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Phenomenological analysis of a topless left-right model

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We reanalyze a topless left-right model due to Ma, after allowing flavor-changing neutral currents which are consistent with all the experimental limits coming from flavor-changing processes. The resulting model can naturally reproduce the observed values of B_d^0-$\bar{B}_d^0$ mixing, where the original model had failed by 8 orders of magnitude. Thus, the model is viable from the viewpoint of low-energy phenomenology. In the case of Z^0 physics, however, we see that the model makes distinctive predictions for some of the $Z^0 \rightarrow f\bar{f}$ partial widths due to the nonstandard couplings of the fermion f. The recent measurements of these partial widths from CERN LEP seem to clearly rule out this model. Indeed the precision measurement of the $Z^0 \rightarrow b\bar{b}$ partial width along with the known properties of b decay seems to rule out any topless model.

I. INTRODUCTION

The top quark is one of the particles predicted by the standard model[1] (SM) that is not yet experimentally observed.[2] Theoretically, the top quark is needed in the SM for the cancellation of chiral anomalies.[3] Experimentally, there are three pieces of data which indicate the existence of an SU(2) partner of the b quark, i.e., the top quark. These are (1) the upper limit on the flavor-changing decay[4,5] $b \rightarrow \mu^+\mu^- X$, (2) the sign and the magnitude of the forward-backward asymmetry[6,7] in $e^+e^- \rightarrow b\bar{b}$, and (3) the observed value of the B_d^0-$\bar{B}_d^0$ mixing.[8,9]

There exist indirect upper and lower limits on the mass of the top quark in the context of the SM.

(1) There is a lower bound of 50 GeV coming from B_d^0-$\bar{B}_d^0$ mixing by equating the experimental value[8] to the result of the SM box diagram with two W exchange.[10]

(2) The ρ parameter of the SM, defined by

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2\theta_W} , \tag{1}$$

is equal to 1 at the tree level. If weak isospin is broken due to large differences in the mass of an isodoublet, 1-loop corrections to the ρ parameter change its value from 1. The upper limit on $\Delta\rho = 1-\rho \leq 0.01$ gives an upper limit on the top-quark mass of 200 GeV.[11] Recent precision measurements of M_Z, Γ_Z, and various partial widths of the Z^0 boson constrain the mass of the top quark to be[12] $m_t \leq 170$ GeV.

Experimentally there are direct production limits $m_t > 77$ GeV from Fermilab Tevatron and ACOL (Antiproton Collector) (Ref. 13) and $m_t > 45$ GeV from CERN LEP.[14]

All the above evidence in favor of the existence of the top quark is from phenomena related to b quark couplings in the SM. That is, they show that, in an SU(2)×U(1) electroweak model, b_L, the left-handed component of the b quark, transforms as an SU(2)-doublet member. Hence it must have a charge $\frac{2}{3}$ quark as its SU(2) partner. In the absence of direct evidence for the top quark, it is pertinent to ask if one can construct a topless model with nonstandard b couplings, which would be consistent with all the above-mentioned phenomena involving the b quark. One such model was proposed by Ma,[15] which is derived from an E_6 grand unified theory (GUT). The fermions in the GUT are grouped into two 27-plets of E_6. Thus this model has two u-type quarks and four d-type quarks. The low-energy interactions in this model are described by an $SU(3)_c \times SU(2)_1 \times SU(2)_2 \times U(1)_H$ gauge group. Under this gauge group, (u_L, d_L) and (c_L, s_L) transform as $SU(2)_1$ doublets and b_L as a singlet. As for the right-handed quarks, (u_R, d_R) form an $SU(2)_2$ doublet as in the case of left-handed quarks. A second $SU(2)_2$ doublet is formed by (c_R, b_R) while s_R is left as an $SU(2)_2$ singlet. The left- and the right-handed components of the fourth d-type quark, which is assumed to be heavy, are singlets under both the SU(2)'s. This model is distinct from the manifest left-right model,[16] because not all the left-handed fermions in this model are doublets under $SU(2)_1$ and singlets under $SU(2)_2$ and vice versa for the right-handed fermions. In order to make this distinction clear the two SU(2)'s are labeled by subscripts 1 and 2 rather than L and R.

This model contains two charged gauge bosons: W_1 (mass M_1) and W_2 (mass M_2) belonging to $SU(2)_1$ and $SU(2)_2$, respectively. W_1 should be identified with the experimentally observed charged gauge boson ($M_1 = 81$ GeV) and is the same as the W boson of the SM. There are also two massive neutral gauge bosons. The lighter one is to be identified with the experimentally observed neutral gauge boson Z^0. The heavier one, denoted by Z_H (mass M_H), is presumed to be very massive with

$M_H \sim M_2$.

All charged-current decays of the b quark in this model are assumed to take place via virtual W_2. Thus the long lifetime of the b quark is interpreted as being due to the large W_2 mass rather than small $b \rightarrow c$ coupling as in the SM. Equating the experimental value of the b-decay width to the prediction of this model gives an upper limit on the W_2 mass[15,17] of about $M_2 \leq 433$ GeV. This limit is compatible with the lower limit on the right-handed gauge boson coming from polarized muon decay[18] $M_2 \geq 405$ GeV and direct production from Tevatron[19,20] $M_2 \geq 380$ GeV.

In the charge $-\frac{1}{3}$ quark sector, some of the quarks belong to SU(2) doublets while the others are SU(2) singlets. In the most general case all the left-handed quarks can mix with one another as can all the right-handed quarks. But mixing between SU(2) singlets and doublet members generates tree-level flavor-changing neutral currents (FCNC's) and there are stringent experimental constraints on these. Ma avoided tree-level FCNC's by assuming that only the quarks with identical gauge quantum numbers mix with each other. Hence the experimental constraints from FCNC processes are trivially satisfied.

