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Neutral Higgs boson from decays of heavy-flavored mesons

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It is shown that the decay of heavy-flavored mesons can be a source of neutral Higgs bosons.

I. INTRODUCTION

The $SU(2)\times U(1)$ non-Abelian gauge theory of weak and electromagnetic interactions, with mass generation through a single Higgs doublet, and three generations of quarks and leptons incorporating the Kobayashi-Maskawa (KM) generalization of the Cabibbo and Glashow-Iliopoulos-Maiani (GIM) formulations of universality, along with the constitutent-quark description of observed (color-singlet) hadrons, will be referred to as the standard model. An important question is the existence or nonexistence of a physical neutral Higgs boson H with the properties of the elementary Higgs scalar of the standard model.[1] Because of its role in mass generation, the couplings of the H to the other particles in the standard model are known, but its mass is an arbitrary parameter.[2] The empirical limit is only that no H exists with mass less than about 18 MeV.[1,3] The reason that data from currently accessible reactions provide so little constraint on the existence of the H is that the H couples very weakly to light particles. To produce the H generally requires reactions involving heavy particles, the gauge vector bosons Z^0, $W^\pm$, or heavy quarks Q. In particular, one looks[1] to reactions involving the production of Z^0 at LEP, or decays of heavy quarkonium ($Q\bar{Q}$). In this paper we discuss another mechanism for production of H, the weak decays of heavy-flavored mesons ($Q\bar{q}$), in particular, the inclusive decay $P\rightarrow H +$ hadrons, where P is the heavy pseudoscalar meson ($Q\bar{q}$). In the standard model the couplings of the H are flavor diagonal, so the process is higher order. The standard model is renormalizable so this higher-order process can be calculated, and the H can couple to heavy (virtual) particles, so the result may not be terribly small. We find that in some cases this branching ratio is greater than the branching ratio for the corresponding quarkonium decay $V(Q\bar{Q})\rightarrow H+\gamma$.

II. CALCULATIONS

For a heavy-flavored meson composed of one very massive quark carrying the flavor and one light antiquark, it is argued[4] that the inclusive *semileptonic* decay rate is just the semileptonic decay rate of the massive quark (spectator model).[5] Note that for the semileptonic process one is effectively replacing $P\rightarrow WX$ by $Q\rightarrow Wq''$ (Fig. 1). In the standard model the H is just as elementary as the W (they each appear as a field in the Lagrangian), so the same argument can be used that the inclusive rate for $P\rightarrow HX$ is given by the rate for $Q\rightarrow Hq'$.

The Feynman diagrams which contribute to $Q\rightarrow Hq'$, through order g^3, are shown in Fig. 2. Individual diagrams contain ultraviolet divergences. Divergences whose coefficients are independent of the intermediate quark (Q_i) mass are eliminated by the generalized GIM mechanism

$$\sum_i C_{iB}C^*_{iA}=\delta_{BA}=0 \quad \text{for} \quad A\neq B\ , \tag{2.1}$$

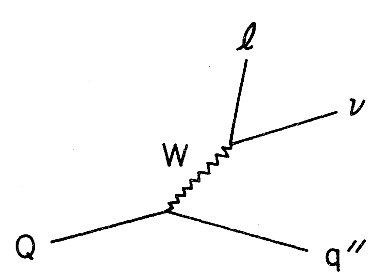

FIG. 1. Semileptonic decay of a heavy quark.

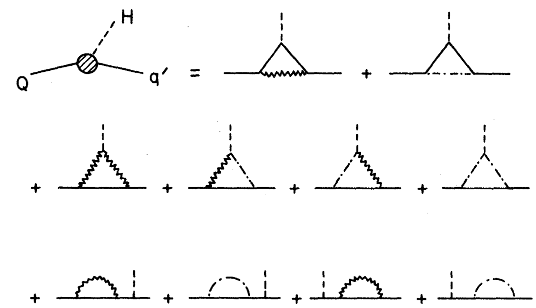

FIG. 2. Feynman diagrams for $Q \to Hq'$. Solid line, quark (Q,Q_i,q'); wiggly line, gauge vector boson ($W^{\pm}$); dashed line, physical Higgs boson (H); dash-dot line, unphysical Higgs boson ($\phi^{\pm}$).

