TeV scale horizontal gauge symmetry and its implications in B-physics (postprint)

Authors: Seong Chan Park, Jing Shu, Kai Wang, Tsutomu T. Yanagida

Date: 2017-08-03T00:00:00+00:00

Abstract

We propose a gauged U(1)H horizontal symmetry around TeV scale that is a subgroup of a SU(3)H horizontal gauge symmetry broken at O(1014) GeV. The breaking generates right-handed Majorana neutrino masses through a SU(3)H sextet scalar. A particular Majorana right-handed neutrino mass matrix explicitly determines the remnant U(1)H at low energy which only couples to b—s and μ— in the gauge eigenstate. The dangerous $K - K^{-}$, $D - D^{-}$ mixing and Bs → μ+μ— are kept to be safe because the relevant couplings are suppressed through high powers of small mixing angles in the fermion rotation matrix. Our analysis which applies to the general case shows that the Tevatron di-muon anomaly can be explained through the Bs and Bd mixing while keeping all the other experimental constraints within 90 % C. L. For the B meson decay, the Bs → μ± is the leading leptonic decay channel which is several orders of magnitude below current experimental bound

Full Text

Preamble

TeV Scale Horizontal Gauge Symmetry and Its Implications in B-Physics
IPMU10-0115

Seong Chan Park,[1] Jing Shu,[1,2] Kai Wang[1] and Tsutomu T. Yanagida[1]

[1] Institute for the Physics and Mathematics of the Universe, University of Tokyo, Kashiwa, Chiba 277-8568, Japan
[2] INPAC, Department of Physics, Shanghai Jiao Tong University, Shanghai, 200240, China

Abstract

We propose a gauged $U(1)_H$ horizontal symmetry around the TeV scale that is a subgroup of an $SU(3)_H$ horizontal gauge symmetry broken at $\mathcal{O}(10^{14})$ GeV. The breaking generates right-handed Majorana neutrino masses through an $SU(3)_H$ sextet scalar. A particular Majorana right-handed neutrino mass matrix explicitly determines the remnant $U(1)_H$ at low energy which only couples to $b-s$ and $\mu-\tau$ in the gauge eigenstate. The dangerous $K-\bar{K}$, $D-\bar{D}$ mixing and $B_s \to \mu^+\mu^-$ processes are kept safe because the relevant couplings are suppressed through high powers of small mixing angles in the fermion rotation matrix. Our analysis, which applies to the general case, shows that the Tevatron di-muon anomaly can be explained through the B_s and B_d mixing while keeping all other experimental constraints within 90% C.L. For B meson decays, the $B_s \to \mu^\pm\tau^\mp$ channel is the leading leptonic decay, which remains several orders of magnitude below current experimental bounds.

Introduction

Horizontal gauge symmetry was proposed as an extension of the SM gauge symmetries to unify all families of quarks and leptons [**?**, **?**]. Given the three families of quarks and leptons, $SU(3)_H$ is the most natural choice for the horizontal gauge symmetry. Interestingly, if one assumes all SM fermions transform as **3** under $SU(3)_H$, the anomaly-free condition requires three generations of right-handed neutrinos $n_{i=1,2,3}$ [**?**], while the right-handed neutrinos also play important roles in explaining the origin of neutrino masses. Therefore, the $SU(3)_H$ horizontal gauge symmetry model provides a natural scheme for the seesaw mechanism [**?**] generating small masses for light neutrinos [**?**]. The Majorana neutrino mass term for the right-handed neutrinos explicitly breaks the $SU(3)_H$, thus it is often believed that the horizontal gauge symmetry should be broken at a very high-energy scale $M_R \sim \mathcal{O}(10^{14})$ GeV. Then it seems impossible to test the $SU(3)_H$ gauge interactions in low-energy experiments. However, this is not always the case, as we will show in detail below. Even if some subgroup of the $SU(3)_H$ remains unbroken, the right-handed neutrinos can still acquire large Majorana masses.

