DO B MESON DECAYS EXCLUDE A LIGHT HIGGS?

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In the standard model with three generations a light Higgs boson will be produced in a quarter of all B meson decays. A novel calculation of the $b \to sH$ matrix element is given. Existing data on B meson decays excludes a light Higgs with mass less than 300 MeV or between 2 and 3.7 GeV. Because of theoretical uncertainties in the branching fractions for Higgs production and decay, a window between 300 MeV and 2 GeV is still allowed. Such a Higgs could be discovered at the $\Upsilon(4S)$ in inclusive $\pi^+\pi^-$ or K^+K^- invariant mass plots, and in events with six kaons.

The singular success of the standard model of electroweak interactions based on the gauge group $SU(2) \times U(1)$ is matched only by our complete ignorance of the mechanism of symmetry breaking. In its simplest incarnation [1], the standard model contains one Higgs doublet. This theory contains in its spectrum a neutral scalar particle, the "Higgs". Its interactions with matter are all completely specified, but its mass, m_H, remains a free parameter. There exist both theoretical and experimental bounds on m_H. From violations of tree-level partial wave unitarity [2], $m_H \lesssim 1$ TeV, while experimentally static nuclear interactions [3,4] set the limit $m_H > 15$ MeV and the study of $K \to \pi + \text{Higgs}$ leads to [5] $m_H > 325$ MeV.

A limit $m_H > 409$ MeV has been deduced from $\eta' \to \eta H$, $H \to \mu^+\mu^-$ [6]. We do not think this is a firm bound because the theoretical branching ratio for $H \to \mu^+\mu^-$ depends on an unknown strong interaction parameter $a-b$ (see eqs. (20) and (21)). If $a-b>2$ the bound on m_H from η' decay is weaker than from K decay. The CUSB collaboration have recently set a limit $m_H > 3.9$ GeV by combining data from radiative decays of Υ and Υ' [7]. However, this limit is subject to theoretical uncertainties from both QCD and relativistic corrections and is therefore tentative until statistics are much improved [8].

In this paper we examine which values of the Higgs mass can be conclusively ruled out from the study of B meson decay. The rate for Higgs production in B-meson decay [9] grows with the fourth power of the top quark mass m_t. Moreover, a theoretical lower bound on the Higgs mass [10], $m_H \gtrsim 6.9$ GeV, can be violated only if $m_t \gtrsim 80$ GeV. Hence if the Higgs is lighter than the B-meson, it is produced copiously in B-decays.

This letter is organized as follows. We first compute the inclusive branching fraction for $B \to HX$ using a novel technique. We then review the theoretical lower bounds on m_H. This is followed by a study of the extent to which a light Higgs can be presently ruled out in B-meson decays. We conclude by summarizing our results.

The quark process underlying the decay of a B meson into a neutral Higgs and strange hadronic matter is $b \to sH$. Since the calculation of the rate for this process is central to our purposes and because this calculation has been the focus of an ongoing controversy [8], we reproduce the result of ref. [9] with a new method.

We first observe that the neutral Higgs field is the pseudo-Goldstone boson associated with scale invariance of the classical lagrangian [4]. This can be seen by observing that aside from the Higgs potential, the classical lagrangian is scale invariant. When we include the potential for the Higgs doublet, (Φ), $V(\Phi) = \lambda(\Phi^2 - v^2)^2$, we see that the Higgs field gets a VEV which breaks this invariance and that the Higgs is the associated Goldstone boson. We employ this

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observation to evaluate the flavor-changing Higgs coupling using PCDC (partially conserved dilatation current), which is very similar in spirit to the standard implementation of PCAC.

