B DECAYS IN THE STANDARD MODEL AND BEYOND

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1. INTRODUCTION

When Vera Lüth asked me to give a talk on B decays in and beyond the standard model (SM), I readily accepted. However, when I sat down and made a list of the topics I would have to cover, I quickly realized that I had bitten off more than I could chew. My list consisted of the following subjects:

- Semileptonic B decays: these are typically described in one of two ways. Either one picks a specific model,[1] or one uses the Heavy Quark Effective Theory (HQET);[2]
- Hadronic B decays: such decays are usually described by the BSW model;[3]
- Right-handed B decays:[4] the suggestion here is that B decays are mediated not by the ordinary W, but rather by a right-handed W_R;
- Rare B decays: included are the flavour-changing neutral-current decays $b \to s\gamma$, $b \to s\ell^+\ell^-$, $b \to s\nu\bar{\nu}$, $b \to sg$, $b \to sq\bar{q}$, and $B^0 \to \ell^+\ell^-$, as well as hadronic penguins ($B \to K\pi$, etc.), and $B^0 \to \gamma\gamma$;
- the decay $B_u^+ \to \ell^+\nu$;
- Exotic B states such as B_c's and Λ_b's;
- B-$\overline{B}$ mixing – x_d and x_s;
- T Violation (triple products),

and I'm sure I've overlooked some other possibilities. Given the length of this list, I realized that I would have to limit myself to a subset of the above topics. I therefore decided to discuss only right-handed B decays, certain rare B decays, B_c decays, B_s^0-$\overline{B_s^0}$ mixing, and T violation. Some of the other subjects, such as HQET and B baryons, are discussed elsewhere in these proceedings.[5]

2. RIGHT-HANDED B DECAYS

Gronau and Wakaizumi[4] (GW) have suggested that B decays might in fact be mediated by a right-handed W_R, instead of the SM left-handed W. This possibility is predicated on two facts. First, the chirality of B decays has not yet been measured. And second, the mass of the W_R could still be relatively small:[6]

$$M_R^g \equiv \left(\frac{g_L}{g_R}\right) M_R > 300 \text{ GeV} , \tag{1}$$

where M_R is the mass of the W_R, and g_L and g_R are the left and right couplings, respectively.

With this in mind, GW have proposed a model in which the SM W doesn't couple to B's at all. They interpret the long B lifetime as being due to the heaviness of the W_R, *not* to the smallness of V_{cb}. That is,

$$\beta_g \equiv \left(\frac{g_R^2}{g_L^2}\right)\left(\frac{M_L^2}{M_R^2}\right) \sim |V_{cb}| = 0.044 \pm 0.006. \tag{2}$$

Phenomenologically, β_g is bounded to be < 0.07.

In order for this model to be viable, the form of the right-handed CKM matrix V^R must take into account a large number of phenomenological constraints involving B's – the B lifetime, $b \to u$ transitions, 2-body B decays, Cabibbo-suppressed B decays, B_d^0-$\overline{B}_d^0$ mixing – as well as the K_L-K_S mass difference, Δm_K. The forms suggested by GW for both the left- and right-handed CKM matrix, consistent with the above data, are

$$V^L = \begin{pmatrix} \cos\theta_c & \sin\theta_c & 0 \\ -\sin\theta_c & \cos\theta_c & 0 \\ 0 & 0 & 1 \end{pmatrix} , \quad V^R = \begin{pmatrix} c^2 & -cs & s \\ s(1-c)/\sqrt{2} & (c+s^2)\sqrt{2} & c/\sqrt{2} \\ -s(1-c)/\sqrt{2} & -(c-s^2)\sqrt{2} & c/\sqrt{2} \end{pmatrix}, \tag{3}$$

in which θ_c is the Cabibbo angle, and $s \equiv \sin\theta^R$, $c \equiv \cos\theta^R$. The magnitude of s is determined from $|V_{ub}/V_{cb}|$, i.e.

$$s = 0.08 \pm 0.02 \; . \tag{4}$$

With this choice of left- and right-handed CKM matrices, all known data can be explained with

$$M_R^g = 300\text{–}600 \text{ GeV}. \tag{5}$$

In fact, strictly speaking, this is not completely true – additional assumptions are necessary. For example, this model demands the existence of a right-handed neutrino with a mass $m(\nu_R) < m_b - m_c$. Furthermore, if the ν_R is very light, muon decay experiments require either that M_R^g be in the upper part of the range of Eq. 5, or that the ν_R be unstable. In addition, from direct searches for right-handed W's at hadron colliders, the limit $M_R > 520$ GeV is obtained for $g_R = g_L$. Thus, if one wants a value for M_R^g in the lower part of the range of Eq. 5, it is necessary that g_R be larger than g_L. Nevertheless, despite these caveats, the model is interesting in the sense that it points out certain aspects of B decays which must be examined in order to fully test the SM.

