LEFT-RIGHT SYMMETRY AND *CP* VIOLATION IN THE B SYSTEM

David LONDON [1]
Deutsches Elektronen Synchrotron - DESY, D-2000 Hamburg, FRG

and

Daniel WYLER
Institut für Theoretische Physik, Universität Zürich, CH-8001 Zurich, Switzerland

Received 16 August 1989

We have reanalyzed a version of LR-symmetric models in which the right-handed W-boson could be light. Excluding several fine-tunings, it is shown that, due to constraints from ϵ, $M_{W_R} \gtrsim 5$ TeV, which would render its effects in heavy quark systems small. If fine-tuning is allowed, there are interesting effects in the B system – in particular, B_s^0–$\bar{B}_s^0$ mixing can be small, the *CP* asymmetry in $B_d \to \psi K_S$ may be negative, and there is *CP* violation in Cabibbo allowed B_s decays.

Although the observed *CP* violation can be incorporated into the standard model in a satisfactory way, new forces might be responsible for this effect. *CP* violation has been considered in a variety of extensions of the standard model such as supersymmetry [1,2], extended Higgs models [1,3], and left–right (LR) symmetric theories [1,4]. In LR symmetric theories, the right-handed Cabibbo–Kobayashi–Maskawa (CKM) matrix, V^R, is usually taken to be equal to the usual left-handed CKM matrix, V^L. Due to the K_L–K_S mass difference, this then implies [5] a large mass M_R for the right-handed charged intermediate vector boson, W_R, whose effects are then suppressed.

Larger effects are possible if M_R is small, or, more precisely, if $\beta_R \equiv (g_R^2/g_L^2)(M_L^2/M_R^2)$ is as large as possible. Here g_L and g_R are the left- and right-handed gauge couplings, respectively, and M_L is the mass of the ordinary W-boson. The effect of the W_R might then be important in B-meson mixings and decays. Furthermore, because of the large number of phases in V^R, *CP* violation could be sizeable in processes where the standard model gives very little contribution, for instance in Cabibbo allowed B_s decays such as $B_s \to \psi\phi$.

In order to remove the constraints on β_R, the restriction $V^R = V^L$ must be relaxed. The right-handed contributions to kaon weak processes are of the form $C_K^R \beta_R$, where C_K^R depends on the elements of the first two columns of V^R (i.e. those elements which involve the d- and s-quarks). The corresponding expressions for mesons with heavy quarks (D, B) are of the same form, with C_K^R replaced by C_D^R (elements involving u- and c-quarks) or C_B^R (d- (or s-) and b-quarks). Kaon physics puts bounds only on $C_K^R \beta_R$; in order to maximize the effects in heavy quark systems, one must therefore choose C_K^R as small as possible.

It is well known that K^0–$\bar{K}^0$ mixing puts strong constraints on β_R [5]. These come from limiting the contributions to ΔM_K of the box diagrams in fig. 1 to the experimental value. These graphs are proportional to $m_i m_j V_{id}^{R*} V_{js}^R$, where i, j=u, c, t are the intermediate quarks. The contributions are clearly small if either $V_{id}^R \simeq 0$ or $V_{js}^R \simeq 0$; i, j=c, t. This condition immediately yields the following two forms for V^R:

[1] Address after October 1, 1989: Nuclear Physics Laboratory, Université de Montréal, Montreal, Quebec, Canada H3C 3J7.

0370-2693/89/$ 03.50 © Elsevier Science Publishers B.V. (North-Holland)

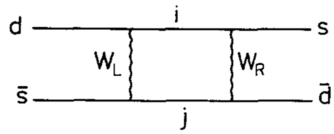

Fig. 1. W_L–W_R contribution to K^0–$\bar{K}^0$ mixing. The internal quarks are i, j = u, c, t. There is also a diagram in which the internal bosons and quarks are interchanged.

$$\text{Model I:} \quad V^R = \begin{pmatrix} e^{i\alpha} & * & * \\ * & c e^{i\sigma} & -s e^{i\gamma} \\ * & s e^{i\theta} & c e^{i\epsilon} \end{pmatrix}, \tag{1}$$

$$\text{Model II:} \quad V^R = \begin{pmatrix} * & e^{i\alpha} & * \\ c e^{i\sigma} & * & -s e^{i\gamma} \\ s e^{i\theta} & * & c e^{i\epsilon} \end{pmatrix}. \tag{2}$$