where the C_{iB} are elements of the KM matrix. The remaining divergences cancel when all of the diagrams are combined. For the process $b \to H+s$ with intermediate t quark, the matrix element obtained from the sum of the ten diagrams in Fig. 2 is

$$M(b \to H_s) = \frac{3g^3}{256\pi^2} C_{tb} C_{ts}^* \frac{m_b m_t^2}{m_W^3} \times \bar{s}(p')(1+\gamma_5) b(p) , \qquad (2.2)$$

where $g^2/8m_W^2 = G_F/\sqrt{2}$. In this calculation we have made the approximation

$$m_W^2, m_t^2 >> m_b^2 >> m_s^2 . \qquad (2.3)$$

No approximation has been made on the ratio $x_t = m_t^2/m_W^2$. Individual diagrams in Fig. 2 contribute complicated functions involving $1/(1-x)$ and $(\ln x)/(1-x)$, but when all the diagrams are added they combine to the simple result (2.2) for arbitrary m_t^2/m_W^2. The factor m_t^2 comes from the operation of the generalized GIM mechanism (2.1), which suppresses flavor-changing neutral currents by the factor m_i^2/m_W^2 for $m_W \to \infty$. The factor m_b appears because the H couples proportional to mass. [Following (2.3) we have dropped terms with $m_s m_t^2/m_W^3$.] For the process $t \to H+c$ (or u), the calculation is more complicated because the heaviest quark is now external rather than internal, and there is also the possibility of a large $m_H \sim m_t$. In this case the result is simple only in the limit of large m_W^2.

$$M(t \to Hc) = \frac{3g^3}{256\pi^2} C_{tb}^* C_{cb} \frac{m_t m_b^2}{m_W^3} \times \bar{c}(p')(1+\gamma_s) t(p) \qquad (2.4)$$

for

$$m_W^2 >> m_t^2 >> m_b^2 >> m_c^2 . \qquad (2.5)$$

Again, the generalized GIM relation has been important; in addition to its role in eliminating ultraviolet divergences it has also eliminated terms m_t^3/m_W^3 which would otherwise appear in (2.4). Comparing (2.2) and (2.4) we see that the enhancement (diminishment of GIM suppression) from large M_t will be more effective in b decays than in t decays.

The rate computed from the square of (2.2) or (2.4), includes an integration over phase space. In the spirit of the heavy quark plus light spectator model, we neglect the masses of the final-state particles, relative to the initial heavy quark. To partly compensate for the approximation thus introduced, we compute the ratio of the inclusive rate for $P \to HX$ to the inclusive semileptonic rate for $P \to e\nu X$, also computed in the spectator model (Fig. 1) with neglect of final-state particle masses in the phase-space integral. The result is

$$\frac{\Gamma(P \to HX)}{\Gamma(P \to e\nu X)} = \frac{|C_{Q_i Q} C_{Q_i q'}^*|^2}{|C_{Qq''}|^2} \frac{27\sqrt{2}}{64\pi^2} \left[\frac{m_{Q_i}}{m_Q} \right]^4 \times \phi(\text{masses}) . \qquad (2.6)$$

The factor ϕ(masses) is a ratio of phase-space factors, normalized to be one in the limit in which all the final-state particle masses are neglected relative to m_Q. Thus $\phi \approx 1$ should be a good approximation except when m_H approaches m_Q. When m_H approaches m_Q, one may include a suppression factor

$$\phi \approx 1 - \frac{m_H^2}{m_Q^2} \qquad (2.7)$$

obtained by keeping m_H in the phase-space integral for $Q \to Hq'$. We note that in the kinematic region in which this suppression factor is important the impulse approximation underlying the spectator model is not justified and the corrections are not easily calculable.