One possibly bothersome aspect of the GW solution (Eq. 3), pointed out by Hou and Wyler[7] (HW), is that V_{cd}^R is unnaturally small ($= 0.0003$). One way to avoid this is to parametrize the right-handed CKM matrix using two angles θ_{12} and θ_{13}. HW propose[7]

$$V^R = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1/\sqrt{2} & 1/\sqrt{2} \\ 0 & -1/\sqrt{2} & 1/\sqrt{2} \end{pmatrix} \begin{pmatrix} c_{13} & 0 & s_{13} \\ 0 & 1 & 0 \\ -s_{13} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} c_{12} & -s_{12} & 0 \\ s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \tag{6}$$

and take $s_{12} \simeq 0.098$, $s_{13} \simeq 0.085$, in which case $V_{cd}^R = s_{12} - c_{12}s_{13} \simeq 0.01$. I will refer to this as solution (I).

HW also point out that even if the $b \to c$ transitions are dominated by right-handed currents, $b \to u$ decays might still be mediated mainly via left-handed currents. They thus arrive at solution (II):

$$V^L \simeq \begin{pmatrix} 1 & \lambda & \delta \\ -\lambda & 1 & \delta \\ -\delta & -\delta & 1 \end{pmatrix} , \quad V^R \simeq \begin{pmatrix} 1 & -\varepsilon & \varepsilon \\ \varepsilon & 1/\sqrt{2} & 1/\sqrt{2} \\ -\varepsilon & -1/\sqrt{2} & 1/\sqrt{2} \end{pmatrix}, \tag{7}$$

with $\lambda = \sin\theta_c$, $\delta \sim 0.05$ and $\varepsilon < 0.01$.

Now, the question is, how can one rule out these models? One of the advantages of a hadron collider, as compared to an asymmetric e^+e^- collider operating at the $\Upsilon(4s)$ resonance, is that one can search directly for new physics. It is more likely that physics beyond the standard model – supersymmetry, extra Higgses, technicolour, etc. – will be first found via direct searches than by looking for indirect signals in B physics. As such, the most straightforward way to rule out models of right-handed B decays is simply to look for, and fail to find, a light W_R.

Another possibility[7] is to look at certain B decays which are suppressed in these models relative to the SM. For example,

$$\begin{aligned} \frac{BR(b \to c\bar{c}d)}{BR(b \to c\bar{c}s)} &= \lambda^2 \simeq 0.05, && \text{(SM)}, \\ &= O(10^{-7}), && \text{(GW)}, \\ &\lesssim O(10^{-4}), && \text{(I,II)}. \end{aligned} \tag{8}$$

In this case, if right-handed currents were responsible for B decays, the ratio of the decay rates of $B \to D^{(*)}D_s^{-(*)}$ and $B \to D^{(*)}D^{-(*)}$ would differ from that of the SM. Similarly,

$$\begin{aligned} \frac{BR(b \to c\bar{u}s)}{BR(b \to c\bar{u}d)} &= \lambda^2 \simeq 0.05, && \text{(SM)}, \\ &\simeq 0.008, && \text{(GW,I)}, \\ &\simeq \varepsilon^2 \lesssim O(10^{-4}), && \text{(II)}. \end{aligned} \tag{9}$$

Here one should compare, for example, $\overline{B} \to D^{(*)}\rho$ and $\overline{B} \to D^{(*)}K^*$.

Finally, there is the possibility of measuring the chirality of B decays. The lepton forward-backward decay asymmetry A_{fb} in the decay $\overline{B} \to D^*\ell^-\overline{\nu}_\ell$ is sensitive to the chirality of the $b \to c$ coupling.[8] However, not being parity-violating, A_{fb} also depends on the chirality of the lepton current, and therefore cannot distinguish models of right-handed B decays from the standard model. On the other hand, experiments at LEP *can* make such a distinction. One looks[9] at the reaction $e^+e^- \to Z^0 \to \Lambda_b X$, in which the Λ_b is highly polarized, its spin carried essentially entirely by the b-quark. The electron energy spectrum in $\Lambda_b \to$ *charm* semileptonic decays is then quite sensitive to the $V \pm A$ nature of the $b \to c$ coupling. In this way it might be possible to rule out models of right-handed B decays at LEP.

3. RARE B DECAYS

3.1 $b \to s\gamma$ (and $b \to d\gamma$)

The flavour-changing decay $b \to s\gamma$ occurs first at one loop, and is dominated at lowest order by the t-quark contribution:

$$\mathcal{M}(b \to s\gamma) = \frac{G_F}{\sqrt{2}} \frac{e}{4\pi^2} \lambda_t F_2(x_t) q^\mu \epsilon^\nu \bar{s} \sigma_{\mu\nu} \left(m_b(1+\gamma_5) + m_s(1-\gamma_5)\right) b \, , \tag{10}$$

in which $\lambda_t \equiv V_{tb}V_{ts}^*$, $x_t = m_t^2/M_W^2$, and

$$F_2(x) = \frac{x}{24(x-1)^4} \left[6x(3x-2)\log x - (x-1)(8x^2+5x-7)\right] . \tag{11}$$

However, this process receives important QCD contributions,[10] as shown in Fig. 1. For example, for $m_t = 150$ GeV, we find

$$\begin{aligned} BR(b \to s\gamma) &= 1.4 \times 10^{-4} \quad \text{(no QCD corrections)}, \\ &= 4.2 \times 10^{-4} \quad \text{(including QCD corrections)}. \end{aligned} \tag{12}$$

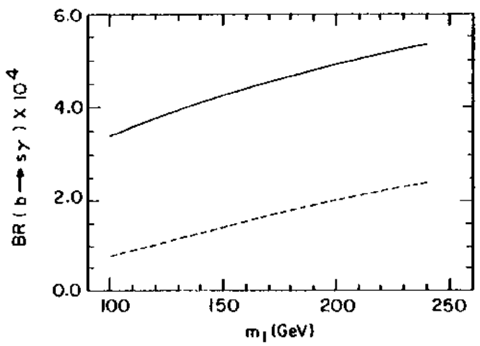

Figure 1: Branching ratio for $b \to s\gamma$ in the SM with (solid line) and without (dashed line) QCD corrections (from Ref. 11 (reproduced by permission)).