Here $c \equiv \cos\phi$, $s \equiv \sin\phi$, and the $*$'s denote terms which are $\ll 1$. (The values of ϕ, α, σ, γ, θ, and ϵ need not be the same in the two models, but we use the same symbols for convenience.) Unitarity requires

$$\sigma - \theta = \gamma - \epsilon . \tag{3}$$

The long B-lifetime implies $|\beta_R s| < 0.06$ [#1]. In addition, in Model II, the doubly Cabibbo suppressed charm decays $c \to d u \bar{s}$ yields $|c\beta_R| \leqslant 0.05$, if they are not observed above the prediction of the standard model [#2]; the strongest limit on such decays of 0.4% for $D^+ \to K^+\pi^+\pi^-$ [6] gives $|c\beta_R| \leqslant 0.06$. Thus, in Model II, $\beta_R < \frac{1}{16}$. B_d^0–$\bar{B}_d^0$ mixing [7] will constrain β_R even further.

These plausible results have been confirmed recently in a careful study by Langacker and Sankar [8]. They consider four different possibilities for the right-handed neutrino: heavy Majorana, heavy Dirac, intermediate mass (10–100 MeV), and very light. Constraints on β_R are obtained in each case using experimental information from ΔM_K, B_d^0–$\bar{B}_d^0$ mixing, B decays, neutrinoless double beta decay, universality, nonleptonic kaon decays, muon decay and astrophysics. They find that the weakest limits on β_R occur when the right-handed neutrino is heavy and Dirac. In this case, when reasonable restrictions on fine-tuning are imposed (for example, the $*$'s in eqs. (1) and (2) are allowed to be as large as 10^{-2}) they obtain

$$\text{Model I:} \quad \beta_R \lesssim 0.04 , \quad \frac{g_L M_R}{g_R} \gtrsim 400\ \text{GeV} ,$$

$$\text{Model II:} \quad \beta_R \lesssim 0.013 , \quad \frac{g_L M_R}{g_R} \gtrsim 700\ \text{GeV} . \tag{4}$$

(The smaller value of β_R in Model II is mainly due to the requirement that the LR contribution to B_d^0–$\bar{B}_d^0$ mixing be no larger than that of the standard model.) However, they do note that if extreme fine-tuning is allowed, the W_R could still be as light as the W_L.

The aim of this paper is to further investigate Models I and II and their possible influence on mixing and CP violation in the B system. We will use the limits on β_R given in eq. (4). In addition, the elements V_{ij}^R, denoted by a $*$ in eqs. (1), (2), will be written as

$$V_{ij}^R = 10^{-2} X_{ij} \exp(i\varphi_{ij}) . \tag{5}$$

We begin with the analysis of ϵ in the K system. Writing

$$\epsilon = \frac{1}{\sqrt{2}} \frac{\text{Im}\, M_{12}}{\Delta M_K} , \tag{6}$$

the contribution of the left–right box diagram of fig. 1 to ϵ is [#3]

$$\epsilon_{LR} = (1.2 \times 10^6)\beta_R \sum_{i,j=u,c,t} \sqrt{x_i x_j}\, \eta_{ij} \times \text{Im}(V_{id}^{L*} V_{is}^R V_{jd}^{R*} V_{js}^L) I(x_i, x_j, \beta) , \tag{7}$$

where $x_i = m_i^2/M_L^2$ and $\beta = M_L^2/M_R^2$, i.e. $\beta_R = (g_R^2/g_L^2)\beta$. In deriving eq. (7), we have used the vacuum saturation approximation ($B_K^{LR} = 1$) and the usual enhancement of 7.7 for the LR matrix element [5]. The η_{ij} are the short-distance QCD corrections, which we will set equal to 1 in the following. (Since, in general, $\eta_{ij} \gtrsim 1$ [9], their inclusion would even strengthen our results.) The function I is given by [10]

$$I(x_i, x_j, \beta) = (1 + \tfrac{1}{4} x_i x_j \beta) I_1(x_i, x_j, \beta) - \tfrac{1}{4}(1+\beta) I_2(x_i, x_j, \beta) , \tag{8}$$

#1 We consider only hadronic B decays since, in order to avoid bounds from muon decay and all mesonic semi-leptonic decays, we assume $m_{\nu_R} > M_B$.