Before the advent of LEP, this model has had two problems[21] in explaining the phenomenology of the b quark: i.e., (i) forward-backward asymmetry in $e^+e^- \rightarrow b\bar{b}$ measured at DESY PETRA;[6] (ii) B_d^0-$\bar{B}_d^0$ mixing observed by the ARGUS and CLEO Collaborations.[8] Note that both phenomena above are cited as experimental evidence for the top quark. Because both b_L and b_R are $\mathrm{SU}(2)_1$ singlets, the axial-vector coupling a_b of the b quark to Z^0 is 0 and hence the forward-backward asymmetry, which arises due to the interference of the γ-Z^0 exchange and is proportional to a_b, vanishes. However, Ma has argued that[22] one can obtain the correct sign and appropriate magnitude for this asymmetry if one considers $e^+e^- \rightarrow b\bar{b}$ via the exchange of scalars in a supersymmetric version of this model. On the other hand, the B_d^0-$\bar{B}_d^0$ mixing in this model is about eight orders of magnitude smaller than the experimental value as pointed out by Webb.[17] The reason for this is twofold. First, all the charged-current interactions of the b quark take place via the heavier W_2 boson, and Δm_{B_d} calculated from the box diagram with two W_2 exchange is suppressed by a factor of $(M_1/M_2)^4$. Secondly, the B_d^0-$\bar{B}_d^0$ mixing is proportional to the squared charm-quark mass because there is no top quark in this model. The model has been ruled out by Webb on the basis of this discrepancy. The failure of the model has been attributed to the absence of the top quark rather than the nonstandard couplings of the b quark since the former is directly responsible for the smallness of B_d^0-$\bar{B}_d^0$ mixing.

In this paper we analyze this model allowing flavor-changing neutral currents, which seems more natural to us theoretically. We find that all the experimental constraints on FCNC's can be satisfied with very little fine-tuning and that the FCNC's can generate a B_d^0-$\bar{B}_d^0$ mixing commensurate with the experimental measurement. Moreover, B_s^0-$\bar{B}_s^0$ mixing in this model is not constrained to be large as it is in the SM. Thus this model is consistent with all the low-energy electroweak phenomena. Coming to Z^0 physics, however, we find that the predictions of this model for Z^0 widths are significantly different from those of the SM, mainly due to the fact that b_L is an $\mathrm{SU}(2)_1$ singlet. And the recent measurements at LEP seem to clearly rule out any such nonstandard couplings for the b quark. Thus we conclude that, while Ma's model has no difficulty in describing the low-energy electroweak phenomena, it seems to be ruled out by the recent measurements of Z^0 widths at LEP. Moreover, the reason for its failure is not the absence of the top quark itself, but the gauge coupling of b quark.

II. THE MODEL

The model proposed by Ma is based on an E_6 GUT. All the known fermions are grouped into two generations (two 27-plets of E_6) in this model.[23] The electroweak interactions are described by a gauge model based on the group $\mathrm{SU}(2)_1 \times \mathrm{SU}(2)_2 \times \mathrm{U}(1)_H$. $\mathrm{SU}(2)_1$ is the same as the SU(2) of the SM and the new hypercharge H is related to the SM hypercharge Y by

$$Y = 2T_3^{(2)} + H \ . \tag{2}$$

The fermion content of this model, along with the low-energy gauge quantum numbers, is

$$(\nu_e, e)_L, (\nu_\mu, \mu)_L : (2, 1, -1) \ , \tag{3}$$

$$(\nu_e, e)_R, (\nu_\mu, \mu)_R : (1, 2, -1) \ , \tag{4}$$

$$\begin{bmatrix} \nu_\tau & \tau^c \\ \tau & N_\tau^c \end{bmatrix}_L, \begin{bmatrix} \nu_E & E^c \\ E & N_E^c \end{bmatrix}_L : (2, 2, 0) \ , \tag{5}$$

$$n_L, n'_L : (1, 1, 0) \ , \tag{6}$$

$$(u'_L, d'_L), (c'_L, s'_L) : (2, 1, \tfrac{1}{3}) \ , \tag{7}$$

$$(u'_R, d'_R), (c'_R, b'_R) : (1, 2, \tfrac{1}{3}) \ , \tag{8}$$

$$b'_L, h_L : (1, 1, -\tfrac{2}{3}) \ , \tag{9}$$

$$s'_R, h_R : (1, 1, -\tfrac{2}{3}) \ . \tag{10}$$

The primes on the quarks in (7)–(10) indicate that the quarks are eigenstates of the gauge group, which mix among themselves to form mass eigenstates. Note that in addition to the known particles this model also contains a charge $-\frac{1}{3}$, heavy quark h, a heavy charged lepton E, and seven extra neutral Weyl fermions. Three of these have $\mathrm{SU}(2)_1$ quantum number $\frac{1}{2}$, which means that they couple to Z^0 with the same strength as the known light neutrinos. This fact is important when considering the total decay width and the invisible partial decay width of Z^0.

Among the observed fermions, the τ lepton and the b quark have nonstandard couplings. Both τ_L and τ_R are $\mathrm{SU}(2)_1$-doublet members. Hence, the $Z^0\bar{\tau}\tau$ coupling is purely vector in nature and the axial-vector coupling a_τ of τ lepton to Z^0 vanishes. The axial-vector coupling of b quark to Z^0, a_b, also vanishes because both b_L and b_R

