B-meson rare decays in two-Higgs-doublet models

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(Received 7 March 1988)

We study $B \to l^+l^-$, $b \to sl^+l^-$, and $B \to Kl^+l^-$ decays in two-Higgs-doublet models. With constraints from $B \to K^*\gamma$ for the parameters in the model, we find that the branching ratio for $B \to l^+l^-$ can be two to three orders of magnitude larger than in the standard model, if the Higgs pseudoscalar is very light. The branching ratios for $b \to sl^+l^-$ and $B \to Kl^+l^-$ can be an order of magnitude larger than the standard-model predictions.

Recently there have been many studies of the B-meson rare decays $B \to l^+l^-$, $B \to Kl^+l^-$, and $b \to sl^+l^-$ (Ref. 1). These processes are interesting for several reasons. (i) They occur through one-loop interactions and can thus provide useful information about heavy quarks in the loop, consequently testing models beyond the tree level. (ii) Since in these processes the decaying b quark is heavy, long-distance effects are expected to be small compared with short-distance effects and so these processes can be computed reliably. (iii) These decays are not as rare as similar kaon processes. Since the B meson is much heavier than the kaon, the branching ratios can be relatively high in the standard model. The quantity $B(B \to \tau^+\tau^-)$, for example, which is proportional to the B-meson mass, is in the range $\approx 10^{-9}$–10^{-6}. The branching ratios of the two semileptonic decays we will be considering are proportional to the fifth power of the B-meson mass. Their values in the standard model are in the ranges $B(B \to Kl^+l^-) \approx 4\times10^{-7}$–$3\times10^{-6}$ and $B(b \to sl^+l^-) \approx 2\times10^{-6}$–$1.5\times10^{-5}$ (Ref. 1). These decays are, therefore, experimentally reachable in the near future.

In this paper we carry out calculations for these decays in both versions of two-Higgs-doublet models.[2] In addition to the standard contributions to these processes, there are contributions due to the extra Higgs particles that exist in these models. We find that the branching ratio for $B \to \tau^+\tau^-$ in one of these models can be as high as 10^{-7}, even when the top-quark mass is as low as 40 GeV. This is an increase by 2 orders of magnitude over the standard-model prediction. The branching ratios for $B \to Kl^+l^-$ and $b \to sl^+l^-$ can be 1 order of magnitude larger.

We will denote the two Higgs doublets in the model by $\phi_1 = (\phi_1^0, \phi_1^-)$ and $\phi_2 = (\phi_2^0, \phi_2^-)$. After spontaneous symmetry breaking at $\sim M_W$, there exists one charged physical Higgs boson H^+ and one neutral physical pseudoscalar Higgs boson P^0. These two particles can introduce new contributions to $B \to l^+l^-$, $B \to Kl^+l^-$, and $b \to sl^+l^-$ through "box-" and "penguin-" type diagrams as shown in Fig. 1. [For the moment we do not distinguish between B_d and B_s; d may be replaced by s in Eqs. (3), (5), (8), and (9) which follow.]

The Lagrangian describing the relevant H^+ interactions is[2]

$$\mathcal{L}_H = \frac{g}{\sqrt{2}M_W} H^+ \bar{U} V_{\mathrm{KM}} (\beta M_u \gamma_- + \delta M_d \gamma_+) D \tag{1}$$

and the P^0 Lagrangian is[3]

$$\begin{aligned}\mathcal{L}_P = {}& \frac{g}{2M_W} iP^0(\beta \bar{U} M_u \gamma_5 U + \delta \bar{D} M_d \gamma_5 D + \delta \bar{l} M_l \gamma_5 l) - \frac{g}{\sqrt{2}M_W} \phi_1^+ \bar{U} V_{\mathrm{KM}} (M_u \gamma_- + M_d \gamma_+) D \\ & + \frac{g}{\sqrt{2}} W_\mu^+ \bar{U} \gamma^\mu \gamma_- V_{\mathrm{KM}} D + \frac{g}{2} W_\mu^+ (P^0 \gamma^\mu H^- - H^- \gamma^\mu P^0) + \frac{g(M_H^2 - M_P^2)}{2M_W} iP^0 H^+ \phi_1^- + \text{H.c.} ,\end{aligned} \tag{2}$$