#2 We thank J.-M. Gérard for pointing this out.

#3 In view of the stringent limits we will disregard the mixing between W_L and W_R.

where

$$I_1(x_i,x_j,\beta)=\frac{x_i\ln x_i}{(1-x_i)(1-x_i\beta)(x_i-x_j)}+(i\leftrightarrow j)$$
$$-\frac{\beta\ln\beta}{(1-\beta)(1-x_i\beta)(1-x_j\beta)},$$
$$I_2(x_i,x_j,\beta)=\frac{x_i^2\ln x_i}{(1-x_i)(1-x_i\beta)(x_i-x_j)}+(i\leftrightarrow j)$$
$$-\frac{\ln\beta}{(1-\beta)(1-x_i\beta)(1-x_j\beta)}. \quad (9)$$

We will set $\beta_R\simeq\beta$ in the following; thus both β and β_R are limited by eq. (4). We note that, because $\beta\ll 1$, I depends on β essentially only through the last term ($\ln\beta$). This β dependence will be kept explicitly in the following. Furthermore, in those graphs in which a t-quark is involved, the explicit dependence of ϵ_{LR} on m_t is linear (ut, ct) or quadratic (tt). Thus the weakest limits on M_R come for small values of m_t. We will therefore take $m_t=50$ GeV (which is the same value used in ref. [9]), corresponding roughly to its lower bound.

Using the Wolfenstein parametrization of the left-handed CKM matrix [11]

$$V^L=\begin{pmatrix}1-\frac{1}{2}\lambda^2 & \lambda & A\rho\lambda^3e^{-i\delta}\\ -\lambda & 1-\frac{1}{2}\lambda^2 & A\lambda^2\\ A\lambda^3(1-\rho e^{i\delta}) & -A\lambda^2 & 1\end{pmatrix}, \quad (10)$$

where $\lambda=0.22$, $A\simeq 1.0$, and $\rho\leqslant 0.9$ [12], we obtain the following contributions to ϵ_{LR} in Model I:

$$\text{cu: } 0.026(\beta/0.04)c\sin(\sigma-\alpha)$$
$$\times[1-0.029\ln(\beta/0.04)], \quad (11)$$
$$\text{tu: } 0.01(\beta/0.04)s[\sin(\theta-\alpha)-\rho\sin(\theta-\alpha-\delta)]$$
$$\times[1-0.12\ln(\beta/0.04)],$$
$$\text{cc: } 0.28(\beta/0.04)X_{cd}c\sin(\sigma-\varphi_{cd})$$
$$\times[1-0.032\ln(\beta/0.04)],$$
$$\text{ct: } 0.13(\beta/0.04)X_{td}c\sin(\sigma-\varphi_{td})$$
$$\times[1-0.12\ln(\beta/0.04)],$$
$$\text{tc: } 0.13(\beta/0.04)X_{cd}s$$
$$\times[\sin(\theta-\varphi_{cd})-\rho\sin(\theta-\varphi_{cd}-\delta)]$$
$$\times[1-0.12\ln(\beta/0.04)],$$
$$\text{tt: } 0.14(\beta/0.04)X_{td}s$$
$$\times[\sin(\theta-\varphi_{td})-\rho\sin(\theta-\varphi_{td}-\delta)]$$
$$\times[1-0.17\ln(\beta/0.04)], \quad \text{(11 cont'd)}$$

where we have taken $m_u=5.6$ MeV, $m_c=1.5$ GeV, and $m_t=50$ GeV. In eq. (11), the notation q_1q_2 indicates that $i=q_1$ and $j=q_2$ in eq. (7). The contributions from other combinations of quarks in the loop are much smaller. In Model II the relevant contributions are