The rate for $b \to d\gamma$ is obtained from that for $b \to s\gamma$ (Eq. 10) by replacing the s-quark variables by d-quark variables. Thus, to lowest order,

$$\frac{BR(b \to d\gamma)}{BR(b \to s\gamma)} = \left|\frac{V_{td}}{V_{ts}}\right|^2 . \tag{13}$$

However, there are additional corrections due to the breaking of $SU(3)_{flavour}$. Estimating these, and taking into account the uncertainty in the magnitude of V_{td}, one finds[11]

$$\begin{aligned} BR(b \to s\gamma) &= 3\text{-}5 \times 10^{-4} \, , \\ BR(b \to d\gamma) &= 0.5\text{-}3 \times 10^{-5} \, . \end{aligned} \tag{14}$$

Although the inclusive decay rate for $b \to s\gamma$ can be calculated with good precision, it is well-known that exclusive decays are poorly understood theoretically:

$$R_B \equiv \frac{BR(B \to K^*\gamma)}{BR(b \to s\gamma)} = 4\text{-}40 \ \%. \tag{15}$$

CLEO has measured both inclusive and exclusive flavour-changing decays:[12]

$$\begin{aligned} BR(b \to s\gamma) &< 8.4 \times 10^{-4} \quad (1991), \\ &< 5.4 \times 10^{-4} \quad (1993), \\ BR(B \to K^*\gamma) &= (4.5 \pm 1.5 \pm 0.9) \times 10^{-5} \quad (1993). \end{aligned} \tag{16}$$

These measurements have important consequences for models of new physics.

First consider models with two Higgs doublets (2HDM). In general, such models will lead to flavour-changing neutral currents. This then requires that the Higgs bosons be very heavy, rendering their effects in B physics unobservable. There are two ways to avoid this, distinguished by the couplings of the fermions and the Higgses. One possibility (model I) is that one Higgs doublet, ϕ_2, gives mass to all fermions, while the other doublet, ϕ_1, decouples. In the other case (model II), one doublet, ϕ_2, couples to all u-type quarks, while the second Higgs doublet, ϕ_1, gives mass to d-type quarks. It is model II which appears in supersymmetric and axion models.

In either of these 2HDM there are new contributions to the decay $b \to s\gamma$, found by replacing the $W^\pm$ in the loop by a charged Higgs, $H^\pm$. In these models, both Higgs doublets acquire vacuum expectation values, denoted v_1 and v_2. We define $\tan\beta \equiv v_2/v_1$, which is apriori completely free. The transition amplitude is then proportional to

$$A_W\left(\frac{m_t^2}{M_W^2}\right) + \lambda A_H^1\left(\frac{m_t^2}{M_{H^\pm}^2}\right) + \frac{1}{\tan^2\beta} A_H^2\left(\frac{m_t^2}{M_{H^\pm}^2}\right) , \tag{17}$$

where A_W and $A_H^{1,2}$ represent the SM and charged-Higgs contributions to the amplitude, respectively. In model I, $\lambda = -1/\tan^2\beta$, while $\lambda = +1$ in model II.

From this we see that in model I, there is an enhancement to the rate for $b \to s\gamma$ only for small values of $\tan\beta$. In model II, the rate is also enhanced for small $\tan\beta$. More importantly, due to the A_H^1 term, the rate is *always* larger than that of the SM. This leads to a lower bound on the mass of the charged Higgs in this model,[13,14] independent of the value of m_t. In Fig. 2, taken from Ref. 13, the constraints on models I and II are shown for $m_t = 150$ GeV, using the 1991 CLEO bound (Eq. 16).

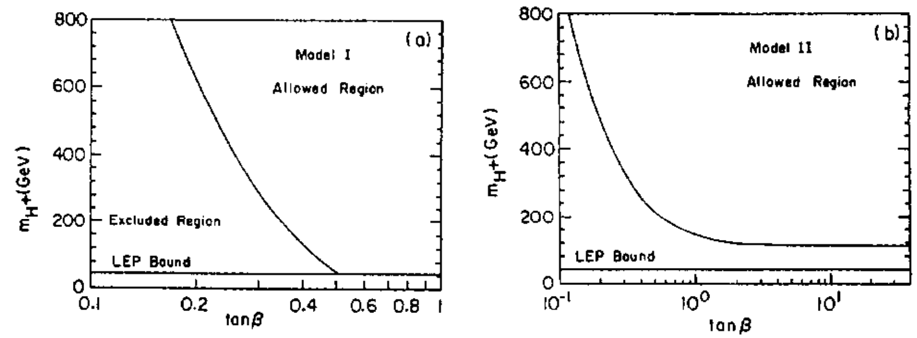

Figure 2: Excluded regions in the $M_{H^\pm}$-$\tan\beta$ plane for models I and II, for $m_t = 150$ GeV, (from Ref. 13 (reproduced by permission)).