In addition to the set of "vertex" diagrams of Fig. 2, there is a set of "bremsstrahlung" diagrams shown in Fig. 3, which are also of order g^3 and involve an H in the decay of Q. Since the final state in the diagrams of Fig. 3 is different than the final state in the diagrams of Fig. 2, it is appropriate to compute the rates separately and compare (or add) them. Calculation of the bremsstrahlung rate requires integration of the square of the matrix element over four-particle phase space, a tedious undertaking. However, a bound may be obtained simply by replacing the momentum dependence of the matrix element by its maximum value (in the kinematically allowed phase space) and using the known[6] simple phase-space integral for all massless final-state particles. Each of these approximations is an overestimate, and the result is

$$\frac{\Gamma_{\text{brem}}}{\Gamma_{\text{vertex}}} \lesssim \frac{|C_{Qq''}|^2}{|C_{Q_iQ}C^*_{Q_iq'}|^2}\left[\frac{m_Q}{m_{Q_i}}\right]^4 . \tag{2.8}$$

Large powers of 2 and π have canceled out in the ratio of the product of the loop integral times the two-particle phase-space integral and the four-particle phase-space integral. For $Q=b$ and $Q_i=t$, (2.8) is entirely negligible, so we will continue with consideration of $B\to HX$ from (2.6).

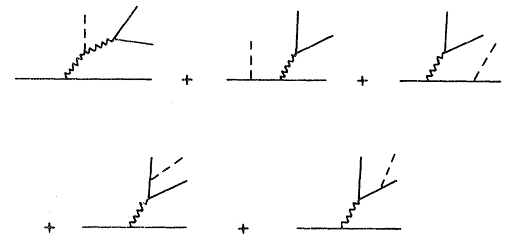

FIG. 3. Feynman diagrams for bremsstrahlung of H. The notation is the same as in Fig. 2. Since these are all diagrams with no loops we work in the U gauge in which there are no unphysical Higgs bosons.

The ratio of KM matrix elements appearing in (2.6) and in (2.8) for B decays is (we use the Particle Data Group[7] convention for the KM angles)

$$\frac{|C_{tb}C^*_{ts}|^2}{|C_{bc}|^2}=\frac{|(c_1s_2s_3+c_2c_3e^{i\delta})(c_1s_2c_3-c_2s_3e^{-i\delta})|^2}{|c_1c_2s_3-s_2c_3e^{i\delta}|^2} . \tag{2.9}$$

It is known that

$$c_1\approx 1,\quad c_\delta\approx\pm 1$$

so that (2.9) reduces to

$$\frac{|C_{tb}C^*_{ts}|^2}{|C_{bc}|^2}\approx(s_2s_3\pm c_2c_3)^2 . \tag{2.10}$$

This number is also probably not very different from one.[8] The constituent b-quark mass is fairly well determined,[9] $m_b\simeq 4.9$ GeV. Then (2.6) gives

$$\frac{\Gamma(B\to HX)}{\Gamma(B\to e\nu X)}\approx 1.5\times 10^{-5}\left[\frac{m_t}{m_b}\right]^4\left[1-\frac{m_H{}^2}{m_b{}^2}\right] \tag{2.11}$$

for $m_b=4.9$ GeV and $m_H\lesssim 4$ to 4.5 GeV. The semileptonic branching ratios for $B\to e\nu X$ and $B\to\mu\nu X$ have been measured[10] with an average value of about 11%. Thus

$$R=\frac{\Gamma(B\to HX)}{\Gamma(B\to \text{all})}\approx 1.5\times 10^{-6}\left[\frac{m_t}{m_b}\right]^4\left[1-\frac{m_H{}^2}{m_b{}^2}\right] . \tag{2.12}$$

Since the experimental[11] lower bound on m_t is close to 20 GeV, this is $\gtrsim 4\times 10^{-4}$, and rapidly increasing for $m_t>20$ GeV.

We compare this to the branching ratio for $\Upsilon\to H+\gamma$ (Ref. 12):

$$\frac{\Gamma(\Upsilon\to H\gamma)}{\Gamma(\Upsilon\to e\bar{e})}\approx\frac{G_Fm_b{}^2}{\sqrt{2}\pi\alpha}\left[1-\frac{m_H{}^2}{m_\Upsilon{}^2}\right]\approx 10^{-2} . \tag{2.13}$$

The branching ratio $\Upsilon\to e^-e^+$ is measured[13] to be about 0.03, so

$$\frac{\Gamma(\Upsilon\to H\gamma)}{\Gamma(\Upsilon\to \text{all})}\approx 3\times 10^{-4} , \tag{2.14}$$

which is less than (2.12) (very much less if m_t is significantly greater than 20 GeV). Of course m_H up to 9 GeV is accessible to the Υ decay, while production of H in B decay is limited to $m_H\lesssim 4.5$ GeV. But given the availability of the $\Upsilon(4s)$ as a "B factory," B decay (2.12) may be the best choice to look for an H of mass less than or equal to about 4 GeV.