where for convenience the 't Hooft gauge has been used. $\phi_1^\pm$ is a fictitious field; M_H and M_P are the masses of H^+ and P^0, respectively; $U = (u,c,t)$, $D = (d,s,b)$, $l = (e,\mu,\tau)$ and $\nu = (\nu_e, \nu_\mu, \nu_\tau)$; V_{KM} is the Kobayashi-Maskawa matrix. The two different models are distinguished by the choices for β and δ. Model (i) has no ϕ_2 Yukawa terms with U, D, and l and so $\beta = -\delta = \zeta/\eta$, where ζ and η are the vacuum expectation values of ϕ_1^0 and ϕ_2^0, respectively. Model (ii) has ϕ_1 coupling to U only and ϕ_2 coupling to D and l and thus, $\beta = 1/\delta = \zeta/\eta$. We first discuss the $B \to l^+l^-$ decay. The standard-model short-distance contribution to this process is[4]

$$\mathcal{L}_{\mathrm{st}} = iM_W^2 \left[\frac{g}{2\sqrt{2}M_W} \right]^4 C(x_i) V_{ib}^* V_{id} M_l f_B \bar{l} \gamma_5 l , \tag{3}$$

where M_l, M_i, and M_W are the lepton, internal quark, and W gauge-boson masses, respectively; $x_i = (M_i/M_W)^2$,

V_{ij} are the Kobayashi-Maskawa matrix elements, and

$$C(x)=\frac{x}{4}+\frac{3x}{4(1-x)}+\frac{3x^2\ln x}{4(1-x)^2} . \tag{4}$$

We have made use of $\langle 0|\bar{d}\gamma_5\gamma^\mu b|B\rangle=if_B P_B^\mu$. In the two-Higgs-doublet model, after calculating Fig. 1(a), we obtain the following effective Lagrangian for $B\to l^+l^-$:

$$\mathcal{L}_{2H}=\frac{i}{16\pi^2}\left[\frac{g}{2\sqrt{2}M_W}\right]^4\left[\frac{M_l^3}{M_H^2}f_B\delta^2(\beta^2M_i^2-\delta^2M_bM_d)I(y_i)V_{ib}^*V_{id}\bar{l}\gamma_5 l\right.$$
$$\left.+\frac{i}{128\pi^2}(g/M_W)^4M_lM_i^2f_B\delta\beta[G_1(x_i,y_i,z)+\beta^2G_2(y_i)]\frac{M_B^2}{M_P^2-M_B^2}V_{ib}^*V_{id}\bar{l}l\right] , \tag{5}$$

where M_B is the B-meson mass, $y_i=(M_i/M_H)^2$, $z=(M_H/M_W)^2$, and

$$I(y)=\frac{1}{1-y}+\frac{y\ln y}{(1-y)^2} ,$$
$$G_1(x,y,z)=\frac{x-2}{2(1-x)}\left[1+\frac{\ln x}{1-x}\right]+\frac{2}{1-z}\left[\frac{\ln x}{1-x}-\frac{\ln y}{1-y}\right]$$
$$+\tfrac{1}{2}(1-z)\left[-\frac{z\ln z}{(1-x)(1-z)^2}-\frac{1}{(1-y)(1-z)}-\frac{y\ln y}{(1-x)(1-y)^2}\right.$$
$$\left.+\frac{y}{2(1-x)(1-y)}-\frac{1}{2(1-z)(1-x)^2}+\frac{\ln y}{2(1-z)(1-y)^2}\right] , \tag{6}$$
$$G_2(y)=\frac{y}{2(1-y)}\left[1+\frac{\ln y}{1-y}\right] .$$

From these effective Lagrangians, we easily obtain the decay rate $\Gamma_{2H}(B\to l^+l^-)$ in the two-Higgs-doublet model,

$$\Gamma_{2H}(B\to l^+l^-)=R\,\Gamma(B\to l^+l^-)_{\rm st} , \tag{7}$$

where, as calculated from Eq. (3),

$$\Gamma(B\to l^+l^-)_{\rm st}=\frac{G_F^4M_W^4}{32\pi^5}f_B^2M_BM_l^2(1-4M_l^2/M_B^2)^{1/2}|C(x_i)V_{ib}^*V_{id}|^2 \tag{8}$$

and

$$R=|M_W^2C(x_i)V_{ib}^*V_{id}|^{-2}\left[\left|\frac{M_l^2I(y)}{16M_H^2}\delta^2(\beta^2M_i^2-\delta^2M_bM_d)+M_W^2C(x_i)V_{ib}^*V_{id}\right|^2\right.$$
$$\left.+(1-4M_l^2/M_B^2)\frac{M_B^4}{(M_P^2-M_B^2)^2}|\tfrac{1}{2}\delta\beta M_i^2(G_1+\beta G_2)V_{ib}^*V_{id}|^2\right] . \tag{9}$$