$$\text{uc: } 0.18(\beta/0.013)c\sin(\alpha-\sigma)$$
$$\times[1-0.028\ln(\beta/0.013)],$$
$$\text{ut: } 0.08(\beta/0.013)s\sin(\alpha-\theta)$$
$$\times[1-0.1\ln(\beta/0.013)],$$
$$\text{cc: } 0.09(\beta/0.013)X_{cs}c\sin(\varphi_{cs}-\sigma)$$
$$\times[1-0.031\ln(\beta/0.013)],$$
$$\text{ct: } 0.05(\beta/0.013)X_{cs}s\sin(\varphi_{cs}-\theta)$$
$$\times[1-0.1\ln(\beta/0.013)],$$
$$\text{tc: } 0.05(\beta/0.013)X_{ts}c$$
$$\times[\sin(\varphi_{ts}-\sigma)-\rho\sin(\varphi_{ts}-\sigma-\delta)]$$
$$\times[1-0.1\ln(\beta/0.013)],$$
$$\text{tt: } 0.05(\beta/0.013)X_{ts}s$$
$$\times[\sin(\varphi_{ts}-\theta)-\rho\sin(\varphi_{ts}-\theta-\delta)]$$
$$\times[1-0.14\ln(\beta/0.013)]. \quad (12)$$

We now require ϵ_{LR} to be less than the experimental value of 2.28×10^{-3}. This can be achieved in various ways.

(1) No fine-tuning of the elements (including phases) of V^R. In this case, which is clearly the most natural, β must be sufficiently small so that each individual contribution to ϵ_{LR} satisfies the above bound. We obtain, from eqs. (11), (12), as particular examples

$$\text{cc: } \beta<2.81\times10^{-4}, \quad M_R>4.8\text{ TeV},$$
$$\text{tt: } \beta<3.62\times10^{-4}, \quad M_R>4.3\text{ TeV}, \quad (13)$$

for Model I, and

$$\text{uc}: \beta < 1.46\times10^{-4}\,, \quad M_R > 6.7\ \text{TeV}\,,$$
$$\text{cc}: \beta < 2.94\times10^{-4}\,, \quad M_R > 4.7\ \text{TeV}\,,$$
$$\text{tt}: \beta < 3.98\times10^{-4}\,, \quad M_R > 4.1\ \text{TeV}\,, \tag{14}$$

for Model II. This result strengthens dramatically the lower limit on M_R which was thought to be avoided by eqs. (1), (2). It can be easily understood. The bounds in eq. (4) come about (in part) by requiring that the contribution from the right-handed sector to ΔM_K be no larger than that of the standard model. However, in this model, the W_R contributes equally to $\text{Re}(M_{12})$ and $\text{Im}(M_{12})$. Therefore the bound on β decreases by a factor ~ 500.

(2) Fine-tunings of the V^R_{ij}. There are several ways to do this. (We discuss here only Model I; Model II is similar.) (a) Cancellations among all the terms in eq. (11). (b) Adjustment of the *CP*-violating phases in V^R so that each term in (11) vanishes. For example, we can set $\sigma\simeq\alpha\simeq\varphi_{cd}\simeq\varphi_{td}$. The two pieces of the tu, tc, ct and tt terms can then either vanish separately or cancel each other. In the former case we have $\theta\simeq\alpha$, $\delta\simeq0$, and, from eq. (3), $\gamma\simeq\epsilon$. This would imply that all *CP* violation resides in the right-handed sector, a rather strange situation. (c) We can set the $X_{ij}\ll 1$. Then the cu and tu contributions imply that the phases must be (O(10%)), which is a relatively weak assumption. (This fine-tuning is clearly not possible in Model II, since the largest contribution, uc, is not proportional to an X_{ij}.) (d) Note that setting $s\simeq 0$ or $c\simeq 0$ would still require a small β or one of the above fine-tunings.

Unlike ϵ, ϵ' does not much constrain the right-handed sector. The s→d penguin graph in the standard model is proportional to $\lambda^5 \sin\delta$ (either c- or t-quark exchange), while the right-handed contributions are $\sim 10^{-2}\beta c \sin\sigma$ (c-quark) and $\sim 10^{-2}\beta s \times \sin\theta$ (t-quark). The ratio is therefore roughly given by $10^{-2}\beta/\lambda^5 \sim 1$. Only if $\sin\delta\simeq 0$ are there any constraints on the V^R_{ij}; otherwise no useful information can be extracted.

We now turn to the effects of the W_R in B-meson physics (there are no sizeable contributions to charm physics). As we have seen, without fine-tuning V^R, due to the bounds on M_R (eqs. (13), (14)), the W_R plays no significant role. We therefore consider the effects when fine-tuned solutions (and a light W_R) are allowed.