We define s^μ as the dilatation current associated with the *unshifted* lagrangian and call the (new improved) energy–momentum tensor $\theta_{\mu\nu}$, where $s_\mu = x^\nu \theta_{\mu\nu}$ so that $\partial^\mu s_\mu = \theta^\mu{}_\mu$. For the new improved energy–momentum tensor, dilation invariance is explicitly broken only by terms in the lagrangian with mechanical dimensions other than four. That is,

$$\theta^\mu{}_\mu(x) = \Delta(x) = (\mathscr{D}-4)\mathscr{L}(x) ,$$

where $\mathscr{D}$ counts the mechanical dimension of the fields. Since both $\partial\cdot s(=\theta^\mu{}_\mu)$ and Δ contains a tensor linear in the Higgs field, they are both good interpolating fields for the Higgs, and they are equal when evaluated for the on-shell field. That is,

$$\langle 0|\theta^\mu{}_\mu(0)|\mathrm{H}\rangle = \langle 0|\Delta(0)|\mathrm{H}\rangle = f_\mathrm{H} m_\mathrm{H}^2 , \tag{1}$$

where we take $f_\mathrm{H} = v = (\sqrt{2}G_\mathrm{F})^{-1/2} = 250$ GeV. We define

$$S(k^2) = \langle \mathrm{s}(p+k)|\Delta(0)|\mathrm{b}(p)\rangle , \tag{2}$$

$$T(k^2) = \langle \mathrm{s}(p+k)|\theta^\mu{}_\mu(0)|\mathrm{b}(p)\rangle , \tag{3}$$

$$Q(k^2) = (k^2-m^2)\langle \mathrm{s}(p+k)|H(0)|\mathrm{b}(p)\rangle , \tag{4}$$

where of course $Q(k^2)$ is the quantity we wish to evaluate. Because Δ is linear in the Higgs field, we obtain the relation

$$S(k^2) = \frac{-vm_\mathrm{H}^2}{k^2-m_\mathrm{H}^2} Q(k^2) , \tag{5}$$

$$Q(k^2) = \frac{-(k^2-m^2)}{vm^2} S(k^2)$$
$$= -\frac{k^2-m^2}{vm^2} T(k^2) , \tag{6}$$

where the last equality applies on-shell.

We now need to compute $S(k^2)$. One could do this explicitly. However, it is much simpler to compute with a PCDC analysis. We assume $Q(k^2)$ behaves smoothly so that we can extrapolate from $k^2=0$ to $k^2=m_\mathrm{H}^2$. The calculation at $k^2=0$ is simpler because of the theorem proven in ref. [4], which states on general grounds that

$$T(k^2) = C(k^2)k^2 \times \bar{\mathrm{u}}(p+k)[m_\mathrm{b}(1+\gamma_5)+m_\mathrm{s}(1-\gamma_5)]\mathrm{u}(p) + \mathrm{O}(m_\mathrm{b}^2)+\mathrm{O}(m_\mathrm{s}^2) . \tag{7}$$

Because of the explicit factor of k^2, T vanishes at $k^2=0$. But we have assumed that we can smoothly extrapolate $Q(k^2)$. Observe that at $k^2=0$

$$Q(0) = -(1/v)T(0) . \tag{8}$$

If relation (6) were true, we would conclude, as in ref. [4], that the process $\mathrm{b}\to\mathrm{sH}$ is very suppressed. Explicit calculations [9] violate this result.

The resolution of this problem parallels [13] the solution of the Sutherland–Veltman paradox for $\pi^0\to\gamma\gamma$. The trace-anomaly, $\partial\cdot s = \theta^\mu{}_\mu = \Delta$ is in fact anomalous [14]. The correct interpolating field for the Goldstone boson (the Higgs field) is

$$\partial\cdot s = \theta^\mu{}_\mu = \Delta + A , \tag{9}$$

where A is the anomaly. Our arguments above are true only if Δ is replaced by $\Delta - A$. Define

$$R(k^2) = \langle \mathrm{s}(p+k)|A|\mathrm{b}(p)\rangle . \tag{10}$$

The correct relation between Q and T is

$$Q(k^2) = -(1/v)[T(k^2)-R(k^2)] .$$

We apply the result at $k^2=0$ and use the result (7) that $T(k^2)|_{k^2=0}=0$ to derive

$$Q(0) = -(1/v)R(0) .$$

Therefore, the computation of the required matrix element has been reduced to the computation of the anomaly. We now outline the calculation in a mass independent subtraction scheme.