Since n_R transform as **3** under $SU(3)_H$, the Majorana neutrino mass term can arise from the vacuum expectation value (vev) of an $SU(3)_H$ sextet χ_6: $\mathcal{L} \supset n_i^T \langle\chi_6\rangle_{ij} n_j$ and $M_R = \langle\chi_6\rangle$. The light neutrino mass is given by the seesaw mechanism as $m_\nu = m_D M_R^{-1} m_D^T$. In order to explain the neutrino oscillation data, suitable $\langle\chi_6\rangle$ and m_D are required. For m_D and the other SM fermion masses, there must exist octet Higgs fields under $SU(3)_H$ to accommodate the correct mass hierarchy in quarks and leptons. In addition, to minimize flavor-changing effects induced by the octet Higgses, we employ a scenario with additional Higgses and singlet fermions [**?**] in which the m_D and quark mass matrices or lepton mass matrix are all independent. By taking a suitable gauge choice of horizontal symmetry, we always choose M_R to be a diagonal matrix:

$$M_R = \langle \chi_6 \rangle = \begin{pmatrix} A & 0 & 0 \\ 0 & B & 0 \\ 0 & 0 & C \end{pmatrix}$$

The $\langle \chi_6 \rangle$ structure explicitly determines the symmetry breaking. The $SU(3)_H$ is completely broken in a generic vacuum with $A \neq B \neq C (\neq A)$. However, with a specific vacuum of $A = B = C$ for instance, the vacuum $\langle \chi_6 \rangle$ is invariant under an $SO(3)$ symmetry and the breaking pattern is $SU(3)_H \to SO(3)_H$. Being a symmetric second-rank tensor under $SU(3)_H$, the sextet χ_6 transforms as $\chi_6 \to U^T \chi_6 U$ where $U = e^{i\epsilon^a T^a}$ and T^a is the generator of the horizontal symmetry. A general scheme to obtain the unbroken symmetry is derived from the condition that if $[T, \langle \chi_6 \rangle] = 0$, then $\langle \chi_6 \rangle$ is invariant under the transformation defined by T.

To illustrate the feature of our proposal, we take a vacuum with $C = -B$. This vacuum $\langle \chi_6 \rangle = \mathrm{diag}(A, B, -B)$ is invariant under the $SU(3)$ generator λ_6 as $T = \lambda_6 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}$. Consequently, one can identify the unbroken $U(1)_H$ gauge symmetry with generator T and it can survive to low energy, for instance $\mathcal{O}(\mathrm{TeV})$, which may lead to interesting predictions in flavor-changing neutral current (FCNC) processes.

It was also observed that if there exist horizontal gauge interactions, CP violation can be realized with only two generations. Explicit examples of CP violation due to $U(1)_H$ and $SU(2)_H$ were discussed in [?]. If the above $U(1)_H$ is broken at low energy, the horizontal gauge boson exchanges can induce additional CP violations [?] at low energies through quark and lepton mixings.

In the last decades, huge experimental efforts have been made in improving the measurements on CP violation in the B meson system. Very recently, the D∅ Collaboration at Tevatron has reported a large charge asymmetry in like-sign di-muon events from both B_s and B_d decays with 6.1 fb^{-1}:

$$A_{sl}^{\mathrm{Exp}} \equiv \frac{N^{++} - N^{--}}{N^{++} + N^{--}} = (-9.57 \pm 2.51(\mathrm{stat.}) \pm 1.46(\mathrm{syst.})) \times 10^{-3},$$

where N^{++} (N^{--}) is the event number for $b\bar{b} \to \mu^+\mu^+X$ ($\mu^-\mu^-X$). Such a large di-muon charge asymmetry has a 3.2σ deviation from the SM prediction $A_{sl}^{\mathrm{SM}} = (-2.3^{+0.5}_{-0.6}) \times 10^{-4}$ [?] and many models have been proposed to account for this anomaly [?, ?, ?, ?]. The CDF measurement gives $A_{sl} = 8.0 \pm 9.0 \pm 6.8 \times 10^{-3}$ [?], using 1.6 fb^{-1} of data, which has a positive value and large uncertainties. Combining the above two results in quadrature (including the systematic uncertainty), we have $A_{sl} \approx -(8.5 \pm 2.8) \times 10^{-3}$.