are $SU(2)_1$ singlets and $Z^0\bar{b}b$ coupling is pure vector. The forward-backward asymmetry in $e^+e^- \to \tau^+\tau^-$ and $e^+e^- \to b\bar{b}$ arising from γ-Z^0 interference is zero in this model because the asymmetry is proportional to the axial-vector coupling of the fermion. Ma has shown,[22] however, that it is possible to obtain the correct sign and magnitude for the forward-backward asymmetry in $e^+e^- \to \tau^+\tau^-, b\bar{b}$ due to scalar exchange in a supersymmetric version of this model. In Refs. 15 and 22, it was assumed that only those quarks with identical gauge quantum numbers mix with each other. Thus, in the left-handed sector, only d_L and s_L mix with each other and in the right-handed sector, only d_R and b_R mix with each other. FCNC's at the tree level are avoided because there is no singlet-doublet mixing. Webb calculated B_d^0-$\bar{B}_d^0$ mixing with this assumption and found that it is eight orders of magnitude too small compared to the experimental value.[17] But we find the assumption, that there is no singlet-doublet mixing, to be very constraining and unnatural. The mass matrix of the charge $-\frac{1}{3}$ quarks must have a very specific form for this assumption to be true. For a general mass matrix, singlet-doublet mixing does occur and as a result the model contains tree-level FCNC's.[2] These must, of course, satisfy the relevant experimental constraints. We show below that, with very little fine-tuning, these tree-level FCNC's can produce a value of Δm_{B_d} commensurate with the experimental value while obeying all the constraints.

The charged-current interactions in this model are given by

$$\mathcal{L}_{\rm CC} = \frac{g}{\sqrt{2}} (J_{1\mu}^+ W_1^{\mu-} + J_{1\mu}^- W_1^{\mu+} + J_{2\mu}^+ W_2^{\mu-} + J_{2\mu}^- W_2^{\mu+}) \tag{11}$$

and the neutral-current interactions by

$$\mathcal{L}_{\rm NC} = \frac{g}{\cos\theta} \left[(J_{1\mu}^3 - \sin^2\theta J_\mu^{\rm em}) Z^{0\mu} + \sqrt{\cos 2\theta} \left[J_{2\mu}^3 + \frac{\sin^2\theta}{2\cos 2\theta} J_\mu^{\rm em} \right] Z_H^\mu \right] . \tag{12}$$

The strength of Z_H coupling to fermions depends on the mass matrix of the neutral bosons. For definiteness, we have taken the strengths of neutral-current interactions of this model to be the same as those given in Ref. 24 for the manifest left-right model. The angle θ is the analog of the weak mixing angle θ_W of SM and is defined by

$$\sin^2\theta = \frac{g'^2}{g^2 + 2g'^2} , \tag{13}$$

where g is the coupling strength for both the SU(2)'s and g' is the coupling strength of $U(1)_H$. Experimentally θ is equal to θ_W because it appears as the coefficient of the electromagnetic current in $Z^0\bar{f}f$ couplings. It also satisfies the relation $M_Z^2 = M_1^2/\cos^2\theta$.

The various currents introduced in (11) and (12) are defined as

$$J_\mu^{\rm em} = \sum_{q=u,c,d,s,b} Q_q (\bar{q}'_L \gamma_\mu q'_L + \bar{q}'_R \gamma_\mu q'_R) , \tag{14}$$

$$J_{1\mu}^- = \bar{u}'_L \gamma_\mu d'_L + \bar{c}'_L \gamma_\mu s'_L , \tag{15}$$

$$J_{2\mu}^- = \bar{u}'_R \gamma_\mu d'_R + \bar{c}'_R \gamma_\mu b'_R , \tag{16}$$

$$J_{1\mu}^3 = \tfrac{1}{2} (\bar{u}'_L \gamma_\mu u'_L + \bar{c}'_L \gamma_\mu c'_L - \bar{d}'_L \gamma_\mu d'_L - \bar{s}'_L \gamma_\mu s'_L) , \tag{17}$$

and

$$J_{2\mu}^3 = \tfrac{1}{2} (\bar{u}'_R \gamma_\mu u'_R + \bar{c}'_R \gamma_\mu c'_R - \bar{d}'_R \gamma_\mu d'_R - \bar{b}'_R \gamma_\mu b'_R) . \tag{18}$$

The gauge eigenstates must be transformed into mass eigenstates in order to get the interactions of the physical states. Without loss of generality, we can assume that the mass eigenstates are identical to the gauge eigenstates for the charge $\frac{2}{3}$ quarks. We can then write the gauge eigenstates (primed states) of the charge $\frac{1}{3}$ quarks in terms of the mass eigenstates (unprimed ones) as

$$d'_L = U_{ud}^L d_L + U_{us}^L s_L + U_{ub}^L b_L , \tag{19}$$

$$s'_L = U_{cd}^L d_L + U_{cs}^L s_L + U_{cb}^L b_L , \tag{20}$$

plus the orthogonal combination for b'_L and

$$d'_R = U_{ud}^R d_R + U_{us}^R s_R + U_{ub}^R b_R , \tag{21}$$

$$b'_R = U_{cd}^R d_R + U_{cs}^R s_R + U_{cb}^R b_R , \tag{22}$$

plus the orthogonal combination for s'_R, where U^L and U^R are 3×3 unitary matrices.

In terms of the mass eigenstates the charged currents become

$$J_{1\mu}^- = \bar{u}_L \gamma_\mu (U_{ud}^L d_L + U_{us}^L s_L + U_{ub}^L b_L) + \bar{c}_L \gamma_\mu (U_{cd}^L d_L + U_{cs}^L s_L + U_{cb}^L b_L) \tag{23}$$

and

$$J_{2\mu}^- = \bar{u}_R \gamma_\mu (U_{ud}^R d_R + U_{us}^R s_R + U_{ub}^R b_R) + \bar{c}_R \gamma_\mu (U_{cd}^R d_R + U_{cs}^R s_R + U_{cb}^R b_R) . \tag{24}$$

From (23) we see that the first two rows of U^L describe the strengths of various left-handed charged currents coupling to W_1. They have the same physical significance as the first two rows of the Kobayashi-Maskawa (KM) matrix V of the SM.[25] The important difference between V and U^L is that the elements of the third row of the KM matrix V (V_{td}, V_{ts}, V_{tb}) describe the strengths of the charged currents $t \xrightarrow[W]{} d,s,b$, respectively, while the elements of the third row of U^L have no physical significance because b'_L does not couple to any charge $\frac{2}{3}$ quark. They can, of course, be expressed in terms of the elements of the first two rows using the unitarity relations between U_{ij}^L's. U^R is the right-handed analog of U^L and the elements of its first and third row describe the strengths of various right-handed currents as written in (24). The second row of U^R has no physical significance because s'_R does not couple to any up-type quark.