For T decay, the factor $(m_t/m_b)^4$ in (2.11) is replaced by $(m_b/m_t)^2$, and the ratio of KM matrix elements is now much less than one because the $T\to e\nu X$ decay is Cabibbo-KM favored. Thus, for

T decay, R (2.12) is uninterestingly small, less than 10^{-8}. The result (2.8) suggests that for $T \to HX$ we should consider the bremsstrahlung contribution. Again, we get a generous upper bound by taking the maximum value of the matrix element and the ultrarelativistic limit of the phase-space integral. This gives

$$\frac{\Gamma_{\text{brem}}(T \to HX)}{\Gamma(T \to e\nu X)} \lesssim \frac{1}{16\pi^2} G_F m_t^2 \qquad (2.15)$$

which is less than 10^{-4} for any $m_t \lesssim 40$ GeV. Considering that this is probably a substantial overestimate and the present experimental inaccessibility of T mesons, we do not pursue this any further.

III. CONCLUSION AND DISCUSSION

When the next generation of e^-e^+ colliders are available one expects to find the Z^0 and perhaps $(t\bar{t})$, and a Higgs boson of mass less than 40 or 50 GeV can probably be found in the decays of these heavy objects.[1] But these machines are several years away and at the present, the heaviest systems we have available are the Υ resonances $(b\bar{b})$ and the B mesons $(b\bar{q})$, copiously produced at the $\Upsilon(4S)$ resonance. We have calculated the branching ratio for the inclusive decay $B \to HX$ and shown that for $m_H \lesssim 4$ GeV that ratio is greater than the branching ratio for the $\Upsilon \to H\gamma$ decay of the Υ. We will discuss briefly the question of the experimental signal for the decay $B \to HX$ and show that data already available put some constraint on the allowed combination of m_H and m_t.

The X in $H+X$ consists of ordinary hadrons, typically, one kaon and some number of pions (we assume that $|C_{td}|^2 << |C_{ts}|^2$ so that the favored quark decay is $b \to H + s$). The decay modes of the H depend on its mass; they are discussed in the papers cited in Ref. 1 and are mostly based on the formula

$$\Gamma(H \to f\bar{f}) = \frac{g^2}{32\pi} m_H \left[\frac{m_f}{m_W}\right]^2 \left[1 - 4\frac{m_F^2}{m_H^2}\right]^{3/2} \times \begin{cases} 1 & \text{for } l\bar{l} \\ 3 & \text{for } q\bar{q} \end{cases} . \qquad (3.1)$$

(i) For $m_H \lesssim 1$ GeV, H will decay into $\mu\bar{\mu}$, $\pi\pi$. For light hadrons (quarks) it is hard to calculate the Higgs-boson decay branching ratios because it is not clear what to use for the light quark masses, but a fair guess[1] may be a 10% branching ratio for $\mu\bar{\mu}$. So the signal is $\mu\bar{\mu}$ pairs, and no $e\bar{e}$ pairs, with the invariant mass squared of the $\mu\bar{\mu}$ pair equal to m_H^2. For this signal, one has to multiply R [Eq. (2.12)] by the guessed 10% branching ratio for $H \to \mu\bar{\mu}$.

(ii) For 1 GeV $\lesssim mH \lesssim 3.5$ GeV, H will decay into $\mu\bar{\mu}$, $s\bar{s}(K\bar{K} + \pi\text{'s})$. If we take a constituent s-quark mass in the range 0.4 to 0.5 GeV, (3.1) gives about a 2% branching ratio for $\mu\bar{\mu}$; with smaller current-algebra s-quark masses one obtains a rather larger branching ratio for $\mu\bar{\mu}$. So for the $\mu\bar{\mu}$ signal in this range of m_H, one may have to multiply R by 0.02. But if the branching ratio for $\mu\bar{\mu}$ is this small, then almost all of the $B \to HX$ decays will contain three K's. The presence of three K's in a B decay is not in itself a signal for the presence of an H, since conventional nonleptonic decays of the B will include three-K events [$b \to cs\bar{c}$ with c, $\bar{c}$ each including a $K(\bar{K})$ in its decay products]. But for $m_H \lesssim 3.5$ GeV, the limited phase-space available implies that the $H \to K\bar{K} + \pi$'s will include a reasonable fraction of two-body (no π's) events. So one can search the $B\bar{B}$ events with four (or six) K's for two K's whose invariant mass squared is m_H^2.