At the Tevatron both B_d and B_s mesons are produced, hence A_{sl}^b is related to the charge asymmetries a_{sl}^d and a_{sl}^s in B_d and B_s decays by:

$$A_{sl}^b = (0.506 \pm 0.043)a_{sl}^d + (0.494 \pm 0.043)a_{sl}^s.$$

New physics (NP) contributions in B_d mixing are strictly constrained (we will show this more explicitly in the parameter fit later), so only large NP contributions to the B_s mixing (compared to the other meson mixings) are allowed. For the NP contribution, if the mixing in the rotation matrix between mass eigenstate and gauge eigenstate is not huge, then one would naturally expect the $U(1)_H$ that maximizes $b-s$ mixing as in Eq. (3). Indeed, for a CKM-like rotation matrix, the gauge boson coupling matrix at tree level in the mass basis behaves like:

$$\begin{pmatrix} \lambda^4 & \lambda^3 & -\lambda \\ \lambda^3 & -\lambda^2 & 1 \end{pmatrix},$$

where λ is the Wolfenstein parameter [**?**] around the order of the Cabibbo angle ($\lambda \approx 0.1$). Clearly, the meson mixings between the first two generations are highly suppressed. The NP also couples to leptons. However, their contributions to the B meson decay branching ratio to electron and muon are highly suppressed (although λ should be replaced by some small mixing of the lepton rotation matrix). Therefore, we focus on the phenomenology in the B meson mixing and decay.

The paper is organized as follows: in Section II, we propose the specific model considered in this paper. In Section III we show phenomenological implications of our model on flavor physics, with subsection III A related to meson mixing and subsection III B related to meson decay. Section IV contains our conclusions.

II. The Model

The model starts with a gauged $SU(3)_H$ model at extremely high energy, with all fermions transforming as **3** under $SU(3)_H$. The particle contents under $SU(3)_H \times SU(2)_L \times U(1)_Y$ are:

$$\begin{aligned} q_L &: (3, 2, \tfrac{1}{3}), & u_R &: (3, 1, \tfrac{4}{3}), \\ \ell_L &: (3, 2, -1), & d_R &: (3, 1, -\tfrac{2}{3}), \\ e_R &: (3, 1, -2), & n_R &: (3, 1, 0) \end{aligned}$$

which is exactly vectorial and the $SU(3)_H$ is therefore an anomaly-free symmetry. It is crucial to have right-handed neutrino triplets, $n_{i=1,2,3}$, for anomaly cancellation [**?**]. The extension to Pati-Salam unification [**?**] may be straightforward.

As we have already discussed the sextet breaking in the introduction, here we focus on the Yukawa interactions for the other SM fermions and the Dirac neutrino mass matrix. In conventional $SU(3)_H$ models, to break the $SU(3)_H$

as well as $SU(2)_L \times U(1)_Y$, one usually introduces one $H : (1, 2, 1)$ and four $\Phi_8 : (8, 2, 1)$ to generate all the SM fermion mass hierarchies. However, the $(8, 2, 1)$ Higgs will induce large FCNC [?] if the Higgs is light. To avoid the too-large FCNC problem, another proposal is to introduce $(8, 1, 0)$ Higgses [?]:

$$\Phi_8^{(i)} : (8, 1, 0), \quad H : (1, 2, 1) \quad (i = u, d, e, \nu)$$

In addition, to generate effective Yukawa couplings, a new set of $SU(2)_L$ singlet fermions is introduced:

$$\begin{array}{ll} U_L : (3, 1, \frac{1}{3}), & U_R : (3, 1, \frac{4}{3}), \\ D_L : (3, 1, -\frac{2}{3}), & D_R : (3, 1, -\frac{2}{3}), \\ E_L : (3, 1, -2), & E_R : (3, 1, -2), \\ N_L : (3, 1, 0), & N_R : (3, 1, 0). \end{array}$$