For model I, we see that there is no $\tan\beta$-independent lower limit on $M_{H^\pm}$ coming from the bound on $b \to s\gamma$. However, in model II, we find that $M_{H^\pm} > 110$ GeV at large $\tan\beta$, with stronger bounds for smaller values of $\tan\beta$. For model II this lower limit has been updated[15] using the 1993 data on $b \to s\gamma$ (Eq. 16): $M_{H^\pm} > 320$ GeV (540 GeV) for $m_t = 120$ GeV (150 GeV). This new lower bound has several important consequences. First, the decay $t \to bH^+$ is no longer allowed. Second, if the two Higgs doublets are part of a supersymmetric theory, the difficult region for Higgs searches is now ruled out (see below, however). Finally, this eliminates most large effects in 2HDM in other rare B decays.

The implications of the limits on $BR(b \to s\gamma)$ are less clearcut for supersymmetric models. If the main new contributions to $b \to s\gamma$ came from the two Higgs doublets, then the constraints would be as described above. However, the situation is more complicated. First, electroweak radiative corrections to the charged-Higgs mass and to the charged Higgs-fermion-fermion vertex can be substantial.[16] These corrections tend to weaken the constraints on the charged-Higgs mass as a function of $\tan\beta$. More importantly, in the minimal supersymmetric standard model (MSSM), the contributions to $b \to s\gamma$ from other supersymmetric particles may not be negligible.[17] In this case there can be cancellations with the charged-Higgs contributions, possibly resulting in a branching ratio for $b \to s\gamma$ which is *smaller* than that of the SM. Thus, it is impossible to say anything concrete regarding the constraints on SUSY models due to $BR(b \to s\gamma)$.

Finally, left-right symmetric models are essentially unconstrained by the limits on $BR(b \to s\gamma)$ (Eq. 16). Models with right-handed B decays predict a rate for $b \to s\gamma$ which is down by a factor of 2 compared to the SM. And in models with manifest left-right symmetry, the W_R must be so heavy that its effects in $b \to s\gamma$ are negligible.

3.2 $b \to s\ell^+\ell^-$

In the SM, at the quark level, the decay $b \to s\ell^+\ell^-$ arises through penguin diagrams with a virtual γ or Z^0, as well as through box diagrams. In addition, in contrast to $b \to s\gamma$, $b \to s\ell^+\ell^-$ receives important long-distance contributions. These effects are dominated by the decays $B \to \Psi(\Psi')X \to \ell^+\ell^- X$, whose branching ratios have been measured by the ARGUS and CLEO collaborations[18] to be $O(10^{-3})$. The long-distance effects are then very important when the $\ell^+\ell^-$ pair has an invariant mass close to that of the Ψ or Ψ'. However, since the long-distance contribution is so much larger than the short-distance contribution, which is estimated to be $O(10^{-5})$ (see below), one has to worry about residual effects in the spectrum away from the Ψ and Ψ' resonances. In other words, the invariant dilepton mass spectrum is important in analysing $b \to s\ell^+\ell^-$.

The short-distance contributions have been calculated:[19,20,21]

$$\begin{aligned} BR(B \to X_s e^+e^-) &= 0.6\text{–}2.5 \times 10^{-5} , \\ BR(B \to X_s \mu^+\mu^-) &= 3.5\text{–}14.0 \times 10^{-6} , \end{aligned} \quad (18)$$

for 100 GeV $< m_t <$ 200 GeV. Note that the m_t-dependence is much more important here than in $b \to s\gamma$. Also note the the UA1 upper limit:[22]

$$BR(B \to \mu^+\mu^- X) < 5 \times 10^{-5} . \quad (19)$$

The short-distance contributions for the inclusive decays $b \to d\ell^+\ell^-$ have also been computed,[21] assuming $|V_{td}/V_{ts}| = 0.21$:

$$\begin{aligned} BR(B \to X_d e^+e^-) &= 2.6\text{–}10.0 \times 10^{-7} , \\ BR(B \to X_d \mu^+\mu^-) &= 1.5\text{–}6.0 \times 10^{-7} . \end{aligned} \quad (20)$$

Again, it must be remembered that the above cross sections are only the short-distance contributions. One can try to also include the long-distance effects, but there are large uncertainties. Nevertheless it is possible to isolate the short-distance contributions by looking at the forward-backward asymmetry in the decay. In Fig. 3 one sees the angular distribution of the decay, for three different values of m_t, in which θ is defined as the angle between the momentum of the B-meson and that of the ℓ^+ in the centre-of-mass frame of the dilepton pair, and $\hat{s}$ is the scaled dilepton invariant mass. This figure is taken from Ref. 23, to which I refer the reader for more details.

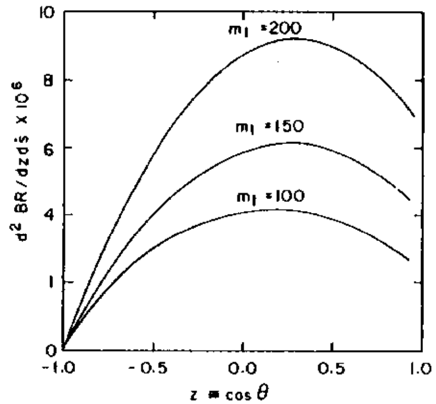

Figure 3: The angular distribution $d^2 BR/dz\,d\hat{s}$ in the decay $b \to s\ell^+\ell^-$, for $\hat{s} = 0.3$ (from Ref. 23 (reproduced by permission)).