Probably the most interesting method for seeing *CP* violation in the B system is through hadronic decay asymmetries [13,14]. Because of mixing a state which starts out as a pure B^0 (or $\bar{B}^0$) will evolve in time into a mixture of B^0 and $\bar{B}^0$. *CP* violation is then manifested by a nonzero value of the asymmetry

$$A_f = \frac{\Gamma(B^0(t)\to f) - \Gamma(\bar{B}^0(t)\to \bar{f})}{\Gamma(B^0(t)\to f) + \Gamma(\bar{B}^0(t)\to \bar{f})}\,, \tag{15}$$

where f is a hadronic final state into which both B^0 and $\bar{B}^0$ can decay, and $\bar{f}$ is its *CP* conjugate. If eq. (15) is integrated over time, then the asymmetry takes the form

$$A_f = -\frac{2x\,\text{Im}\,\lambda_f}{2+x^2+x^2|\rho_f|^2}\,. \tag{16}$$

Here, x is the B^0–$\bar{B}^0$ mixing parameter ($x=\Delta M/\Gamma$),

$$\rho_f = \frac{A(\bar{B}^0\to f)}{A(B^0\to f)}\,, \tag{17}$$

and

$$\lambda_f = \frac{q}{p}\rho_f\,, \tag{18}$$

where

$$\frac{q}{p} = \sqrt{\frac{M^*_{12}-\frac{1}{2}i\Gamma^*_{12}}{M_{12}-\frac{1}{2}i\Gamma_{12}}}\,. \tag{19}$$

In the standard model, $\Gamma_{12}\ll M_{12}$ [15], so that q/p is a pure phase. If the final state is not a *CP* eigenstate (e.g. $B_d\to D^+\pi^-, B_s\to D_s^+\pi^-$), then A_f depends on the hadron dynamics in the ρ_f term, which makes reliable calculation of the asymmetry very difficult. The most useful final states are those which are *CP* eigenstates, in which case $|\rho_f|=1$, and A_f takes the familiar form [#4]

$$A_f = -\frac{x}{1+x^2}\,\text{Im}\,\lambda_f\,. \tag{20}$$

In these cases, the standard model makes some definite predictions [14,16]. For Cabibbo allowed B_d decays, $A(B_d\to\psi K_S)$ is expected to be positive. In addition, $A(B_d\to\psi K_S)$ should be equal to $A(B_d\to\psi\pi^0)$. In the case of Cabibbo allowed B_s decays, there is a particularly strong prediction. Both $A(B_s\to\psi\phi)$

[#4] If it is possible to do time-dependent measurements, then $\text{Im}\,\lambda_f$ may be measured directly.

and $A(B_s\to\psi K_S)$ are expected to be $\simeq 0$! (For Cabibbo suppressed B_d and B_s decays, the standard model allows all values of the asymmetries, so that by themselves, with regards to finding new physics, these processes are somewhat less interesting.)

How might these predictions be influenced by the presence of a light W_R? Let us first examine Model I. In this case the possibly relevant four-fermi interactions are

$$b\to c\bar{c}s\,,\quad b\to c\bar{u}d \tag{21}$$

(the transition $b\to u$ is suppressed in both models). In Model I, because $V^R_{td}\simeq 0$, B^0_d–$\bar{B}^0_d$ mixing is not affected by the presence of the W_R. Thus, apart from the option $\delta\simeq 0$ discussed above, the quantity q/p for B_d's is unchanged. (We have checked that, even in the presence of right-handed interactions, the relation $\Gamma_{12}\ll M_{12}$ still holds.) However, the amplitudes for the decays in eq. (21) are now modified. For the decays $b\to c\bar{c}s$ and $\bar{b}\to\bar{c}c\bar{s}$ (which correspond to the interesting decay $B_d\to\psi K_S$), they are

$$A(b\to c\bar{c}s)\sim\mathcal{M}_L\lambda^2+\mathcal{M}_R\beta cs\,e^{i(\gamma-\sigma)}\,,$$

$$A(\bar{b}\to c\bar{c}\bar{s})\sim\mathcal{M}_L\lambda^2+\mathcal{M}_R\beta cs\,e^{i(\gamma-\sigma)}\,, \tag{22}$$