These singlet fermions form invariant Dirac masses and act as messengers to generate the necessary Yukawa interactions. We take the up-type quark mass matrix as an example. Since the octet Higgs is no longer an $SU(2)_L$ doublet, $q_L \bar{u}_R \Phi$ is forbidden and the up-quark Yukawa interactions only arise as:

$$\mathcal{L} \supset \bar{U}_L \Phi_8^{(u)} \bar{u}_R + M_U \bar{U}_L U_R + \bar{q}_L \bar{U}_R H + \lambda_u \bar{q}_L \bar{u}_R H,$$

where $\bar{q}_L \bar{u}_R H$ is universal. After integrating out the heavy fermion fields U_L, U_R, the effective up-quark Yukawa coupling reduces to:

$$\mathcal{L}_{\text{eff}} \supset \bar{u}_R (\lambda_u \delta_{ij} + (\langle \Phi_8^{(u)} \rangle M_U^{-1})_{ij}) q_j.$$

The same mechanism also applies to the mass generation of down-type quarks, charged leptons, and Dirac neutrinos. By assigning the $\langle \Phi_8^{(i)} \rangle$ independently, the mixings and masses in different fermion sectors are completely independent of each other and one can easily accommodate hierarchies and mixings in SM fermions and the Dirac neutrinos. This also enables us to choose the Dirac neutrino mass matrix other than a nearly-diagonal structure.

After electroweak symmetry breaking, the effective Yukawa coupling of $\Phi^{(u)}$ also arises: $\langle H \rangle M_U^{-1} \bar{u}_R \Phi_8^{(u)} u_L$. Then, the $\Phi_8^{(i)}$ exchanges induce FCNCs in general. We have checked that they satisfy the strongest constraint from $K - \bar{K}$ mixing marginally. However, actual effects depend on the mass spectrum of the $\Phi_8^{(i)}$ and hence we do not discuss them in this paper.

Another consequence is that both up and down quark mass matrices become Hermitian: $m_u^\dagger = m_u$, $m_d^\dagger = m_d$. Thus, the CP violation in strong interactions due to quark mass matrices, $\arg\{\det(m_u) \det(m_d)\}$, is absent at least at tree level [?]. In addition, the Hermitian mass matrices also require the rotations

U_L, U_R in the mass diagonalization $U_L^\dagger m_u U_R$ to be equal: $U_L = U_R$. In this case, the horizontal gauge boson couples to vector currents of quarks and leptons. As a consequence, pseudo-scalar bosons like B_s or B_d do not decay to a pair of leptons. However, this is only the result of our specific choice of mass generation model for quarks and leptons. In the following analyses, we assume more generic rotation matrices and take U_R and U_L to be independent of each other to estimate the predictions.

III. Phenomenological Implications in Flavor Physics

The horizontal gauge interaction is real but family-dependent. After mass diagonalization, other flavor violation entries as well as new CP violation can arise. The Lagrangian of gauge interactions is:

$$\mathcal{L}_H = g_H \bar{q}'_L \gamma^\mu q'_L \left(L Z'_\mu + L \leftrightarrow R \right) = g_H \bar{q}_i \gamma^\mu (V_q^L)^\dagger T V_q^L q_i Z'_\mu + L \leftrightarrow R,$$

where V_q^L stands for the rotation for left-handed q-type quarks and T is the generator of $U(1)_H$ interaction given in Eq. (3).