The renormalization structure of the theory may be analyzed in the symmetric phase without loss of generality. Integrating the anomalous trace identity (9) over space–time [11,12,14]

$$\left(-\sum_{k=1}^{n-1} p_k\cdot\frac{\partial}{\partial p_k} - 4(n-1) + \sum_{k=1}^{n} d_k\right) G^{(n)}(\{p\}) = -\mathrm{i}G_\Delta^{(n)}(0;\{p\}) - \mathrm{i}G_A^{(n)}(0;\{p\}) . \tag{11}$$

Here $G^{(n)}(\{p\})$ is any Green's function which depends on momenta $\{p\}=\{p_1,...,p_n\}$, where $\sum_{k=1}^n p_k = 0$, $G_\Delta^{(n)}(0;\{p\})$ ($G_A^{(n)}(0;\{p\})$) is the corresponding Green's function with an insertion of the operator Δ (A) at zero momentum, and d_k is the scaling dimension of the k-th field in $G^{(n)}$. Dimensional analysis then gives

$$\left(\mu\frac{\partial}{\partial\mu}+\sum_i m_i\frac{\partial}{\partial m_i}\right)G^{(n)}(\{p\})$$
$$=-\mathrm{i}G_{\Delta}^{(n)}(0;\{p\})-\mathrm{i}G_{A}^{(n)}(0;\{p\})\,, \qquad (12)$$

where μ is the renormalization scale and m_i are renormalized masses. There are no mass terms in the standard model in the symmetric phase, but it is useful to include them here since the method can then be used for other theories (e.g. QED). In a mass independent scheme

$$\sum_i m_i\frac{\partial}{\partial m_i}G^{(n)}(\{p\})=-\mathrm{i}G_{\Delta}^{(n)}(0;\{p\})\,. \qquad (13)$$

To see this, notice that eq. (13) always holds for bare quantities, and in mass independent schemes the renormalization functions, Z, Z_g and Z_m, depend only on the cut-off and the dimensionless coupling constants. Therefore

$$\mu\frac{\partial}{\partial\mu}G^{(n)}(\{p\})=-\mathrm{i}G_{A}^{(n)}(0;\{p\})\,. \qquad (14)$$

Following ref. [15] we write a general ansatz for the anomaly $\mathscr{A}=\sum_i c_i\mathscr{O}_i$ where the sum runs over all operators $\mathscr{O}_i$ of dimension four or less and the coefficients c_i are determined from (14) and the renormalization group equation

$$\left(\mu\frac{\partial}{\partial\mu}+\beta\frac{\partial}{\partial g}+m\gamma_m\frac{\partial}{\partial m}+n\gamma\right)G^{(n)}=0\,.$$

We find that the terms in $\mathscr{A}$ quadratic in the down-type quarks are (up to terms that vanish by the equations of motion)

$$\mathscr{A}=v(\bar{\mathrm{d}},\bar{\mathrm{s}},\bar{\mathrm{b}})\left[\tfrac{1}{2}(1+\gamma_5)\beta^{\mathrm{D}}+\tfrac{1}{2}(1-\gamma_5)\beta^{\mathrm{D}\dagger}\right]\begin{pmatrix}\mathrm{d}\\ \mathrm{s}\\ \mathrm{b}\end{pmatrix}.$$

Here $\beta^{\mathrm{D}}=\mu(\partial/\partial\mu)\lambda^{\mathrm{D}}$ is the beta-function for the Yukawa couplings λ^{D} that are associated with the masses of the down-type quarks [16]. One finds [15]

$$\beta^{\mathrm{D}}=(\tfrac{3}{2}/16\pi^2)\lambda^{\mathrm{U}}\lambda^{\mathrm{U}\dagger}\lambda^{\mathrm{D}}+\ldots\,,$$

where the "dots" indicate terms that are diagonalized as we go to a basis where the quark mass matrices are diagonal. Making this rotation we obtain