R is a measure of the deviation of the two-Higgs-doublet model from the standard model. From the expression for R we see that the dominating contribution is the second term in Eq. (9) (that is, the one from P^0-induced effects) if the P^0 is very light. In our later discussion we will take M_P to be zero. A motivation for having a very light P^0 is connected with a possible solution of the strong CP problem.[5] The first term contains the H^+-induced contributions, which we see to be suppressed by $(M_l/M_H)^2$. From the way the vacuum expectation values (VEV's) enter into Eq. (9) it is clear that the new contribution from model (ii) is smaller than model (i), if $\xi/\eta>1$. We will place emphasis on model (i) for this reason. We notice that the decay width is approximately proportional to the lepton mass squared. Thus, we deduce that roughly

$$\Gamma(B\to e^+e^-):\Gamma(B\to\mu^+\mu^-):\Gamma(B\to\tau^+\tau^-)\approx M_e^2:M_\mu^2:M_\tau^2 . \tag{10}$$

Also, since $B_d\to\tau^+\tau^-$ and $B_s\to\tau^+\tau^-$ are dominated by the top-quark internal loop, we have that

$$\frac{B(B_d\to\tau^+\tau^-)}{B(B_s\to\tau^+\tau^-)}\approx\frac{|V_{\rm dt}|^2}{|V_{\rm st}|^2}\approx\frac{s_1^2s_2^2}{|V_{\rm cb}|^2}<1 , \tag{11}$$

where $s_{1,2}\equiv\sin\theta_{1,2}$ (KM angles). We thus find that $B(B_s\to\tau^+\tau^-)$ will be the largest branching ratio among

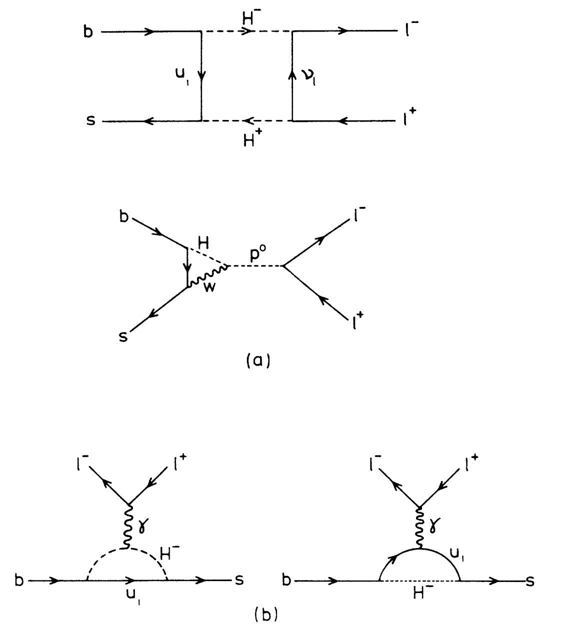

FIG. 1. (a) Generic diagrams for new contributions to $B \to l^+ l^-$ in the two-Higgs-doublet models. (b) Generic diagrams which together with those in (a) give all new contributions to $b \to s l^+ l^-$.

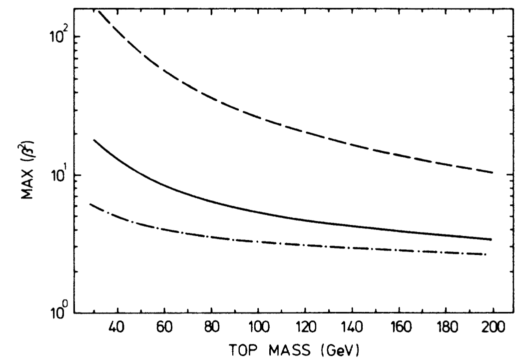

FIG. 2. The upper bound on β^2 as a function of M_t due to $B \to K^* \gamma$. The dotted-dashed, solid, and dashed lines correspond to $M_H = 30$, 100, and 500 GeV, respectively.