As mentioned earlier, the constraints from $b \to s\gamma$ on two-Higgs-doublet models preclude large enhancements to $b \to s\ell^+\ell^-$. As to supersymmetric models, in Ref. 24, it is found that the rate for $b \to s\ell^+\ell^-$ can be greater than that of the SM by up to a factor of 2, when the electroweak symmetry is broken radiatively. On the other hand, this reference predates the recent CLEO bounds on $b \to s\gamma$, and I'm not sure how their inclusion would change the predictions of SUSY models for $b \to s\ell^+\ell^-$. The feeling seems to be that the CLEO data probably now precludes SUSY enhancements to $b \to s\ell^+\ell^-$, but this should be checked.[25]

Another type of new physics which could lead to an enhancement of the rate for $b \to s\ell^+\ell^-$ is extended technicolour. In fact, for certain models, specifically those which include a "techni-GIM" mechanism, the enhancement is too large.[26] In such models, barring delicate fine-tuned cancellations, the prediction for $BR(B \to \mu^+\mu^- X)$ is $O(10^{-4})$, which is in conflict with the UA1 bound (Eq. 19). These models therefore appear to be ruled out. On the other hand, extended technicolour models without a GIM mechanism are still allowed – they predict $BR(B \to \mu^+\mu^- X) = 1\text{–}3 \times 10^{-5}$, an enhancement of roughly a factor of 4 compared to the SM.

3.3 $b \to s\nu\bar{\nu}$

Although the decay $b \to s\nu\bar{\nu}$ has negligible QCD corrections, it is very sensitive to the value of m_t. In the SM, its branching ratio is calculated to be[19,21,27]

$$\sum_i BR(b \to s\bar{\nu}_i\nu_i) = 2.8\text{–}13.0 \times 10^{-5} \quad (21)$$

for 100 GeV $< m_t <$ 200 GeV.

This branching ratio is not expected to be significantly affected by the presence of new physics. In two-Higgs-doublet models, any possible effects are already ruled out by the $b \to s\gamma$ measurement, and the inclusion of supersymmetric particles[24] is not expected to lead to any enhancement.

3.4 $B_s^0 \to \mu^+\mu^- / \tau^+\tau^-$

In order to deduce the form of the operator leading to the decay $B_s^0 \to \ell^+\ell^-$, one notes the following points. First, the s-b matrix element is

$$\langle 0|\bar{s}\gamma^\mu\gamma_5 b|B\rangle = f_B P_B^\mu \ . \qquad (22)$$

This is because the matrix element of $\bar{s}\gamma^\mu b$ vanishes due to considerations of parity (the B_s^0 is a pseudoscalar) and $\bar{s}\sigma^{\mu\nu}b$ won't work since there aren't enough Lorentz vectors to construct a scalar. Second, $P_B^\mu \bar{u}_\ell\gamma_\mu v_\ell = 0$, which means we need a helicity flip in the leptonic current. Thus, the operator describing the decay $B_s^0 \to \ell^+\ell^-$ is

$$\mathcal{O} \sim \bar{s}\gamma^\mu\gamma_5 b\,\bar{\ell}\gamma_\mu\gamma_5\ell \ . \qquad (23)$$

The helicity flip means, of course, that the final answer will depend on the lepton mass.

The branching ratio for the decay $B_s^0 \to \mu^+\mu^-$ is given in Fig. 4 as a function of m_t,[27,28] for $f_{B_s} = 200$ MeV, $\tau_{B_s} = 1.49$ psec, and $|V_{ts}| = 0.042$. For $m_t = 150$ GeV, this gives

$$\begin{aligned} BR(B_s^0 \to \mu^+\mu^-) &= 2\times 10^{-9} \ , \\ BR(B_s^0 \to \tau^+\tau^-) &= 4\times 10^{-7} \ . \end{aligned} \qquad (24)$$

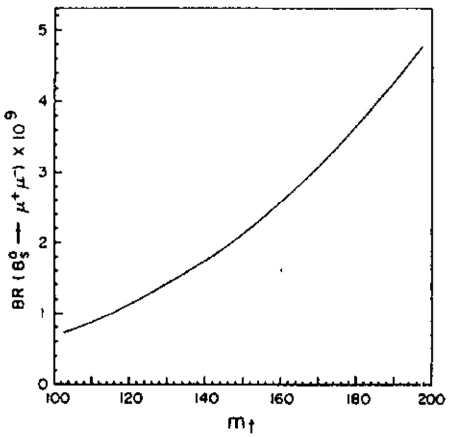

Figure 4: Standard model branching ratio for $B_s^0 \to \mu^+\mu^-$ as a function of m_t, assuming $f_{B_s} = 200$ MeV, $\tau_{B_s} = 1.49$ psec, and $|V_{ts}| = 0.042$.