$$Q(0)=(3/256\pi^2)V_{\mathrm{st}}^{\dagger}V_{\mathrm{tb}}g^3(m_{\mathrm{t}}/M_{\mathrm{W}})^2$$
$$\times\bar{\mathrm{u}}(p+k)\left(\frac{m_{\mathrm{b}}}{M_{\mathrm{W}}}(1+\gamma_5)+\frac{m_{\mathrm{s}}}{M_{\mathrm{W}}}(1-\gamma_5)\right)\mathrm{u}(p)\,, \qquad (15)$$

where V is the KM matrix [17] and $8G_{\mathrm{F}}/\sqrt{2}=g^2/M_{\mathrm{W}}^2$. This is in agreement with the result of the explicit calculation in ref. [9]. Moreover, it explains the surprising cancellation of logarithms of $m_{\mathrm{t}}/M_{\mathrm{W}}$ observed there, since the beta-function has a power-series expansion in $\lambda_{\mathrm{t}}\propto m_{\mathrm{t}}/M_{\mathrm{W}}$. Now that we have calculated the $\mathrm{b}\to\mathrm{sH}$ matrix element we can compute the light Higgs branching ratio

$$\frac{\Gamma(\mathrm{b}\to\mathrm{sH})}{\Gamma(\mathrm{b}\to\mathrm{ce}\bar{\nu}_{\mathrm{e}})}=\frac{27\sqrt{2}}{64\pi^2}G_{\mathrm{F}}m_{\mathrm{b}}^2$$
$$\times\frac{(1-m_{\mathrm{H}}^2/m_{\mathrm{b}}^2)^2}{f(m_{\mathrm{c}}/m_{\mathrm{b}})}\left|\frac{V_{\mathrm{st}}^{\dagger}V_{\mathrm{tb}}}{V_{\mathrm{cb}}}\right|^2\left(\frac{m_{\mathrm{t}}}{m_{\mathrm{b}}}\right)^4, \qquad (16)$$

where $f(m_{\mathrm{c}}/m_{\mathrm{b}})\simeq0.5$ is the phase space factor for $\mathrm{b}\to\mathrm{ce}\bar{\nu}_{\mathrm{e}}$. We have normalized the rate by the semileptonic rate to reduce the dependence on m_{b} and because empirically $|V_{\mathrm{st}}^{\dagger}V_{\mathrm{tb}}/V_{\mathrm{cb}}|\simeq1$. With this rate and branching ratios computed below, we can now determine the extent to which a light Higgs can be excluded from the study of B meson decays.

If all quarks are much lighter than the W, Linde and Weinberg have shown that the desired SU(2) $\times$U(1) breaking vacuum is lower in energy than that without weak breaking only if $m_{\mathrm{H}}<6.9$ GeV [10]. This bound can be strengthened because in the standard cosmological model, the universe must actually evolve to this SU(2)$\times$U(1) breaking minimum. Guth and Weinberg have shown that this occurs in the absence of heavy quarks, only if $m_{\mathrm{H}}>9.7$ GeV [18]. However, these bounds are contingent upon the absence of heavy quarks. A sufficiently heavy ($\gtrsim M_{\mathrm{W}}$) quark introduces radiative corrections with opposite sign to the gauge boson contribution so that a light Higgs is allowed. In the standard model with three generations, the Higgs can be arbitrarily light only if the top quark mass exceeds about 80 GeV (if we also assume a standard cosmology). However, even if the top quark is light, the scalar potential of the standard model does have a local minimum with the desired

$SU(2) \times U(1)$ breaking for an arbitrarily light Higgs boson. Although these vacua have higher energy than those preserving $SU(2) \times U(1)$, they are sufficiently long lived (at least for [19] $m_H > 260$ MeV) that once the universe reaches this state, it will not tunnel to the lower energy vacuum on cosmological time scales. Hence from the viewpoint of acceptable particle physics today, there is nothing wrong with these vacua. The problem with them is purely cosmological: no plausible cosmology has been found which leads to these vacua. Certainly the standard big bang cosmology does not.