Flavor-changing interactions in the SM can only be measured via electroweak charged current interactions. Therefore, for the SM fermion rotation matrices, only the left-handed ones get constrained from the CKM matrix: $(V_u^L)^\dagger V_d^L = V_{\mathrm{CKM}}$, and one cannot determine V_u^L and V_d^L respectively. The other rotations are completely unknown. For simplicity of the discussion here, we will assume that all magnitudes of the left-handed mixings are CKM-like but the complex phases are $\mathcal{O}(1)$, and unconstrained right-handed mixings have a similar structure. Therefore, we have the mixing matrix in the mass eigenstates as:

$$(G')^{u/d}_{L/R} = (V^{u/d}_{L/R})^\dagger T (V^{u/d}_{L/R}) \to G' \sim \begin{pmatrix} \lambda^4 & \lambda^3 & -\lambda \\ \lambda^3 & -\lambda^2 & 1 \end{pmatrix}.$$

This $U(1)_H$ gauge interaction maximizes the mixing between the second and third generations. The mixing magnitude in B_s, B_d, and K^0 is of order $(1 : \lambda^2 : \lambda^6)$. The $D^0 - \bar{D}^0$ mixing is also suppressed at order λ^6 compared with B_s mixing. This $U(1)_H$ is consistent with the phenomenological constraints among different meson mixings. If one assumes the lepton doublet and right-handed singlet rotations are similar to the quark sector, one can also compute the leptonic decays of mesons. For instance, the $B_s \to \mu^+\mu^-$ decay partial width has a λ^4 suppression.

A. Meson Mixing

At the energy scale m_b, the effective Hamiltonian responsible for neutral meson mixing (and in particular $B_s - \bar{B}_s$ mixing) through the tree-level exchange of Z' is:

$$\mathcal{H} = C_{ij}^{LL}(m_b)\mathcal{O}_{ij}^{LL} + C_{ij}^{RR}(m_b)\mathcal{O}_{ij}^{RR} + C_{ij}^{LR}(m_b)\mathcal{O}_{ij}^{LR} + \tilde{C}_{ij}^{LR}(m_b)\tilde{\mathcal{O}}_{ij}^{LR},$$

where the $\Delta F = 2$ operators are given by:

$$\begin{aligned}
\mathcal{O}_{ij}^{LL} &= \bar{q}_i\gamma^\mu P_L q_j\, \bar{q}_i\gamma^\mu P_L q_j,\\
\mathcal{O}_{ij}^{RR} &= \bar{q}_i\gamma^\mu P_R q_j\, \bar{q}_i\gamma^\mu P_R q_j,\\
\mathcal{O}_{ij}^{LR} &= \bar{q}_i\gamma^\mu P_L q_j\, \bar{q}_i\gamma^\mu P_R q_j,\\
\tilde{\mathcal{O}}_{ij}^{LR} &= \bar{q}_i P_L q_j\, \bar{q}_i P_R q_j,
\end{aligned}$$

and the Wilson coefficients at the $M_{Z'}$ scale are ($\tilde{C}_{ij}^{LR}(M_{Z'}) = 0$):

$$\begin{aligned}
C_{ij}^{LL}(M_{Z'}) &= \frac{g_H^2}{M_{Z'}^2}(G'_{ij})_L^2,\\
C_{ij}^{RR}(M_{Z'}) &= \frac{g_H^2}{M_{Z'}^2}(G'_{ij})_R^2,\\
C_{ij}^{LR}(M_{Z'}) &= \frac{g_H^2}{M_{Z'}^2}(G'_{ij})_L(G'_{ij})_R,
\end{aligned}$$

where g_H is the horizontal gauge coupling at the $M_{Z'}$ scale and $M_{Z'}$ is the horizontal gauge boson mass.