In two-Higgs-doublet models, there can be an enhancement to the rate for $B_s^0 \to \mu^+\mu^-, \tau^+\tau^-$ by as much as one to two orders of magnitude.[29] Since this decay proceeds through the loop-induced exchange of a neutral Higgs scalar, the constraint on the $M_{H^\pm}$ from $b \to s\gamma$ is unimportant. In extended technicolour models without a GIM mechanism, the rate can also be an order of magnitude bigger than that of the SM.[30] (Recall that extended technicolour models with a GIM mechanism are already in conflict with date from $B \to \mu^+\mu^- X$.) Finally, light leptoquarks could also enhance the rate for $B_s^0 \to \mu^+\mu^-$.

3.5 $B_s \to \gamma\gamma$

All that I will say about this process[28,21] is that in the SM $BR(B_s \to \gamma\gamma) = 1.5\times 10^{-8}$ for $m_t = 150$ GeV and $f_{B_s} = 200$ MeV.

3.6 *Hadronic penguins*

The predictions for the exclusive rates of penguin-induced hadronic B decays are highly model dependent. However, it is important to measure the branching ratios of such decays for several reasons. First, this will give us some idea as to the importance of penguin contributions[31] in CP-violating hadronic B asymmetries. Also, we will be able to test different models of exclusive decays and hence gain some information regarding QCD effects in B decays.

Some examples of such penguin-induced decays[3,32] and their predicted branching ratios (taken from Ref. 32) are given in Table 1. These specific final states have been chosen since the signal consists only of charged particles, so that these processes might be observable at hadron colliders.

Mode	Branching Ratio
$B^+ \to K^0\pi^+$	1.06×10^{-5}
$K^+\phi$	1.12×10^{-5}
$K^{*0}\pi^+$	0.58×10^{-5}
$K^{*+}\phi$	3.12×10^{-5}
$B_d^0 \to K^{*0}\phi$	3.12×10^{-5}
$K^{*0}\rho^0$	0.62×10^{-5}

Table 1: Some exclusive penguin-induced hadronic B decays and their predicted branching ratios (from Ref. 32).

4. B_c PHYSICS

One particularly interesting piece of B physics which is likely to be studied at hadron colliders is the $B_c = (\bar{b}c)$ system (for further discussion regarding B_c physics, see Refs. 33 and 34). The mass of the B_c has been calculated,[34,35] using potential models, to be $\simeq 6.25$ GeV. Its production cross-section is about $\sigma(B_c)/\sigma(b\bar{b}) \sim 10^{-3}$. This leads to[36]

1.3×10^3	B_c's per year (10^7 sec.) at	LEP,
2.0×10^6		TeVatron,
1.1×10^7		LHC (fixed target),
1.1×10^{11}		LHC.

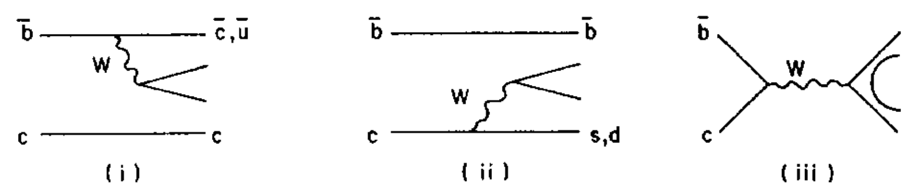

Figure 5: The three mechanisms for B_c decay: (i) c-spectator, (ii) b-spectator, (iii) annihilation.

The main reason that B_c mesons are so interesting is that there are three mechanisms for their decay, shown in Fig. 5. Examples of these different decays are:

$$\begin{aligned} c\text{-spectator:}\quad & B_c^+ \to \Psi e^+ \nu_e \,, \\ & B_c^+ \to \eta_c e^+ \nu_e \,, \\ & B_c^+ \to \Psi \pi^+ \,, \\ & B_c^+ \to D^+ \overline{D^0} \,, && (25) \\ b\text{-spectator:}\quad & B_c^+ \to B_s^{(*)} e^+ \nu_e \,, \\ & B_c^+ \to B^{(*)} e^+ \nu_e \,, \\ & B_c^+ \to B_s^{(*)} \rho^+ \,, \\ & B_c^+ \to B^{(*)+} \overline{K^0} \,, && (26) \\ \text{annihilation:}\quad & B_c^+ \to \tau^+ \nu_\tau \,, \\ & B_c^+ \to D^{(*)+} K^0 \,. && (27) \end{aligned}$$

Note that, unlike B_u's, B_d's and B_s's, the annihilation decays of the B_c are expected to be important. There are a number of reasons for this. First, helicity suppression is ineffective if there are heavy particles (e.g. τ, D, ...) in the final state. Second, in the B_c system, such decays are unsuppressed by CKM factors. And finally, f_{B_c} is expected to be large.

The relative importance of these three different decay mechanisms have been estimated. Using quark and spectator models, and taking $\tau_{B_c} \simeq 5 \times 10^{-13}$ sec., the inclusive branching ratios for each of these three types of decay are predicted to be:[37]

$$\begin{aligned} c\text{-spectator:}\quad & 37\%, \\ b\text{-spectator:}\quad & 45\%, \\ \text{annihilation:}\quad & 18\%. && (28) \end{aligned}$$

Assuming $\tau_{B_c} \simeq 9 \times 10^{-13}$ sec., QCD sum rules give:[38]

$$\begin{aligned} c\text{-spectator:}\quad & 48\%, \\ b\text{-spectator:}\quad & 39\%, \\ \text{annihilation:}\quad & 13\%. && (29) \end{aligned}$$

For a more complete discussion of these relative inclusive branching ratios, see Ref. 34.