In the rest of this letter we will take the viewpoint of the standard hot big bang scenario, so that a Higgs boson can only be produced in B decay if the top quark mass is larger than 80 GeV [#1]. Hence the branching ratio for $B \to HX_s$ is given by [#2]

$$\frac{\Gamma(B \to HX_s)}{\Gamma(B \to All)} \geqslant 0.26 \left(1 - \frac{m_H^2}{m_B^2}\right)^2 . \qquad (17)$$

This branching ratio is very large and can be used to exclude certain values for the Higgs mass. Other values are excluded only if certain assumptions are made. Below we analyze each energy regime below the $\tau\bar{\tau}$ threshold.

If $2m_e < m_H < 2m_\mu$, then 26% of B decays would have a hard e^+e^- pair carrying half the energy of the B. This is more than twice the observed rate for single electrons, and at the very least would give too large a semi-leptonic branching ratio to electrons [21]. For $2m_\mu < m_H < 2m_\pi$ the Higgs decays nearly always to $\mu^+\mu^-$. This is excluded, for example, by the TASSO collaboration who give a limit for the inclusive branching ratio $B(B \to X\mu^+\mu^-) < 0.02$ [22].

For $m_H > 2m_\pi$, it is important to know the branching ratios of the Higgs $\mu^+\mu^-$, $\pi^+\pi^-$ and other open channels. However, because of the intrinsic uncertainties in the strong interaction calculation, we can only reliably compute these branching ratios at high energy, where perturbative QCD calculations are trustworthy, and at low energies, where the chiral lagrangian can be applied. At the quark level, the lowest order effective lagrangian for Higgs couplings to light quarks (u, d, s) and leptons (e, μ) and to photons and gluons is given by

$$\mathscr{L}_{\text{eff}} = -\sum_i m_i \bar{f}_i f_i (1 + H/v) + (\alpha_s N / 12\pi) G^A G^A H/v$$
$$+ (\alpha/8\pi) \tfrac{13}{3} FFH/v , \qquad (18)$$

where i runs over the light fermions f_i of mass m_i, G^A and F are the field strengths for gluons and photons and N is the number of heavy quarks, which we take to be three (c, b, t). We restrict our analysis to a Higgs lighter than 3.7 GeV, so that the decays to τ and c are not important. For m_H close to this upper value, the charm quark cannot really be considered to be heavy. This introduces known corrections [23] to eq. (18), but does not change our conclusions.

For sufficiently large m_H, which we take to be $m_H > 2$ GeV, the decay rates can be calculated in the spectator quark model. We find

$$\Gamma_{\mu\bar{\mu}} : F_{s\bar{s}} : \Gamma_{gg} = m_\mu^2 : 3m_s^2 : (\alpha_s/\pi)^2 m_H^2 , \qquad (19)$$

which are plotted as branching ratios in fig. 1 using $\alpha_s/\pi = 0.1$ and $m_s = 150$ MeV.

For small Higgs boson masses, the hadronic decay rates can be calculated using the chiral lagrangian

$$\mathscr{L}_{\text{C.L.}} = \tfrac{1}{4} f^2 \operatorname{Tr} \partial^\mu \Sigma \partial_\mu \Sigma^\dagger$$
$$+ f^2 [\tfrac{1}{2} \operatorname{Tr} \mu M (1 + H/v) \Sigma^\dagger + \text{h.c.}]$$
$$+ (H/v) \tfrac{1}{8} f^2 (a \operatorname{Tr} \partial^\mu \Sigma \partial_\mu \Sigma^\dagger + b \operatorname{Tr} \Sigma \Box \Sigma^\dagger + \text{h.c.}) , \qquad (20)$$

where $f = 93$ MeV, $\Sigma = \exp(2i\pi/f)$, $\pi = \Sigma_a T^a \pi^a$ and T^a are the SU(3) matrices, $\operatorname{Tr}(T^a T^b) = \delta^{ab}/2$. M is the quark mass matrix and the parameter μ is chosen to give the observed meson masses. The last term corresponds to the HG^AG^A term in the effective lagrangian of the quark and gluon theory. Notice that there are two real unknown strong interaction parameters a and b.