where $\mathcal{M}_L$ and $\mathcal{M}_R$ are the relevant matrix elements of the left- and right-handed operators [#5]. The effect of eq. (22) is to modify the asymmetry in the decay $B_d\to\psi K_S$ in an interesting way. In the standard model, $\rho_f=1$ for this process. In Model I, setting $\mathcal{M}_L=\mathcal{M}_R$, we have

$$\rho_f=\frac{1-(\beta/\lambda^2)sc\,e^{i(\gamma-\sigma)}}{1-(\beta/\lambda^2)sc\,e^{-i(\gamma-\sigma)}}\,. \tag{23}$$

ρ_f is still a pure phase, and for $\beta_{sc}\gtrsim\lambda^2$ this phase can be sizeable. (Note that the phase γ is free – it is unaffected by any of the previous fine-tunings.) Therefore $A(B_d\to\psi K_S)$ may be significantly altered by the presence of the W_R; it can even be negative. Furthermore, the asymmetry $A(B_d\to\psi\pi^0)$ is unaffected (since the W_R does not participate in the quark decay $b\to c\bar{c}d$). Thus one may have $A(B_d\to\psi K_S)\neq A(B_d\to\psi\pi^0)$ in this model.

[#5] One might worry that there is no interference between $\mathcal{M}_L$ and $\mathcal{M}_R$ since, on the quark level, with massless quarks, amplitudes involving left- and right-handed quarks do not interfere. This does not apply on the meson level, as can be easily seen from the fact that the ψ is a coherent sum of left- and right-handed quarks.

The case of B_s-mesons is particularly interesting. In addition to the absence of *CP* violation in Cabibbo allowed B_s decays, the standard model also predicts a large value for the B^0_s–$\bar{B}^0_s$ mixing parameter, x_s [17], which would suppress the time-integrated asymmetries in eq. (20). Using eq. (7) (adapted for the B system), one gets for the off diagonal matrix element of the B_s mass matrix

$$M_{12}=M^{st}_{12}[1-9.3s^2\,e^{i(\gamma-\theta)}-10.0cs\,e^{i(\epsilon-\theta)}]\,, \tag{24}$$

where M^{st}_{12} is the standard model contribution. The dominant left–right contributions are the ct and tt intermediate states and $\beta=0.04$, $m_t=50$ GeV and $m_c=1.5$ GeV were used (relative to M^{st}_{12}, the LR contributions decrease as m_t increases). Surprisingly, $x_s=2|M_{12}|/\Gamma$ can be considerably smaller than in the standard model if the terms in eq. (24) cancel. (Small values of x_s are usually thought to be a sign of four generation models [18].) Nonzero *CP* asymmetries in Cabibbo allowed B_s decays are now possible. Not only is ρ_f in the decay $B_s\to\psi\phi$ complex (eq. (24)), but q/p for the B_s system (eqs. (19), (24)) is also complex (both are essentially equal to 1 in the standard model). In addition, as in the case of B_d decays, in general one would expect $A(B_s\to\psi\phi)\neq A(B_s\to\psi K_S)$ since $\rho_f=1$ for the latter process.

Due to the quark decay $b\to c\bar{u}d$, the W_R in Model I will also affect processes such as $B_d\to D^+\pi^-$ and $B_s\to D_s^+\pi^-$. However, as mentioned above, since these final states are not *CP* eigenstates, the values of the *CP* asymmetries are very uncertain in the standard model. Thus, even if *CP* violation were seen here, it would be very difficult to ascertain whether or not a light W_R played any role.

In Model II, the possible quark decays are

$$b\to c\bar{c}d\,,\quad b\to c\bar{u}s\,. \tag{25}$$

The analysis of Model II is similar to that of Model I, although, since M_R is larger here, the effects are somewhat smaller. First of all, in this case, the phase q/p for B_d's may be changed [#6], but for the B_s system the W_R does not affect q/p (it remains equal to 1). Secondly, ρ_f retains its standard model value ($=1$)

[#6] Of course, the elements of V^R must be chosen so that the observed value of x_d is reproduced.

for the decays $B_d \to \psi K_S$ and $B_s \to \psi\phi$, but, due to eq. (25), may be complex for the processes $B_d \to \psi\pi^0$ and $B_s \to \psi K_S$. Thus $A(B_d \to \psi K_S)$ may be different from its standard model value, and, as in Model I, $A(B_d \to \psi K_S) \neq A(B_d \to \psi\pi^0)$ in general. Unlike Model I, Model II still gives $A(B_s \to \psi\phi) \simeq 0$, although $A(B_s \to \psi K_S)$ may be nonzero.