There are several particularly interesting decay modes of the B_c which involve a Ψ in the final state. The decay $B_c^+ \to \Psi\pi^+$ is likely to be the discovery mode. Its branching ratio is estimated to be 2×10^{-3} and it permits the full reconstruction of the B_c. The decay $B_c^+ \to \Psi\mu^+\nu_\mu$ has a large branching ratio (1-4 $\times 10^{-2}$) and its signal is three leptons coming from the same vertex. In fact, $BR(B_c \to \Psi + X)$ is estimated to be (19-24)%, which means that the B_c probably can be seen at CDF.

Given a sufficiently large sample of B_c's, it is even possible to look for CP violation in the B_c system.[39] In order to have a non-zero CP-violating decay-rate asymmetry, it is necessary to choose a final state which can be reached via two different weak amplitudes. For example, the decay $B_c^+ \to D^0 K^+$ has two contributions with different CKM matrix elements – a c-spectator tree diagram and a $\bar{b} \to \bar{s}$ penguin diagram. Another example is the processes $B_c^\pm \to D^0 D_s^\pm$ and $B_c^\pm \to \overline{D^0} D_s^\pm$. By measuring these decay rates and the rate for $B_c^\pm \to D_{CP}^0 D_s^\pm$, where D_{CP}^0 is identified by its decay to a CP eigenstate, the angle γ of the unitarity triangle can in principle be extracted.[40] (Unfortunately, this particular example is probably experimentally unfeasible, due to the tiny product branching ratios.)

5. B_s^0-$\overline{B_s^0}$ MIXING

The measurement of B_s^0-$\overline{B_s^0}$ mixing[41] is important for a number of reasons:

- The mixing parameter $x_s \equiv (\Delta M)_{B_s}/\Gamma_{B_s}$ is expected to be large (> 3). If found to be small, this would be a smoking gun for new physics.
- x_s can be used in conjuction with x_d to get a handle on V_{td}:

$$\frac{x_d}{x_s} \sim \frac{f_{B_d}^2 B_{B_d}}{f_{B_s}^2 B_{B_s}} \left| \frac{V_{td}}{V_{ts}} \right|^2 . \qquad (30)$$

  The ratio of hadronic matrix elements is usually known better than each individual one. Thus, the measurement of x_s would enable us to extract $|V_{td}|$ with better precision.
- An accurate knowledge of x_s is needed to extract the CP-violating angle γ in B_s^0 decays.

In the SM, B_s^0-$\overline{B_s^0}$ mixing is dominated by t-quark exchange in the box diagram, leading to

$$x_s = \tau_{B_s} \frac{G_F^2}{6\pi^2} M_W^2 M_{B_s} \left(f_{B_s}^2 B_{B_s} \right) \eta_{B_s} y_t f_2(y_t) |V_{ts}^* V_{tb}|^2 \,, \qquad (31)$$

in which $y_t \equiv m_t^2/M_W^2$ and

$$f_2(x) = \frac{1}{4} + \frac{9}{4}\frac{1}{(1-x)} - \frac{3}{2}\frac{1}{(1-x)^2} - \frac{3}{2}\frac{x^2 \ln x}{(1-x)^3} . \qquad (32)$$

Taking

$$\begin{aligned} |V_{ts}| = |V_{cb}| &= 0.042 \pm 0.005 \,, \\ \tau_{B_s} = \tau_B &= 1.49 \pm 0.04 \text{ psec} \,, \\ \eta_{B_s} = \eta_B &= 0.55 \,, \\ M_{B_s} &= 5.38 \text{ GeV} \,, \end{aligned} \qquad (33)$$

this gives

$$x_s = (175 \pm 21) \frac{f_{B_s}^2 B_{B_s}}{(1 \text{ GeV})^2} y_t f_2(y_t). \qquad (34)$$

For 89 GeV $\leq m_t \leq$ 182 GeV, the function $y_t f_2(y_t)$ is in the range 0.88–2.72, and is equal to 2.03 for the "central" value of m_t, 150 GeV.

A consensus has not yet been reached regarding the value of $f_{B_s}^2 B_{B_s}$. Potential models and QCD sum rules tend to give smaller values, while lattice calculations give larger values. I will therefore consider two ranges for $f_{B_s}^2 B_{B_s}$:

$$\begin{aligned} (I):\quad & f_{B_s}\sqrt{B_{B_s}} = 180 \pm 35 \text{ MeV}, \\ (II):\quad & f_{B_s}\sqrt{B_{B_s}} = 225 \pm 25 \text{ MeV}. \end{aligned} \qquad (35)$$

These lead to the following "central" values for x_s (taking $m_t = 150$ GeV):

$$\begin{aligned} (I): \quad x_s &= 11.5, \\ (II): \quad x_s &= 18.0. \end{aligned} \tag{36}$$

The "1σ" lower limits on x_s are

$$\begin{aligned} (I): &\quad x_s > 3.3, \\ (II): &\quad x_s > 6.6. \end{aligned} \tag{37}$$

Clearly there is a large theoretical uncertainty regarding the hadronic matrix elements. For example, lattice estimates give[42]

$$\begin{aligned} f_{B_d} &= 188\text{–}246 \text{ MeV}, \\ f_{B_s} &= 204\text{–}241 \text{ MeV}. \end{aligned} \tag{38}$$

However, the error on the ratio of these two quantities is considerably smaller:[43]

$$\frac{f_{B_s}^2 B_{B_s}}{f_{B_d}^2 B_{B_d}} = 1.19 \pm 0.10 \ . \tag{39}$$

This is why a precise measurement of x_s can be used, along with x_d, to extract V_{td} (see Eq. 30).