We expect the chiral lagrangian to give reliable results for $m_H \lesssim 800$ MeV. In this case the dominant hadronic decays of the Higgs are to $\pi^+\pi^-$ and $\pi^0\pi^0$. We calculate these to lowest non-trivial order and find the $\mu^+\mu^-$ branching ratio

[#1] It is possible to evade this by postulating the existence of a minimum in the non-perturbative region where λ, the quartic scalar coupling, is large. We think this is unreasonable since, in this case, there is no small parameter to make the Higgs light and one expects $m_H \simeq O(v)$.

[#2] Note that when m_H is not negligible compared with m_B the correction factor is the square of $(1 - m_H^2/m_B^2)$. One power comes from the square of the matrix element, and one from phase space. The matrix element suppression is frequently forgotten [7,9,20].

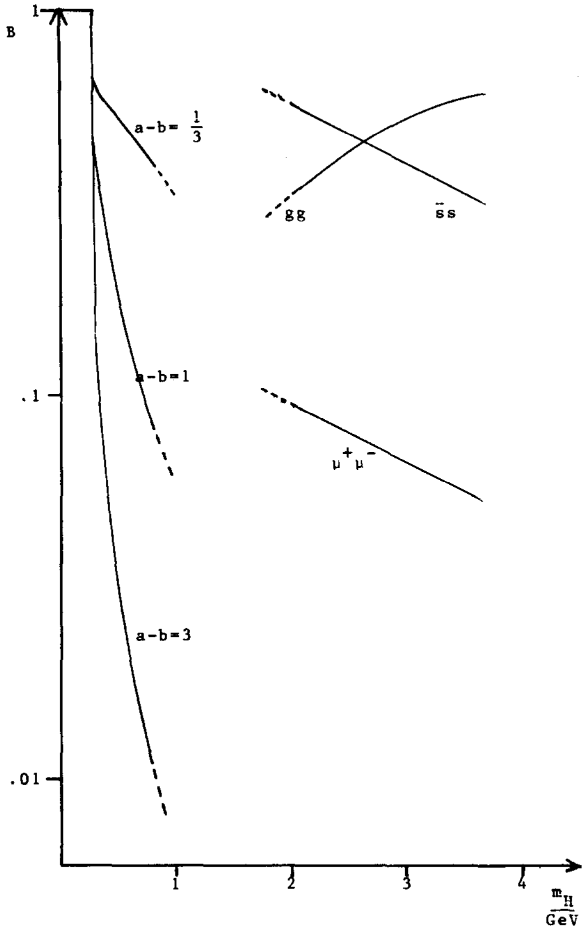

Fig. 1. Higgs branching ratios to $\mu^+\mu^-$ in the low energy regime and to $\mu^+\mu^-$, gg and $\bar{s}s$ in the higher energy regime. Note the high sensitivity to the parameter $a-b$ of the chiral lagrangian. In the quark model calculation the gg rate was calculated with $\alpha_s/\pi=0.1$.

$$
\begin{aligned}
B_{\mu^+\mu^-} = & m_\mu^2(1-4m_\mu^2/m_H^2)^{3/2} \\
& \times\{\tfrac{3}{16}m_H^2[2m_\pi^2/m_H^2+(a-b)(1-2m_\pi^2/m_H^2)]^2 \\
& \times(1-4m_\pi^2/m_H^2)^{1/2} \\
& +m_\mu^2(1-4m_\mu^2/m_H^2)^{3/2}\}^{-1} .
\end{aligned}
\tag{21}
$$

The result is plotted as a function of m_H for $a-b=\frac{1}{3}$, 1 and 3 in fig. 1. The result is clearly very sensitive to the unknown parameter $a-b$ [#3].

[#3] M. Voloshin [24] finds a particular value of $a-b=\frac{4}{9}$. We do not think this result should be trusted because of his use of perturbative QCD in this low energy process.