The matrix elements U_{ud}^L and U_{us}^L are determined to a very good accuracy from β decay and charged kaon decay. The values are equal to the corresponding elements of the KM matrix, $U^L_{ud}=V_{ud}=0.974$ and $U^L_{us}=V_{us}=0.22$. From the charm decays into nonstrange and strange particles one obtains $U^L_{cd}\simeq -0.22$ and $U^L_{cs}\simeq 0.97$. The lifetime of b quark gives the constraint

$$[2|U^L_{ub}|^2+|U^L_{cb}|^2+\beta^2(2|U^R_{ub}|^2+|U^R_{cb}|^2)]^{1/2}\simeq 0.035\ , \tag{25}$$

where $\beta=M_1^2/M_2^2$ and the factor of 2 comes from the relative phase-space factor between the decay channels. U^L_{ub} and U^L_{cb} arise due to the mixing between SU(2)$_1$-singlet b'_L and -doublet members d'_L and s'_L, respectively. This singlet-doublet mixing should be very small to satisfy the limits on the tree-level FCNC's. As we shall see below, these limits imply $U^L_{ub}, U^L_{cb}\sim 10^{-3}$. Then $\Gamma(b\xrightarrow{W_1}c,uX)$ is very small and the first two terms on the left-hand side of (25) can be neglected. The smallness of the charmless b decays implies that $|U^R_{ub}|^2 \ll |U^R_{cb}|^2$ and (25) reduces to

$$\beta|U^R_{cb}|\simeq 0.035\ . \tag{26}$$

The searches for right-handed currents in polarized μ^+ decay[18] give the lower limit $M_2\geq 405$ GeV or $\beta\leq 0.04$. For $|U^R_{cb}|\simeq 1$, (26) implies $\beta=0.035$, which is consistent with the upper limit from muon decay and is the smallest value β can take. Note that the mass of W_2 in this model has to be within a very narrow range of

$$405\text{ GeV}\leq M_2\leq 433\text{ GeV} \tag{27}$$

to satisfy the constraints from b decay and muon decay. The ratio of b-decay rates into noncharm to charm final states $(\Gamma\to u)/(\Gamma\to c)\leq 0.04$ implies that $U^R_{ub}\leq 0.14$. The elements U^R_{us} and U^R_{cs} arise due to mixing between right-handed SU(2)$_2$-singlet and -doublet quarks. We will see in the next section that these also have to be small $(\sim 10^{-3})$ to satisfy the experimental limits on flavor-changing processes. The unitarity of U^R determines $U^R_{ud}\simeq 1$ and $U^R_{cd}\simeq -U^R_{ub}$.

In the right-handed sector one could have made the choice of pairing s'_R into an SU(2)$_2$ doublet with u_R, leaving d'_R as a singlet. This choice gives $U^R_{us}\simeq 1$ and $U^R_{ud}\leq 10^{-3}$. This is acceptable as far as the flavor-changing phenomena are concerned but runs into problems in describing the exclusive modes of the b decays. For this choice the charm-strange exclusive channel would dominate over the charm-nonstrange channel, whereas experimentally, the situation is just reverse.[26] This justifies the choice of pairing (u_R,d'_R) in an SU(2)$_2$ doublet.

Here we are interested in the neutral-current interactions of only the charge $-\frac{1}{3}$ quarks. The electromagnetic current $J^{\rm em}_\mu$ is diagonal in all the three flavors, and it remains flavor diagonal when the gauge eigenstates are transformed into mass eigenstates. On the other hand, the SU(2) neutral currents $J^3_{1\mu}$ and $J^3_{2\mu}$ each have a flavor term "missing." When these currents are rewritten in terms of the mass eigenstates, they each have a flavor-diagonal piece and a flavor-changing piece. The flavor-diagonal parts are

$$\begin{aligned}J^3_{1\mu}(\text{fd})=-\tfrac{1}{2}[&\bar d_L\gamma_\mu(|U^L_{ud}|^2+|U^L_{cd}|^2)d_L\\&+\bar s_L\gamma_\mu(|U^L_{us}|^2+|U^L_{cs}|^2)s_L\\&+\bar b_L\gamma_\mu(|U^L_{ub}|^2+|U^L_{cb}|^2)b_L]\ ,\end{aligned} \tag{28}$$

$$\begin{aligned}J^3_{2\mu}(\text{fd})=-\tfrac{1}{2}[&\bar d_R\gamma_\mu(|U^R_{ud}|^2+|U^R_{cd}|^2)d_R\\&+\bar s_R\gamma_\mu(|U^R_{us}|^2+|U^R_{cs}|^2)s_R\\&+\bar b_R\gamma_\mu(|U^R_{ub}|^2+|U^R_{cb}|^2)b_R]\end{aligned} \tag{29}$$

and the flavor-changing parts are

$$\begin{aligned}J^3_{1\mu}(\text{fc})=-\tfrac{1}{2}[&\bar d_L\gamma_\mu(U^{l*}_{ud}U^L_{us}+U^{L*}_{cd}U^L_{cs})s_L\\&+\bar d_L\gamma_\mu(U^{L*}_{ud}U^L_{ub}+U^{L*}_{cd}U^L_{cb})b_L\\&+\bar s_L\gamma_\mu(U^{L*}_{us}U^L_{ub}+U^{L*}_{cs}U^L_{cb})b_L]+\text{H.c.}\ ,\end{aligned} \tag{30}$$

$$\begin{aligned}J^3_{2\mu}(\text{fc})=-\tfrac{1}{2}[&\bar d_R\gamma_\mu(U^{R*}_{ud}U^R_{us}+U^{R*}_{cd}U^R_{cs})s_R\\&+\bar d_R\gamma_\mu(U^{R*}_{ud}U^R_{ub}+U^{R*}_{cd}U^R_{cb})b_R\\&+\bar s_R\gamma_\mu(U^{R*}_{us}U^R_{ub}+U^{R*}_{cs}U^R_{cb})b_R]+\text{H.c.}\end{aligned} \tag{31}$$