It is possible to get smaller values of x_s if one invokes physics beyond the SM. Examples of such new physics are: a fourth generation,[44] non-minimal SUSY models,[45] fine-tuned left-right symmetric models[46] and models with Z-mediated flavour-changing neutral currents.[47] However, none of these is particularly compelling.

6. T VIOLATION

The last topic I wish to briefly discuss is T violation. By this I do not mean CP violation, which is discussed elsewhere,[48] but rather triple-product correlations. There are two examples of these which have been discussed in the literature, having to do with the decays $B \to V_1 V_2$ (V_1 and V_2 are spin-1 mesons) and $B \to D^* \ell \nu_\ell$. I won't go into very much detail regarding either of these decays, preferring instead to simply sketch out the salient features.

Consider first the decay[49]

$$B(p) \to V_1(k, \epsilon_1) V_2(q, \epsilon_2), \tag{40}$$

in which the particles are specified by their 4-momenta (p, k, q) and their polarizations (ϵ_1, ϵ_2). The most general decay amplitude can be written

$$\mathcal{M} = a\, \epsilon_1 \cdot \epsilon_2 + \frac{b}{m_1 m_2} (p \cdot \epsilon_1)(p \cdot \epsilon_2) + i \frac{c}{m_1 m_2} \epsilon^{\alpha\beta\gamma\delta} \epsilon_{1\alpha}\, \epsilon_{2\beta}\, k_\gamma\, p_\delta \ , \tag{41}$$

in which

$$\begin{aligned} a &= |a| e^{i(\delta_a + \phi_a)} \ , \\ b &= |b| e^{i(\delta_b + \phi_b)} \ , \\ c &= |c| e^{i(\delta_c + \phi_c)} \ , \end{aligned} \tag{42}$$

where $\delta_{a,b,c}$ and $\phi_{a,b,c}$ are the strong phases and the weak phases, respectively. The corresponding amplitude for the decay of the antiparticle is

$$\overline{\mathcal{M}} = \bar{a}\, \epsilon_1 \cdot \epsilon_2 + \frac{\bar{b}}{m_1 m_2} (p \cdot \epsilon_1)(p \cdot \epsilon_2) - i \frac{\bar{c}}{m_1 m_2} \epsilon^{\alpha\beta\gamma\delta} \epsilon_{1\alpha}\, \epsilon_{2\beta}\, k_\gamma\, p_\delta \ , \tag{43}$$

in which $\bar{a}, \bar{b}, \bar{c}$ are identical to a, b, c (Eq. 42), except that the $\phi_{a,b,c}$ change sign.

Now, the asymmetry

$$A_B = \frac{N_{\text{events}}(\vec{k} \cdot \vec{\epsilon}_1 \times \vec{\epsilon}_2 > 0) - N_{\text{events}}(\vec{k} \cdot \vec{\epsilon}_1 \times \vec{\epsilon}_2 < 0)}{N_{TOT}} \tag{44}$$

can be written

$$A_B \sim \text{Im}(ac^*) \sim |ac| \sin(\delta + \phi), \tag{45}$$

where $\delta \equiv \delta_a - \delta_c$ and $\phi \equiv \phi_a - \phi_c$. If we imagine measuring a similar asymmetry $A_{\overline{B}}$ for the antiparticle decay, then we can obtain

$$\begin{aligned} A_B + A_{\overline{B}} &\sim |ac| \cos\delta \sin\phi \ , \\ A_B - A_{\overline{B}} &\sim |ac| \sin\delta \cos\phi \ . \end{aligned} \tag{46}$$

The useful thing about such asymmetries, particularly the sum $A_B + A_{\overline{B}}$, is that they are sensitive to the weak phases only, i.e. they do not vanish if $\delta = 0$. On the other hand, the question of how to relate phases at the meson level to phases at the quark level, and of how to calculate strong phases, introduces much theoretical uncertainty and model dependence.[50] Still, the signals would be interesting to look for. Some possible decay modes are: $\overline{B_d^0} \to \rho^{*+} K^{*-}$, $B^- \to \Psi K^{*-}$ and $\overline{B_s^0} \to \Psi \phi$.

Another interesting process is the decay $B \to D^* \ell \nu_\ell$, in which the D^* decays further to $D\pi$.[51] The triple product $\vec{p}_\ell \cdot (\vec{p}_{D^*} \times \vec{p}_D)$ is T-violating. There are a variety of asymmetries one can measure which depend on this triple product (I refer the reader to Ref. 51 for more details). Again, to go from the quark-level calculation to the meson-level measurement introduces hadronic uncertainties and model dependence. However, this triple product vanishes in the SM, so that this would be another way of looking for CP violation from new physics.

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