In order to calculate the B physics observables, one must take into account the running effects of the four operators above. The relation between these four operators at the $M_{Z'}$ and m_b scales is presented in Appendix A. After obtaining the Wilson coefficients at the m_b scale, by using the relevant hadronic matrix elements [**?**]:

$$\begin{aligned}
\langle B_q|\mathcal{O}_{bq}^{LL/RR}|\bar{B}_q\rangle &\approx \frac{m_{B_q} f_{B_q}^2}{3} B_{B_q}^{LL/RR},\\
\langle B_q|\mathcal{O}_{bq}^{LR}|\bar{B}_q\rangle &\approx -\frac{5}{12}\left(\frac{m_{B_q}}{m_b+m_q}\right)^2 m_{B_q} f_{B_q}^2 B_{B_q}^{LR},\\
\langle B_q|\tilde{\mathcal{O}}_{bq}^{LR}|\bar{B}_q\rangle &\approx \frac{1}{2}\left(\frac{m_{B_q}}{m_b+m_q}\right)^2 m_{B_q} f_{B_q}^2 \tilde{B}_{B_q}^{LR},
\end{aligned}$$

we use $m_{B_q}^2/(m_b+m_q)^2 \approx 1$ and assume $B_{B_q}^{LR} \approx \tilde{B}_{B_q}^{LR} \approx B_{B_q}^{RR} = B_{B_q}^{LL} \equiv B_{B_q}$. Then we obtain:

$$M_{12}^q \equiv \langle B_q|\mathcal{H}|\bar{B}_q\rangle = -\frac{2}{3} f_{B_q}^2 m_{B_q} B_{B_q}\left[C_{bq}^{LL}(\mu) + C_{bq}^{RR}(\mu) - \frac{1}{4}C_{bq}^{LR}(\mu) - \frac{1}{4}\tilde{C}_{bq}^{LR}(\mu)\right].$$

From the discussion above, the flavor off-diagonal coupling between the horizontal gauge boson Z' and the first two generation quarks is highly suppressed (for a CKM-like rotation matrix, it is at least λ^3 suppressed), so we will neglect the new physics contributions to the $K-\bar{K}$ and $D-\bar{D}$ mixings. The Z'-b-s and Z'-b-d couplings, on the other hand, are either unsuppressed or λ suppressed, hence we expect large new physics contributions to modify the magnitudes and phases of $M_{12}^{d/s}$, where $M_{12}^{d/s}$ are off-diagonal mixing matrix elements in Eq. (23). We can parametrize such effects by:

$$M_{12}^{d/s} \equiv (M_{12}^{d/s})_{\rm SM}\Delta_{d/s}, \quad \Delta_s \equiv |\Delta_{d/s}|e^{i\phi_\Delta}.$$

The experimentally measured observables are summarized as follows: $\Delta m_{d/s}$ and $\Delta\Gamma_{d/s}$ measure the mass and decay width difference between the heavy and light mass eigenstates of the $B_{d/s}$ mesons. $a_{sl}^{d/s}$ is the charge asymmetry in semileptonic $B_{d/s}$ decays. β_d or β_s measure the time-dependent CP-violating phases in the hadronic B decay channels $B_d \to J/\psi K_S$ or $B_s \to J/\psi\phi$. They are shifted by the CP-violating phases in B_d or B_s mixing:

$$\begin{aligned}
\Delta m_{d/s} &= \Delta m_{d/s}^{\rm SM}|\Delta_{d/s}|, \\
\cos(\phi_{d/s}^{\rm SM} + \phi_\Delta) &= \frac{\Delta\Gamma_{d/s}}{2\Delta m_{d/s}}, \\
S_{\psi K} &= \sin(2\beta_d + \phi_\Delta^d), \\
\beta_s^{\rm Exp} &= \beta_s - \phi_\Delta^s.
\end{aligned}$$

The theoretical inputs are listed in Table I. All decay constants and bag parameters are taken from Ref. [**?**] which uses more recent decay constants and bag parameters with much smaller uncertainties. All other SM inputs are from [**?**] or [**?**]. Notice that we use the calculations for $\Delta m_{B_s} = 2|M_{12}^s|_{\rm SM}$ and ϕ_s from Ref. [**?**].