We do not know how to reliably calculate the Higgs branching ratios in the region between the validity of chiral perturbation theory and the quark model (800 MeV $\lesssim m_H \lesssim 2$ GeV). It is likely that the extrapolation is smooth, but without knowing $a-b$, it is hardly worth attempting.

It is straightforward to exclude the range 2 GeV $< m_H < 3.7$ GeV by using μ pair data from the ARGUS and CLEO collaborations. The ARGUS collaboration [25] have obtained a high-statistics sample of $B\to X\mu^+\mu^-$ while measuring the branching ratio for $B\to X\psi$ at the $\Upsilon(4S)$. They required the momentum of the $\mu^+\mu^-$ pair to be less than 2 GeV, so this data can only be used for Higgs heavier than 2.3 GeV, i.e. to exclude 2.3 GeV $< m_H < 3.7$ GeV. In their spectrum for the invariant mass of the $\mu^+\mu^-$ pair, a typical bin of width 50 MeV has 15 events, whereas a light Higgs would give $\simeq 1000(1-m_H^2/m_B^2)^2$. Hence the region 2.3 GeV $< m_H < 3.7$ GeV is easily excluded [#4].

The CLEO collaboration [26] have searched directly for $B\to HX$, $H\to\mu^+\mu^-$ and have a limit of $<5\times10^{-4}$ on the product of the branching ratios for $m_H > 3.2$ GeV. In this region with values of α_s and m_s as above, we predict the product of the branching ratios to be greater than 5×10^{-3}, giving an order of magnitude discrepancy between the theoretical prediction and experiment.

For 2 GeV $< m_H < 3.2$ GeV CLEO conclude, from an analysis with only one identified muon, that the product of branching ratios is less than [26] 8×10^{-3}. Using eq. (17) together with $B\mu^+\mu^-$ from the figure we find that this rules out the region 2 GeV $< m_H < 2.8$ GeV. The upper limit is sensitive to the value of α_s used in calculation $H\to gg$. For $\alpha_s=0.25$ (instead of 0.31 used in fig. 1) the upper limit would stretch to 3.2 GeV.

We now consider the region $2m_\pi < m_H < 2$ GeV. This region can only be excluded if assumptions are made. For example, the TASSO limit [22] of

[#4] The ARGUS collaboration also had a cut on the angle θ between the μ^+ and μ^-: $|\cos\theta| < 0.9$. This cut will not significantly affect the number of events from the decay of a Higgs with mass between 2 and 3.7 GeV. The Higgs is heavy enough that the μ^+ and μ^- must have a large opening angle. They will only be going back to back if the Higgs is produced at rest in the laboratory. Although the momentum spectrum of the Higgs is not known it will not be sufficiently peaked around zero to cut the signal by two orders of magnitude.

$B(B\to X\mu^+\mu^-)<0.02$ excludes $2m_\mu<m_H<800$ MeV if $|a-b|<1$. As $|a-b|$ increases the upper limit drops; for $|a-b|=3$, the excluded region does not go much beyond $2m_\pi$. For $m_H>500$ MeV, the CLEO limit [26] of 8×10^{-3} on the product of branching ratios for $B\to HX$, $H\to\mu^+\mu^-$ is a factor 2.5 stronger than the TASSO limit, but the uncertainty in $a-b$ still prevents any strong conclusion.

The CLEO collaboration have recently reported on a search for the Higgs meson in exclusive two-body decays of the B meson [27] $B\to KH$, K^*H. Excellent limits on various branching ratio products have been placed. For example for 0.3 GeV $<m_H<$ 2.8 GeV:

$$B(B\to HK^-)B(H\to\mu^+\mu^-)<1.5\times10^{-4}, \tag{22a}$$

$$B(B\to HK^-)B(H\to\pi^+\pi^-)<3\times10^{-4}. \tag{22b}$$

For 0.3 GeV $<m_H<$ 1 GeV, $B(H\to\mu^+\mu^-)+B(H\to\pi^+\pi^-)+B(H\to\pi^0\pi^0)=1$, so (22) together with isospin give

$$B(B^-\to HK^-)<6\times10^{-4}.$$

Using eq. (17) gives

$$R=\frac{\Gamma(B^-\to K^-H)}{\Gamma(B^-\to X_s^-H)}<2.3\times10^{-3}. \tag{23}$$