Substituting the expressions of the flavor-diagonal and flavor-changing currents from (28)–(31) in the neutral-current interaction Lagrangian (12), we find the couplings of fermions to Z^0 to be

$$\begin{aligned}\mathcal{L}_{Z^0}&=\frac{g}{\cos\theta}\{[J^3_{1\mu}(\text{fd})-\sin^2\theta J^{\rm em}_\mu]+[J^3_{1\mu}(\text{fc})]\}Z^{0\mu}\\&=[J_{Z\mu}(\text{fd})+J_{Z\mu}(\text{fc})]Z^{0\mu}\end{aligned} \tag{32}$$

and those to Z_H to be

$$\begin{aligned}\mathcal{L}_{Z_H}&=\frac{g\sqrt{\cos 2\theta}}{\cos\theta}\left[\left[J^3_{2\mu}(\text{fd})+\frac{\sin^2\theta}{2\cos 2\theta}J^{\rm em}_\mu\right]+[J^3_{2\mu}(\text{fc})]\right]Z^\mu_H\\&=[J_{H\mu}(\text{fd})+J_{H\mu}(\text{fc})]Z^\mu_H\ .\end{aligned} \tag{33}$$

III. FCNC PHENOMENOLOGY

The flavor-nondiagonal couplings of Z^0 and Z_H to quarks, given in (32) and (33), can give rise to flavor-changing transitions at the tree level. Here, we calculate these tree-level FCNC contributions to Δm_K, $B(b\to\mu^+\mu^- X)$, Δm_{B_d}, and Δm_{B_s}, and compare them to the experimental values.

A. Δm_K

Tree-level FCNC interactions typically give rise to large values for the K_L-K_S mass difference. Hence the measured small value gives stringent bounds on the FCNC couplings and on the masses of the particles mediating the FCNC transitions. In the present case, Δm_K

coming from tree diagrams with Z^0 exchange is

$$\Delta m_K|_{\mathrm{tree}(Z^0)} = \frac{\sqrt{2}G_F f_K^2 m_K}{6} \mathrm{Re}(U_{ud}^{L*}U_{us}^L + U_{cd}^{L*}U_{cs}^L)^2 . \tag{34}$$

Demanding that this should be less than the experimental value[27] yields the constraint

$$(U_{ud}^{L*}U_{us}^L + U_{cd}^{L*}U_{cs}^L)^2 \le 10^{-7} . \tag{35}$$

That is, the Glashow-Iliopoulos-Maiani (GIM) cancellation[28] should take place to one part in 10^4. This is a rather delicate fine-tuning but is consistent with the present experimental values. In fact, for $|U_{ub}^L|, |U_{cb}^L| \simeq 10^{-3}$, the unitarity of the matrix U^L implies that

$$U_{ud}^{L*}U_{us}^L + U_{cd}^{L*}U_{cs}^L \le 10^{-5} , \tag{36}$$

which makes the contribution of tree diagram with Z^0 exchange to Δm_K negligible.

From the tree diagrams with Z_H exchange, we have

$$\Delta m_K|_{\mathrm{tree}(Z_H)} = \frac{\sqrt{2}G_F f_K^2 m_K}{6} \cos 2\theta \frac{M_Z^2}{M_H^2} \times \mathrm{Re}(U_{ud}^{R*}U_{us}^R + U_{cd}^{R*}U_{cs}^R)^2 . \tag{37}$$

We take $M_Z^2/M_H^2 \simeq M_1^2/M_2^2 \simeq 0.035$ to get an order-of-magnitude estimate. Again, we demand that the value given by (37) should be less than the experimental value. Then we get the condition

$$(U_{ud}^{R*}U_{us}^R + U_{cd}^{R*}U_{cs}^R)^2 \le 10^{-6} . \tag{38}$$

The above condition is satisfied if $U_{us}^R \sim 10^{-3} \sim U_{cs}^R$. One may attribute the smallness of U_{us}^R, U_{cs}^R as well as that of U_{ub}^L, U_{cb}^L to the fact that they correspond to mixing of an SU(2)-doublet member with a singlet. The smallness of U_{us}^R and U_{cs}^R guarantees that the Δm_K from the W_1-W_2 box diagrams is negligibly small. Thus the strong bounds obtained previously on M_2 from Δm_K in manifest left-right models[29] are not relevant in this model.