(iii) For $3.5 \lesssim m_H \lesssim 4$ to 4.5 GeV, H will decay into $\tau\bar{\tau}$, $c\bar{c}$. Now (3.3) gives 25 to 30% for the branching ratio to $\tau\bar{\tau}$.

There is already an experimental limit[14] on the branching ratio for the inclusive decay $B \to X\mu\bar{\mu}$,

$$\frac{\Gamma(B \to Xl\bar{l})}{\Gamma(B \to \text{all})} < 7.4 \times 10^{-3} \qquad (3.2)$$

for $l = e$ or μ, assumed equal. Before confronting this limit with (2.12), two model dependent assumptions which went into (3.2) should be considered. First, (3.2) was obtained by combining separate limits on $B \to Xe\bar{e}$ and $B \to X\mu\bar{\mu}$, assumed equal. Since an H heavy enough to decay into $\mu\bar{\mu}$ effectively never decays into $e\bar{e}$, we should compare (2.12) to the branching ratio limit for $B \to X\mu\bar{\mu}$ which is slightly larger than (3.2). Second, there was an experimental cut requiring each μ to have an energy greater than or equal to 1 GeV, so the number (3.2) includes a momentum acceptance factor of 1/0.17 which was calculated assuming the momentum spectrum from $b \to X\mu\bar{\mu}$. Our proposed mechanism is $b \to Hs$ followed by $H \to \mu\bar{\mu}$.

With this mechanism, for $m_H \gtrsim 2$ GeV, most of the μ's will have energy greater than 1 GeV (not all, because the H is recoiling against the s in the rest frame of the b, which is nearly the laboratory system). For illustrative purposes we will take an acceptance factor of 1/0.6, for $m_H \gtrsim 2$ GeV and $m_s{}^2 \ll m_H{}^2$, and consider

$$(1.5\times10^{-6})\left[\frac{m_t}{m_b}\right]^4\left[1-\frac{m_H{}^2}{m_b{}^2}\right]B(H\to\mu\bar{\mu})$$
$$\lesssim 2\times10^{-3} \quad \text{for } 2\lesssim m_H 3.5 \text{ GeV} . \qquad (3.3)$$

In this range of m_H we consider a conservative $B(H\to\mu\bar{\mu})\approx0.02$ and an optimistic $B(H\to\mu\bar{\mu})\approx 0.2$, discussed above. For $m_H \lesssim 2$ GeV, we have to take into account the decreasing fraction of μ's with (laboratory) energy greater than 1 GeV. In the range $3.5 \lesssim m_H \lesssim 4-4.5$ GeV, the H will decay into $\tau\bar{\tau}$ 25–30 % of the time. The probability for the subsequent decay of the τ and $\bar{\tau}$ to give two charged leptons (e or μ) is about $\frac{2}{5}\times\frac{2}{5}=0.16$. So in this range we replace the B in (3.3) by about 0.05. In Fig. 4 we show a rough sketch of the region of combined m_H, m_t values disallowed by these considerations. Because of the tradeoff between the $\mu\bar{\mu}$ and $K\bar{K}$ decay modes of the H, it is possible that the m_H, m_t limits illustrated here could be considerably improved by an experimental analysis which considered both these modes. Also, because the signal is $\mu\bar{\mu}$ pairs which reconstruct to a fixed invariant mass, it would seem possible to relax the momentum cut and increase the sensitivity to a lighter Higgs boson.