Because we have added the branching ratios to pions and muons, the only theoretical uncertainty is the branching fraction R of B into exclusive modes. Naively, one expects $R\sim(\alpha_s/4\pi)^2\sim10^{-3}$ since a hard gluon must be exchanged with the spectator quark for a bound state to form. Using the quark model R has been estimated to be 0.07 by Haber et al. [20]. However, these calculations are very uncertain, especially for the light Higgs case where the non-relativistic approximation cannot be justified and where there is high sensitivity to the component of the meson wavefunction with large relative quark momentum. Using the methods of Grinstein et al. [28] we find

$$R=\frac{0.35}{(1-m_H^2/m_B^2)}\exp\left[-\frac{1.8}{\kappa^2}\left(1-1.22\frac{m_H^2}{m_B^2}\right)\right]. \tag{24}$$

For $\kappa=1$ this gives $R=0.06$ in good agreement with Haber et al. However, there is exponential sensitivity to the phenomenological parameter κ^2. Although we do not know how much κ^2 differs from unity it could easily be by a factor of 2 (as, for example, occurs in K decays [28]) giving $R=0.009$. Using these estimates it is tempting to say that eq. (23) rules out $0.3<m_H<1$ GeV. However, given the lack of validity of the non-relativistic quark model, and the exponential sensitivity of R on κ^2 we do not believe this to be a firm basis for excluding this region of Higgs mass. The CLEO data on the two-body decays of B mesons could also exclude the mass range 1 GeV $<m_H<$ 2 GeV [21]. This would require knowing Higgs branching ratios as well as R. In this mass region, their data sets limits:

$$B(H\to\mu^+\mu^-)<8\times10^{-4}/R, \tag{25a}$$

$$B(H\to\pi^+\pi^-)<1.2\times10^{-3}/R, \tag{25b}$$

$$B(H\to K^+K^-)<8\times10^{-4}/R, \tag{25c}$$

where we have used eq. (17). These exclusive modes can only be used to exclude the Higgs when we have better understanding of the Higgs branching fractions and of R.

It is remarkable that even with $B\to XH$ having a branching ratio as large as 0.26, it is only possible to definitely exclude $m_H<300$ MeV and 2 GeV $\lesssim m_H\lesssim$ 3.7 GeV. The window between 300 MeV and 2 GeV largely reflects our inability to calculate reliable Higgs decay branching ratios in this region. Above 3.7 GeV, $H\to\bar\tau\tau$, $\bar cc$ dominate, for which experiments have low sensitivity.

In this paper, we have obtained the one loop amplitude for $b\to sH$ by considering the Higgs to be a dilaton and studying the anomaly in the dilatation current. The result is proportional to the β function for the down-quark Yukawa couplings $\beta^D\propto\lambda^U\lambda^{U\dagger}\lambda^D$ and agrees with the result from the usual Feynman diagram expansion. The corresponding rate for $b\to sH$ is proportional to m_t^4. In the three-generation standard model the Higgs can only be light enough to appear in B decay if $m_t\gtrsim80$ GeV, giving a branching ratio for $B\to X_sH$ in excess of 26%. Present data then excludes Higgs masses below 3.7 GeV, with the exception of a window between 300 MeV and 2 GeV. This window can only be excluded using present data with further assumptions. It is important to conduct a Higgs search to study this mass range. A high statistics inclusive invariant mass plot for $\pi^+\pi^-$ and K^+K^- at the $\Upsilon(4S)$ would be conclusive, since the theoretical interpretation is not subject to large uncertainties. For a Higgs mass between 1 GeV and 2 GeV a search can be made at the $\Upsilon(4S)$ for events with six kaons. Backgrounds are effectively removed by requiring two K pairs to have the same invariant mass.

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