the family of $B \to l^+ l^-$ decays. Note that Eqs. (10) and (11) are approximately valid for both the standard model and the two-Higgs-doublet model. To calculate the branching ratio for $B_s \to \tau^+ \tau^-$, we normalize the total B_s width to the semileptonic decay width of B_s (Ref. 6),

$$\Gamma(B_s \to \text{all}) \approx \Gamma(b \to \text{all}) = \frac{|V_{cb}|^2 G_F^2 m_b^5}{192\pi^3 B_{SL}} f(m_c/m_b) X_{QCD}$$
$$\approx 3.3 \frac{|V_{cb}|^2 G_F^2 m_b^5}{192\pi^3}, \tag{12}$$

where B_{SL} is the semileptonic branching fraction of the B meson, and

$$f(x) = 1 - 8x^2 + 8x^6 - x^8 - 24x^4 \ln x ,$$
$$X_{QCD} = 1 - [2\alpha_s(M_b)/3\pi](\pi^2 - \tfrac{25}{4}) . \tag{13}$$

Using $M_B = 5.3$ GeV, $M_b = 4.8$ GeV, and $f_B = 0.1$ GeV we obtain

$$B(B_s \to \tau^+ \tau^-) = 3.5 \times 10^{-7} |C(x_t)|^2 R . \tag{14}$$

This branching ratio depends on the top-quark mass, the charged-Higgs-boson mass, and the ratio of VEV's. If one assumes that the top quark has the standard branching ratio for the $t \to b e^+ \nu_e$ decay then the top-quark mass is bounded by experimental data from UA1: $M_t > 44$ GeV (Ref. 7). However, it is possible to lower this bound in the two-Higgs-doublet model, provided that the decay mode $t \to H^+ b$ is kinematically allowed.[8] In this case it is possible for the bound to be as low as the KEK TRISTAN value, $M_t > 26$ GeV (Ref. 9). An upper bound for M_t of 200 GeV can be obtained from neutral-current analyses.[10] We will take M_t as a free parameter ranging from 25 to 200 GeV. It is understood that when $M_t < M_H + M_b$ the UA1 lower bound is to be invoked. Constraints on β can be obtained from several sets of experimental data. From B_S^0-B_L^0 mass-difference data it can be shown that $\beta^2 < 4.1 M_H/M_t - 12 M_H/M_t$ provided that $(M_t/M_H)^2 \ll 1$ (Ref. 11). However, this bound depends on other parameters, for example, the bag factor B, which makes this bound unreliable. In Ref. 12, by the use of data from $B \to K^* \gamma$ and by assuming $B(b \to s\gamma) \le 2 \times B(B \to K^* \gamma)$, it was found that

$$\beta^2 \le (1.95 - |F_2|)/|H| , \tag{15}$$

where

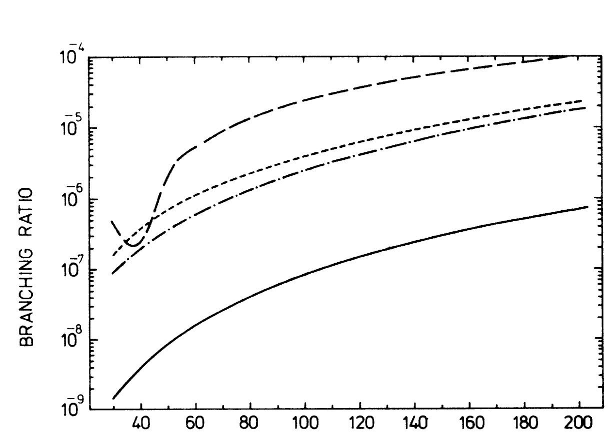

FIG. 3. The solid line is the standard-model prediction for $B(B_s \to \tau^+ \tau^-)$. The dotted-dashed, short-dashed, and long-dashed lines are the upper bounds on this branching ratio in model (i) for $M_H = 30$, 100, and 500 GeV, respectively.