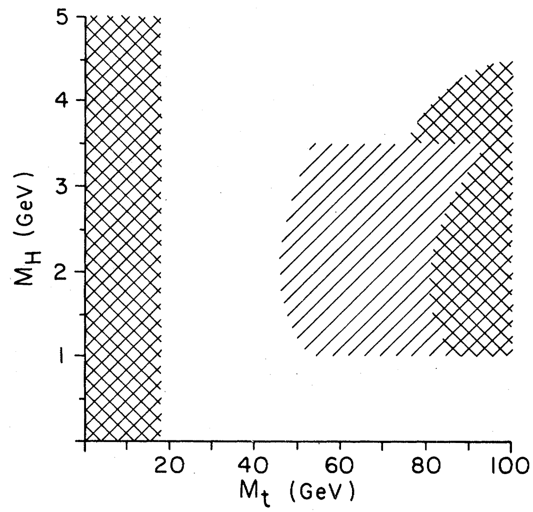

FIG. 4. Disallowed values of m_t, m_H. The single-hatched and double-hatched areas for $1.0 \lesssim m_H \lesssim 3.5$ GeV come from the assumptions $B(H\to\mu\bar{\mu})=0.2$ and $B(H\to\mu\bar{\mu})=0.02$, respectively. We emphasize that these results are illustrative, not optimal, because they are obtained by confronting (2.12) with an experiment (Ref. 14) which was not designed nor analyzed for this purpose.

[1]For reviews of the phenomenology of the Higgs boson see J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. B106, 292 (1976); M. K. Gaillard, Comments Nucl. Part. Phys. 8, 31 (1978); G. Barbiellini *et al.*, ECFA/LEP working group Report No. SSG/9/4 (unpublished).

[2]There is a theoretical argument based on the stability of the broken symmetry vacuum with respect to radiative corrections that m_H should not be less than about 7 GeV. However, the argument fails if there exist very heavy fermions or if one only requires that the tunneling time out of the false vacuum be greater than the current age of the universe. S. Weinberg, Phys. Rev. Lett. 36, 294 (1976); A. D. Linde, Pis'ma Zh. Eksp. Teor. Fiz. 23, 73 (1976) [JETP Lett. 23, 64 (1976)]; Phys. Lett. 70B, 306 (1977).

[3]R. Barbieri and T. E. O. Ericson, Phys. Lett. 57B, 270 (1975).

[4]M. K. Gaillard, B. W. Lee, and J. L. Rosner, Rev. Mod. Phys. 47, 277 (1975); J. L. Cortes, Phys. Rev. D 24, 2529 (1981).

[5]Most of the available data on D-meson decays is consistent, within sizable errors, with the spectator-model calculation of the inclusive semileptonic rates. (The situation with regard to the nonleptonic decays is in a state of flux.)

[6]E. Byckling and K. Kajante, *Particle Kinematics* (Wiley-Interscience, New York, 1973).

[7]Particle Data Group, Rev. Mod. Phys. 52, S1 (1980).

[8]A fairly popular prejudice is for all sines to be equal to small mass ratios and all cosines close to one. Some phenomenological investigators are V. Barger, W. F. Long, and S. Pakvasa, Phys. Rev. Lett. 42, 1585 (1979); R. E. Shrock, S. B. Treiman, and L.-L. Wang, *ibid.* 42, 1589 (1979); S. Pakvasa, S. F. Tuan, and J. J. Sakurai, Phys. Rev. D 23, 2799 (1981).

[9]J. L. Richardson, Phys. Lett. 82B, 272 (1979); R. S. Willey, Phys. Rev. D 26, 3254 (1982).

[10]C. Bebek *et al.* (CLEO), Phys. Rev. Lett. 46, 84 (1981); K. Chadwick *et al.* (CLEO), *ibid.* 46, 88 (1981); L. J. Spencer *et al.* (CUSB), *ibid.* 47, 771 (1981).

[11]D. P. Barber *et al.* (Mark J), Phys. Rev. Lett. 44, 1722 (1980).
[12]F. Wilczek, Phys. Rev. Lett. 39, 1304 (1977).
[13]B. Niczyporuk *et al.* (LENA), Phys. Rev. Lett. 46, 92 (1981); A. Brody *et al.* (CLEO), contribution to the 1981 International Symposium on Lepton and Photon Interactions at High Energies, Bonn (unpublished).
[14]Chadwick *et al.* (Ref. 10); J. Green *et al.* (CLEO), contribution to the 1981 International Symposium on Lepton and Photon Interactions at High Energies, Bonn (unpublished).