B. $B(b \to \mu^+\mu^- X)$

Experimentally the branching ratio for this process is measured to be $B(b \to \mu^+\mu^- X) \le 10^{-3}$. In this model, this decay can occur at the tree level because both Z^0 and Z_H have flavor-conserving as well as flavor-changing couplings to fermions. The rate for $b \to \mu^+\mu^- X$ can be calculated from the couplings in (32) for the diagram with Z^0 exchange. Comparing this rate with the rate for the dominant b-decay mode $b \xrightarrow{W_2} c\mu^-\bar{\nu}$, we obtain

$$\frac{\Gamma(b \xrightarrow{Z^0} \mu^+\mu^- X)}{\Gamma(b \xrightarrow{W_2} c\mu^-\bar{\nu})} \simeq \left[\frac{M_2^2}{M_1^2}\right]^2 \frac{g_{Z\mu}^2}{|U_{cb}^R|^2 P_C} \times (|U_{ud}^{L*}U_{ub}^L + U_{cs}^{L*}U_{cb}^L|^2 + |U_{us}^{L*}U_{ub}^L + U_{cs}^{L*}U_{cb}^L|^2) , \tag{39}$$

where $g_{Z\mu}^2 = 0.126$ is the sum of the squares of left and right chiral couplings of $\bar{\mu}\mu$ to Z^0 and $P_c \simeq 0.5$ is the phase-space factor for the $(b \to c\mu^-\bar{\nu})$ decay.[30] The branching ratio for $(b \xrightarrow{W_2} c\mu^-\bar{\nu})$ is about 0.12 (Ref. 31). Therefore, the experimental limit on $B(b \to \mu^+\mu^- X)$ is satisfied if the ratio in (39) is less than 0.01. This gives us the constraint

$$|U_{ud}^{L*}U_{ub}^L + U_{cd}^{L*}U_{cb}^L|^2 + |U_{us}^{L*}U_{ub}^L + U_{cs}^{L*}U_{cb}^L|^2 \le 10^{-5} , \tag{40}$$

which is automatically satisfied if

$$U_{ub}^L, U_{cb}^L \le 3 \times 10^{-3} . \tag{41}$$

It was mentioned above that these singlet-doublet mixing elements should be of the order of 10^{-3} to satisfy the limits on FCNC's. Note that such small values imply small decay rate for b via W_1 and guarantee an almost exact GIM cancellation between $U_{ud}^{L*}U_{us}^L$ and $U_{cd}^{L*}U_{cs}^L$ in (36).

From the couplings in (33), we get the rate for $b \xrightarrow{Z_H} \mu^+\mu^- X$ to be

$$\frac{\Gamma(b \xrightarrow{Z_H} \mu^+\mu^- X)}{\Gamma(b \xrightarrow{W_2} c\mu^-\bar{\nu})} \simeq \left[\frac{\cos 2\theta}{\cos^2\theta} \frac{M_2^2}{M_H^2}\right]^2 \frac{g_{H\mu}^2}{|U_{cb}^R|^2 P_C} \times (|U_{ud}^{R*}U_{ub}^R + U_{cd}^{R*}U_{cb}^R|^2 + |U_{us}^{R*}U_{ub}^R + U_{cs}^{R*}U_{cb}^R|^2) , \tag{42}$$

where $g_{H\mu}^2 \simeq 1$ is the sum of squares of the left and right chiral couplings of $\bar{\mu}\mu$ to Z_H. We saw that the constraints from Δm_K force $U_{us}^R, U_{cs}^R \le 10^{-3}$. Hence the second term in the final set of parentheses in (42) is of the order of 10^{-6}. The unitarity of U^R together with the smallness of U_{us}^R and U_{cs}^R implies

$$U_{ud}^{R*}U_{ub}^R + U_{cd}^{R*}U_{cb}^R \le 10^{-5} , \tag{43}$$

as in the case of (36). The first term in the final set of parentheses in (42) then becomes $\sim 10^{-10}$. Therefore the constraint on $b \xrightarrow{Z_H} \mu^+\mu^- X$ is trivially satisfied.

C. Δm_{B_D} and Δm_{B_S}

The main reason we have allowed flavor-nondiagonal couplings in this model is that they can lead to appreciable B_d^0-$\bar{B}_d^0$ mixing.[9] As mentioned earlier,[17] if the FCNC were forbidden, B_d^0-$\bar{B}_d^0$ mixing in this model would be eight orders of magnitude too small. However, now both Z^0- and Z_H-exchange tree diagrams can give rise to Δm_{B_d}. We shall see below that the resulting B_d^0-$\bar{B}_d^0$ mixing is consistent with the experimental value[8] of $x_d \simeq 0.7$. Moreover, in the standard model the unitarity of the three-generational Kobayashi-Maskawa (KM) matrix constrains Δm_{B_s} to be 5–20 times larger than Δm_{B_d}. Hence the B_s^0-$\bar{B}_s^0$ mixing in the SM is maximal, a conclusion that is in potential conflict with the result of the Mark II experiment,[10,32] which suggests $x_s \le 2$ for

$x_d \simeq 0.7$. This can be avoided in this model because the unitarity constraints relating Δm_{B_d} and Δm_{B_s} are not very stringent. It is possible to choose values of U^L_{ub} and U^L_{cb} such that the B^0_d-$\bar{B}^0_d$ mixing is commensurate with the experimental value and the B^0_s-$\bar{B}^0_s$ mixing is small enough to avoid conflict with the Mark II result.

From the Z^0 exchange diagrams we get

$$\Delta m_{B_d}|_{\mathrm{tree}(Z^0)} = \frac{\sqrt{2}G_F f_B^2 m_B}{6} |U^{L*}_{ud}U^L_{ub} + U^{L*}_{cd}U^L_{cb}|^2 , \quad (44)$$

and

$$\Delta m_{B_s}|_{\mathrm{tree}(Z^0)} = \frac{\sqrt{2}G_F f_B^2 m_B}{6} |U^{L*}_{us}U^L_{ub} + U^{L*}_{cs2}U^L_{cb}|^2 . \quad (45)$$

We can reproduce $x_d = \Delta m_{B_d}/\Gamma_b \simeq 0.7$ and $x_s = \Delta m_{B_s}/\Gamma_b \simeq 2$ if

$$U^{L*}_{ud}U^L_{ub} + U^{L*}_{cd}U^L_{cb} \simeq 1\times 10^{-3} , \quad (46)$$

$$U^{L*}_{us}U^L_{ub} + U^{L*}_{cs}U^L_{cb} \simeq \sqrt{3}\times 10^{-3} . \quad (47)$$

Equations (46) and (47) give us the solutions

$$U^L_{ub} \simeq 1.3\times 10^{-3} , \quad (48)$$

$$U^L_{cb} \simeq 1.5\times 10^{-3} , \quad (49)$$

which are consistent with the constraints on these singlet-doublet mixing elements obtained earlier.

