat <1. In view of the tendency of the Zweig rule for suppressing disconnected quark diagrams to improve when heavier quarks are involved, we find it reasonable to believe that $\Gamma(V' \to V\pi\pi) \lesssim 5$ keV. In this case the estimates (1.2) and (1.3) would be essentially unaffected, except possibly for the relative amounts of J''_B and J''_T. For an alternative discussion of this point, see the recent review by Gottfried [40].

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