In conclusion, we have shown that, in the absence of fine-tuning, it is not possible to have a light right-handed W. Although one can evade the constraints from ΔM_K by taking $V^R \neq V^L$ and choosing a particular form for V^R, constraints from ϵ will, in general, require $M_R \gtrsim 5$ TeV. If fine-tuned solutions are allowed, then not only can the W_R be light, but there may be interesting effects in the B system. In one model, B_s^0–$\bar{B}_s^0$ mixing can be smaller than B_d^0–$\bar{B}_d^0$ mixing. Also, the pattern of *CP* asymmetries predicted by the standard model may be altered. The asymmetry in Cabibbo allowed B_s decays (e.g. $B_s \to \psi\phi$) may be nonzero; $A(B_d \to \psi K_S)$ may be negative; and the predicted relations $A(B_d \to \psi K_S) = A(B_d \to \psi\pi^0)$ and $A(B_s \to \psi\phi) = A(B_s \to \psi K_S)$ may be violated.

We thank Francesca Borzumati, Jean-Marc Gérard and Paul Langacker for helpful discussions. D.L. also thanks Paul Langacker for the hospitality of the University of Pennsylvania, where part of this work was done.

References

[1] For a review, see L. Wolfenstein, Ann. Rev. Nucl. Part. Sci. 36 (1986) 137, and references therein.
[2] S. Bertolini, F. Borzumati and A. Masiero, Phys. Lett. B 194 (1987) 545; Phys. Lett. B 198 (1987) 590(E).
[3] J.F. Donoghue and E. Golowich, Phys. Rev. D 37 (1988) 2542.
[4] G. Ecker and W. Grimus, Z. Phys. C 30 (1986) 293.
[5] G. Beall, M. Bander and A. Soni, Phys. Rev. Lett. 48 (1982) 848.
[6] Particle Data Group, G.P. Yost et al., Review of particle properties, Phys. Lett. B 204 (1988) 1.
[7] ARGUS Collab., H. Albrecht et al., Phys. Lett. B 192 (1987) 245;
CLEO Collab., A. Jawahery, Proc. XXIV Intern. Conf. on High energy physics (Munich, FRG, 1988).
[8] P. Langacker and S.U. Sankar, University of Pennsylvania preprint Print-89-0255 (PENN) (1989).
[9] G. Ecker and W. Grimus, Nucl. Phys. B 258 (1985) 328.
[10] G. Ecker and W. Grimus, Phys. Lett. B 153 (1985) 279.
[11] L. Wolfenstein, Phys. Rev. Lett. 51 (1983) 1945.
[12] For a recent analysis, see C.S. Kim and A.D. Martin, Phys. Lett. B 225 (1989) 186.
[13] A. Carter and A.I. Sanda, Phys. Rev. Lett. 45 (1980) 952; Phys. Rev. D 23 (1981) 1567;
I.I. Bigi and A.I. Sanda, Nucl. Phys. B 193 (1981) 85; B 281 (1987) 41;
I. Dunietz and J. Rosner, Phys. Rev. D 34 (1986) 1404;
Y. Azimov, V. Khoze and M. Uraltsev, Yad. Fiz. 45 (1987) 1412;
D. Du, I. Dunietz and D. Wu, Phys. Rev. D 34 (1986) 3414.
[14] P. Krawczyk, D. London, R.D. Peccei and H. Steger, Nucl. Phys. B 307 (1986) 19.
[15] A.J. Buras, W. Slominski and H. Steger, Nucl. Phys. B 238 (1984) 529; B 245 (1984) 369;
J. Hagelin, Nucl. Phys. B 193 (1981) 123.
[16] D. London, DESY preprint DESY-89-048 (1989).
[17] G. Altarelli and P. Franzini, Z. Phys. C 37 (1988) 271.
[18] J. Hewett and T. Rizzo, Mod. Phys. Lett. A 3 (1988) 975.