The Z_H-exchange diagrams give

$$\Delta m_{B_d}|_{\mathrm{tree}(Z_H)} = \frac{\sqrt{2}G_F f_B^2 m_B}{6}\cos 2\theta \times \frac{M_Z^2}{M_H^2} |U^{R*}_{ud}U^R_{ub} + U^{R*}_{cd}U^R_{cb}|^2 \quad (50)$$

and

$$\Delta m_{B_s}|_{\mathrm{tree}(Z_H)} = \frac{\sqrt{2}G_F f_B^2 m_B}{6}\cos 2\theta \times \frac{M_Z^2}{M_H^2} |U^{R*}_{us}U^R_{ub} + U^{R*}_{cs}U^R_{cb}|^2 . \quad (51)$$

Substituting the limit from (43) on $U^{R*}_{ud}U^R_{ub} + U^{R*}_{cd}U^R_{cb}$ in (50), we see that the Z_H-exchange diagram gives a negligible contribution to Δm_{B_d}. The contribution of (51) to Δm_{B_s} is also negligible because $U^R_{us}, U^R_{cs} \le 10^{-3}$. Thus, the Z^0 exchange tree diagrams can naturally account for the experimental values of B^0_d-$\bar{B}^0_d$ and B^0_s-$\bar{B}^0_s$ mixing, while those from Z_H exchange are negligible.

To summarize, we find that Ma's model can satisfy the limits coming from Δm_K and $B(b\to\mu^+\mu^- X)$ as well as give values of Δm_{B_d} and Δm_{B_s} commensurate with the experimental measurements provided the mixings between SU(2) singlets and doublet members [for both the SU(2)'s] are of the order of 10^{-3}.

IV. COMPARISON WITH LEP RESULTS ON Z^0 PHYSICS

In the preceding section we saw that Ma's model with FCNC's is viable from the point of view of low-energy phenomenology; i.e., it can be made consistent with all the limits from flavor-changing interactions. Moreover, it gives values for B^0_d-$\bar{B}^0_d$ and B^0_s-$\bar{B}^0_s$ mixings which are consistent with the experimental data. This brings us to the final round of phenomenological tests—i.e., comparing the predictions of this model with the recent LEP measurements on Z^0 resonance. As can be seen from Eqs. (3)–(10), the $SU(2)_1$ couplings of the τ lepton and the b quark in this model are strikingly different from those in the SM. Thus the predictions of this model for the partial decay widths $Z^0\to\tau^+\tau^-, Z^0\to b\bar{b}$ and the total width are very different from the SM predictions. Unlike low-energy phenomena such as $e^+e^-\to b\bar{b}$ asymmetry and B^0_d-$\bar{B}^0_d$ mixing, these differences cannot be patched up by any compensatory contributions such as scalar exchange or FCNC's, because the partial widths $Z^0\to f\bar{f}$ on the resonance directly measure the coupling of Z^0 to the fermion f. Thus the recent measurements of the Z^0 partial widths at LEP provide an unambiguous test for this model. As we shall see below, the LEP results seem to rule out this model. Indeed the very recent measurements of the $(Z^0\to b\bar{b})$ partial width by ALEPH and L3 Collaborations seem to rule out any topless model and hence provide the most model-independent evidence for the top quark so far.

The partial width of Z^0 into a fermion pair $f\bar{f}$ is given by

$$\Gamma(Z^0\to f\bar{f}) = \frac{G_F M_Z^3}{3\pi\sqrt{2}}\mathcal{C}(g^2_{Lf}+g^2_{Rf}) = \frac{G_F M_Z^3}{3\pi\sqrt{2}} g^2_{Zf} , \quad (52)$$

where g_{Lf} and g_{Rf} are the couplings of the left- and right-handed fermions to Z^0. The color factor $\mathcal{C}$ is equal to 1 for leptons and to $3[1+\alpha_s(M_Z^2)/\pi]=3.105$ for quarks, where we have taken $\alpha_s(M_Z^2)=0.11$. The mass of the Z^0 has been measured to a very good accuracy at LEP, giving $M_Z = 91.154\pm 0.032$ GeV. G_F is measured from muon decay and is known to a great accuracy. Therefore the common factor in the partial widths is known to an accuracy of better than 0.1%:

$$\frac{G_F M_Z^3}{3\pi\sqrt{2}} = 663 \text{ MeV} . \quad (53)$$

To evaluate the partial widths we also need the value of $\sin^2\theta_W$, we take it to be $\sin^2\theta_W = 0.23$ (Ref. 12).

We have mentioned in Sec. II that the τ lepton and the b quark have nonstandard couplings to Z^0 in this model. Therefore let us first consider the partial widths $\Gamma(Z^0\to\tau^+\tau^-)$ and $\Gamma(Z^0\to b\bar{b})$. The chiral couplings of the τ lepton in the SM are

$$g_{L\tau} = -\tfrac{1}{2}+\sin^2\theta_W, \quad g_{R\tau} = \sin^2\theta_W . \quad (54)$$

For $\sin^2\theta_W = 0.23$, we get

$$g_{Z\tau}^2=0.1258 \text{ and } \Gamma(Z^0\to\tau^+\tau^-)=83.4 \text{ MeV} . \quad (55)$$

In the present model both τ_L and τ_R are $SU(2)_1$ doublets. Hence

$$g_{L\tau}=-\tfrac{1}{2}+\sin^2\theta=g_{R\tau} , \quad (56)$$

yielding

$$g_{Z\tau}^2=0.1458 \text{ and } \Gamma(Z^0\to\tau^+\tau^-)=96.7 \text{ MeV} . \quad (57)$$

Comparing the partial width values in (55) and (57) with the experimental measurement of $\Gamma(Z^0\to\tau^+\tau^-)|_{\text{expt}}=82.8\pm2.4$ MeV (Ref. 33), we see that the SM prediction agrees perfectly with the experimental value while the prediction of Ma's model is more than *six* standard deviations away.