$$H(y)=\beta^2 H_1(y)+\delta\beta H_2(y) ,$$

$$H_1(y)=\frac{y}{12(1-y)^3}\left[Q_i\left[2+5y-y^3+\frac{6y\ln y}{1-y}\right]-\left[-1+5y+2y^2+\frac{6y^2\ln y}{1-y}\right]\right] , \tag{16}$$

$$H_2(y)=\frac{y}{2(1-y)^3}[-Q_i(3-4y+y^2+2\ln y)+1-y^2+2y\ln y] ,$$

$$F_2(x)=-Q\left[-\frac{1}{4}+\frac{1-5x-2x^2}{4(1-x)^3}-\frac{3x^2\ln x}{2(1-x)^4}\right]+2T_3\left[-\frac{3}{4}+\frac{3-9x}{4(1-x)^2}-\frac{3x^2\ln x}{2(1-x)^3}\right] ,$$

where Q_i is the electric charge of the internal fermion, and Q and T_3 are the electric charge and the weak charge of the external quark. This bound is stronger than the one from B_S^0-B_L^0 mass difference. Constraints on β can also be obtained from $K_L\to\mu^+\mu^-$ by doing an analysis similar to $B\to l^+l^-$. However, because of long-distance effects we can only obtain an upper bound on β which is slightly weaker than bounds obtained from $B\to K^*\gamma$. Consequently in our later discussions we will use constraints on β from $B\to K^*\gamma$. The bounds are plotted as a function of M_t in Fig. 2. This plot relates to model (i) since not only is the new contribution to the decay larger than in model (ii), but also the above bound is more stringent due to a larger form factor H. Throughout the paper we plot all quantities of interest as functions of M_t with some representative values of M_H: 30, 100, and 500 GeV.

Using the bounds from Fig. 2, the upper bounds on $B(B_s\to\tau^+\tau^-)$ in model (i) and the standard-model prediction are plotted in Fig. 3. The bounds, in general, are increasing functions of M_H and M_t. Note also that in the two-Higgs-doublet model $B(B_s\to\tau^+\tau^-)$ can be as large as 10^{-4}, and values of the order of 10^{-7} are possible even for small M_t (≈ 30 GeV), which is an increase of 2 orders of magnitude over the standard-model result. Thus, this decay may provide a good test of this model.

We now discuss the decays $b\to sl^+l^-$ and $B\to Kl^+l^-$. The effective photon-exchange Lagrangian in two-Higgs-doublet models, from Fig. 1(b), is[12]

$$\mathcal{L}=\frac{eg^2}{32\pi^2 M_W^2}\sum_i V_{ia}V_{ib}^*(\bar{B}V'^{\mu}A)A_\mu ,$$
$$V'^{\mu}=iq_\nu\sigma^{\mu\nu}(M_a\gamma_+ + M_b\gamma_-)(F_2-H)+(q^2\gamma^\mu-q^\mu \not{q})(F_1+\beta^2 y_i G)\gamma_- . \tag{17}$$

F_1 and F_2 are contributions due to the W boson in the standard model and are given in Ref. 4. H and G are contributions[12] due to H^+. H is given in Eq. (16) and $E(y)=\beta^2 yG(y)$ is given by

$$E(y)=\frac{y\beta^2}{36(1-y)^3}\left[Q_i\left[-16y+29y^2-7y^3-\frac{6y(2-3y)\ln y}{1-y}\right]+\left[-2+7y-11y^2-\frac{6y^3\ln y}{1-y}\right]\right] . \tag{18}$$

Diagrams in Fig. 1(a) also contribute to this process, but is suppressed by a factor $(M_l/M_W)^2$ compared with the contribution from Fig. 1(b). We will neglect this term in our subsequent discussion. We also neglect the small contribution due to Z^0 exchange, to obtain[1]

$$\Gamma(b\to sl^+l^-)=\frac{G_F^2 M_b^5}{192\pi^3}\left[\frac{g^2}{16\pi^2}\right]^2 T ,$$
$$T=|\lambda_i A_i|+|\lambda_i B_i|+4s_W^2\lambda_i(A_i+B_i)\lambda_j F'_{2j}+16s_W^4[\ln(M_b/2M_l)-\tfrac{2}{3}]\,|\lambda_i F'_{2i}|^2 , \tag{19}$$

where $s_W\equiv\sin\theta_W$ ($\theta_W=$ Weinberg angle) and

$$B_i=-s_W^2[F'_{1i}+2C^z(x_i)],\quad A_i=C(x_i)+B_i ,$$
$$F'_{2i}=F_2(x_i)-H(y_i),\quad F'_{1i}=F_1(x_i)+E(y_i) , \tag{20}$$
$$\lambda_i=V_{is}^* V_{ib} .$$