In the case of the b quark, the chiral couplings in the SM are

$$g_{Lb}=-\tfrac{1}{2}+\tfrac{1}{3}\sin^2\theta_W, \quad g_{Rb}=\tfrac{1}{3}\sin^2\theta_W , \quad (58)$$

giving

$$g_{Zb}^2=0.5747 \text{ and } \Gamma(Z^0\to b\bar{b})=381 \text{ MeV} . \quad (59)$$

Both b_L and b_R are $SU(2)_1$ singlets in this model and the chiral couplings are

$$g_{Lb}=\tfrac{1}{3}\sin^2\theta_W=g_{Rb} , \quad (60)$$

with

$$g_{Zb}^2=0.0365 \text{ and } \Gamma(Z^0\to b\bar{b})=24.2 \text{ MeV} . \quad (61)$$

We see from (59) and (61) that the prediction of Ma's model for $\Gamma(Z^0\to b\bar{b})$ is one-sixteenth that of the SM. Very recently ALEPH and L3 Collaborations[34] have measured this partial width to be 353 ± 48 MeV. The Mark II Collaboration has also reported a similar value, but with much larger error.[34] Thus the SM prediction agrees perfectly with the experimental result whereas the prediction of Ma's model is in very strong conflict.

One can also compare the predictions of Ma's model with the measured hadronic width and total width of Z^0 because these have been measured at LEP to a very good accuracy.[33] The SM prediction for the hadronic partial width is 1735 MeV whereas the prediction in Ma's model is 1378 MeV. The latter is smaller by about 350 MeV because of the smaller decay rate into $b\bar{b}$ pairs. Comparing these values with the experimental measurement $\Gamma(Z^0\to\text{hadrons})=1804\pm44$ MeV, we see that the SM prediction is within 1.5 standard deviations while the prediction of Ma's model is off by nearly *ten* standard deviations. Although the fractional discrepancy between the prediction of this model and the data is less striking than for $Z^0\to b\bar{b}$ partial width, the hadronic width is measured more precisely and discrepancy between the model and the data is even worse.

Because of the smaller hadronic width, the prediction of Ma's model (MM) for the total decay width is smaller than that of the SM and is in conflict with the experimental data. The relevant numbers are

$$\Gamma_Z(\text{SM})=2482 \text{ MeV} , \quad (62)$$

$$\Gamma_Z(\text{MM})=2139 \text{ MeV} \quad \text{assuming 3 neutrino flavors} , \quad (63)$$

$$\Gamma_Z(\text{expt})=2534\pm45 \text{ MeV} . \quad (64)$$

The SM prediction is consistent with the experimental value within 1.2 standard deviations but Ma's model prediction is off by nearly *eight* standard deviations. The total width in Ma's model can be made larger if one assumes that more than three flavors of neutrinos are produced. It was mentioned in the discussion of the model that there are six flavors of neutrinos in the model which couple to Z^0 with canonical coupling. If we assume that 5 neutrino flavors are light enough to be produced in Z^0 decay then the total width in the model becomes 2471 MeV, which is in good agreement with the experimental result on Γ_Z. But it is in conflict with the measured value for the number of neutrinos,[33] $N_\nu=3.01\pm0.15$.

Thus the LEP measurements of Γ_Z, $\Gamma(Z^0\to\tau^+\tau^-)$, N_ν, $\Gamma(Z^0\to\text{hadrons})$, and $\Gamma(Z^0\to b\bar{b})$ rule out the topless model of Ma convincingly. The precision measurements of the LEP experiments, which are in excellent agreement with the predictions of the SM, are in strong disagreement with the predictions of this model for partial decay widths into $\tau^+\tau^-$, hadrons, and particularly $b\bar{b}$. These disagreements arise due to the nonstandard couplings of $\tau^+\tau^-$ and $b\bar{b}$ to Z^0. The LEP measurement of the $Z^0\to b\bar{b}$ partial width seems to rule out, indeed, any model with nonstandard b-quark couplings to the W and Z^0 bosons as we see below.

We assume the interaction of quarks and leptons with the standard W, Z bosons to be an $SU(2)\times U(1)$ gauge interaction, for which there is overwhelming phenomenological support.[2] Thus the $Z^0\to b\bar{b}$ partial width is proportional to $g_{Lb}^2+g_{Rb}^2$, where g_{Lb} and g_{Rb} are the couplings of b_L and b_R to Z^0, respectively; and we have seen that its measured value implies either b_L or b_R to be a member of an SU(2) doublet. In either case it has to be paired up with a heavier quark (i.e., top), since its pairing with any of the lighter quarks[35] would result in a b-decay rate nearly 400 times larger than the observed value. This provides by far the clearest evidence for the existence of a top quark. Moreover, the doublet assignment for b_R is also ruled out by the observed properties of b decay. The observed lepton spectrum in the semileptonic decay $b\to l\bar{\nu}X$ clearly shows that the couplings are either left handed or right handed for both the quark and the lepton vertices.[36] However, there are numerous phenomenological examples of evidence showing that the leptons and quarks of the first two generations form left-handed SU(2) doublets[2] for charged-current weak interaction. Thus the only solution, consistent with the above observations, is the standard SU(2)-doublet (-singlet) assignment for the left- (right-)handed b quarks. In short, the measured $Z^0\to b\bar{b}$ partial width along with the known properties of b decay provides by far the most unambiguous evidence for the presence of a top quark as the weak SU(2) partner of the left-handed b quark.

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