The form factors C and C^z can be found in Ref. 4. Setting E and H to zero restores the standard-model result. To obtain the decay amplitude for $B\to Kl^+l^-$ we use

$$A(B\to Kl^+l^-)=\langle Kl^+l^-|\mathcal{L}|B\rangle . \tag{21}$$

We thus have[13]

$$\Gamma(B\to Kl^+l^-)=\frac{G_F^2 M_b^5}{192\pi^3}\left[\frac{g^2}{16\pi^2}\right]^2\frac{f_+(0)}{8}\times[\,|\lambda_i a_i|^2+|\lambda_i C(x_i)|^2] ,$$
$$a_i=C(x_i)-2s_W^2[F'_1+2C^z(x_i)]-2(M_b/M_B)^2 s_W^2 F'_{2i} . \tag{22}$$

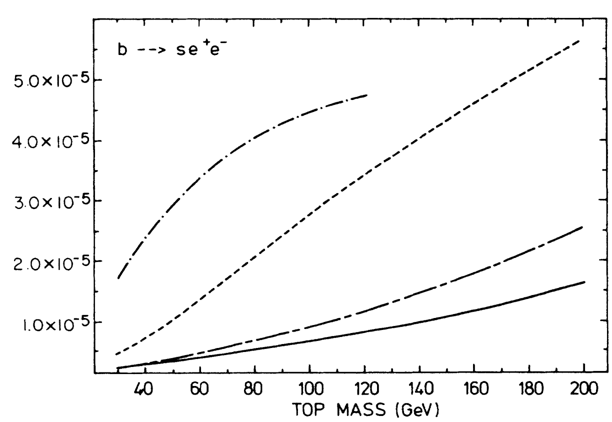

FIG. 4. Variation of $B(b\to se^+e^-)$ with respect to M_t and M_H for a fixed value of $\beta^2=3$. The short-dashed–long-dashed, dashed, and dotted-dashed lines correspond to $M_H=500$, 100, and 30 GeV in model (i). The 30-GeV line is cut off at $M_t\approx 120$ GeV because of the bound from $B\to K^*\gamma$ (see Fig. 2). The solid line is the standard-model prediction.

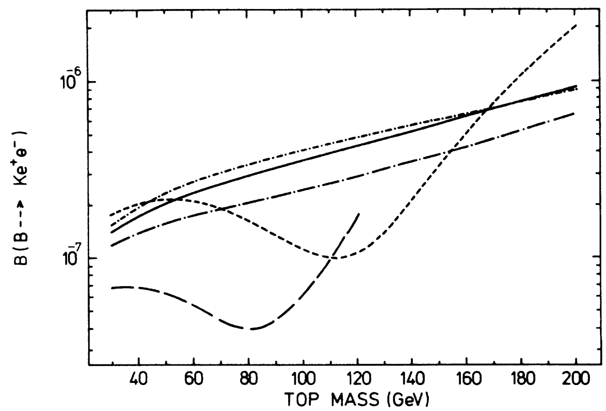

FIG. 5. The solid line is the standard model $B(B \to Ke^+e^-)$. The long-dashed and dotted-long-dashed lines are predictions in model (i) for $M_H=30$ and 100 GeV, respectively. The short-dashed and dotted-short-dashed lines are predictions in model (ii) for $M_H=30$ and 100 GeV, respectively. $\beta^2=3$ for all the two-Higgs-doublet model lines.

To obtain Eq. (22) we have used

$$\langle K \mid \bar{s}\gamma^{\mu} b \mid B \rangle = f_+(q^2)(P_B+P_K)_\mu + f_-(q^2)(P_B-P_K)_\mu \ , \tag{23}$$

$$\langle K \mid \bar{s} i q_\nu \sigma^{\mu\nu} b \mid B \rangle = h(q^2)(P_{B\mu}P_{K\nu}-P_{B\nu}P_{K\mu})q^\nu \ .$$

Assuming B^* vector-meson dominance, nonrelativistic bound-state quark dynamics, and the spectator model, we have[14]

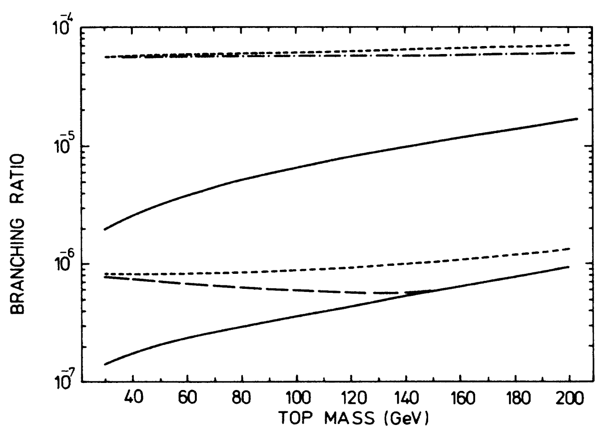

FIG. 6. The upper and lower solid lines are the standard-model predictions for $B(b \to se^+e^-)$ and $B(B \to Ke^+e^-)$, respectively. The upper short-dashed line is the upper bound on this branching ratio for both $M_H=100$ and 500 GeV. The dotted-dashed line next to that is the upper bound for $M_H=30$ GeV. The two lower short-dashed and long-dashed lines are upper bounds on $B(B \to Ke^+e^-)$ for $M_H=30$ and 500 GeV, respectively.

$$f_+(q^2) = \left[\frac{M_{B^*}^2}{M_{B^*}^2 - q^2} \right] f_+(0) \ ,$$

$$h(q^2) = 2(M_s+M_b) f_+(q^2)/M_{B^*}^2 \ , \tag{24}$$

$$f_+(0) = \frac{f_B M_B}{f_K M_{B^*}} \ .$$

Normalizing the total width of the B meson to the semileptonic decay width, we obtain the branching ratios

$$B(b \to sl^+l^-) = 1.93 \times 10^{-6} T \ , \tag{25}$$

where T is given in Eqs. (19) and (20) and

$$B(B \to Kl^+l^-) = 2.4 \times 10^{-7} [\,|a|^2 + |C(x)|^2] f_+^2(0) \ . \tag{26}$$

The results are plotted in Figs. 4–6.

For $b \to sl^+l^-$, the situation is similar to that of $B_s \to \tau^+\tau^-$ in that the enhancement in model (i) is much larger than that in model (ii), so we again focus on model (i). Figure 4 shows the variation of $B(b \to se^+e^-)$ with M_H and M_t for a fixed value $\beta^2=3$. The curve for $M_H=30$ GeV is cut at $M_t \approx 120$ GeV due to the bound in Fig. 2.

For $B \to Kl^+l^-$, however, there is a partial cancellation between the H^+ contribution and the W contribution in model (i), whereas in model (ii) they add. This effect is exactly opposite to that in $b \to sl^+l^-$. This occurs because the H form factor in model (i) always has the opposite sign to that in model (ii). Note also the sign difference of the F_2 terms in Eqs. (19) and (22). To illustrate this we show the deviations from $B(B \to Ke^+e^-)_{\rm st}$ in both models in Fig. 5, for a common value $\beta^2=3$. Note the difference between predictions for small M_H (≈ 30 GeV) and relatively large M_H (>100 GeV). The explanation lies in the fact that for relatively large M_H the form factor E is negligible compared to H, while for small M_H it is not. The curve for $M_H=30$ GeV in model (i) is cut at $M_t=120$ GeV due to the bound in Fig. 2. Finally, Fig. 6 shows the upper bounds for $B(b \to se^+e^-)$ and $B(B \to Ke^+e^-)$ in model (i). The upper bound on $B(b \to se^+e^-)$ is essentially independent of M_H, because the numerical factor $16s_W^4[\ln(M_b/2M_e)-\frac{2}{3}]$ appearing in the term $|F_2|^2$ in Eq. (19) is so large that the maximum branching ratio is well approximated by taking the maximum value of F_2. From the experimental bound on $B \to K^*\gamma$ we have $F_2 \le 1.95$. For $B \to Ke^+e^-$ this also holds for relatively large M_H (>100 GeV). However, for small M_H (≈ 30 GeV) the form factor E actually dominates, because its numerical factor is about twice as large as that of F_2 and also because $E > H$ for small M_H [see Eq. (22)].

In conclusion, then, we have demonstrated that present bounds on the ratio of VEV's β in the two-Higgs-doublet model permit a substantial increase in the branching ratios of rare B-meson decays compared with standard-model values. We await with interest the new generation of experiments which may discover evidence for an extended Higgs sector.

X.-G.H. would like to thank Professor B. H. J. McKellar for valuable discussions. We would like to acknowledge support by the Australian Research Grants Committee.

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