All experimental measurements used to compare with our model outputs are listed in Table II. For the like-sign dimuon charge asymmetry, we use $A_{sl}^b \approx -(8.5 \pm 2.8) \times 10^{-3}$, which combines the D∅ measurements with the CDF measurements. For $\beta_s^{\rm Exp}$ and $\Delta\Gamma_s$ measured by both CDF and D∅ [**?**, **?**, **?**, **?**], we use the combined results.

The experimental constraints on the parameter space of our model are presented in Fig. 1 and Fig. 2. The parameters $|G'_{bs}|$ and $\mathrm{Arg}(G'_{bs})$ are quite similar to the parameter h_2 in Ref. [**?**] except for an overall factor. For illustration, we choose parameters for B_d couplings $|G'_{bd}| = |G'_{bs}|/20$ and $\mathrm{Arg}(G'_{bd}) = \mathrm{Arg}(G'_{bs})$, in which there are sizable contributions to A_{sl}^b while the a_{sl}^d contribution is negligible. In contrast to Ref. [**?**], it is clear from the upper right plot in Fig. 1 that there is no region allowed by all experimental constraints within 1σ. The best-fit region for the phase $\mathrm{Arg}(G'_{bs}) \subset (\pi/2, 3\pi/4)$ is quite consistent with the one found in Ref. [?]. However, since Ref. [?] essentially marginalizes over Γ_{12} in the range $0 - 0.25$ ps^{-1} and uses the best-fit points where $\Delta\Gamma_s$ is about 2.5 times larger than the prediction, the goodness of fit in our result is reduced significantly compared to the one in Ref. [?].

B. Meson Decays

The new horizontal gauge boson can also mediate meson hadronic decays or leptonic decays at tree level. In this section, we choose to discuss the two leading processes, $b \to s\bar{c}c$ and $b \to s\mu^{\pm}\tau^{\mp}$ respectively, to illustrate how meson decays constrain the horizontal gauge interaction. Other transitions always have additional powers of λ suppression.

The effective $\Delta F = 1$ Hamiltonian responsible for neutral meson decay is:

$$\mathcal{H}_{\Delta F=1} = C_3 Q_3 + C_5 Q_5 + \tilde{C}_3 \tilde{Q}_3 + \tilde{C}_5 \tilde{Q}_5,$$

and for leptonic decays like $B^0_{d,s} \to \ell^+\ell^-$:

$$\mathcal{H}_{\ell}^{\text{eff}} = C_9 Q_9 + C_{10} Q_{10} + \tilde{C}_9 \tilde{Q}_9 + \tilde{C}_{10} \tilde{Q}_{10},$$

where $i, j = e, \mu, \tau$ are lepton flavor indices.

In the case of $b \to s\bar{c}c$ transition, the SM contribution at tree level is induced via weak charged current with a CKM factor $V_{bc}V_{sc}^* \sim \lambda^2$. Reading from the effective couplings, the horizontal gauge boson mediated $b \to s\bar{c}c$ has a factor of $G_{bs}G_{cc} \sim \lambda^2$. The SM and horizontal gauge interaction contributions are at the same order of λ, and one can simply compare their couplings and gauge boson masses to estimate the ratio. As discussed in the previous section, the $U(1)_H$ is broken at $\mathcal{O}$(TeV), which results in a suppression factor of $(g_H/g)^4(m_W^4/M_{Z'}^4) \sim 10^{-8}$. Consequently, the contribution to $b \to s\bar{c}c$ from the new horizontal gauge boson is completely negligible.

In the case of neutral meson mixing, the SM leading contribution is from box diagrams while the horizontal gauge boson contribution is at tree level. Therefore, even if the new horizontal gauge boson is of order TeV, it can still significantly change ΔM. For decay processes where SM tree-level contributions exist, the above argument always applies, so we won' t discuss constraints from such decays.

Since the horizontal gauge boson also couples to leptons, there is again tree-level contribution to meson leptonic decays. Within the SM framework, B_s pure leptonic decays are realized via electroweak penguin diagrams with $Z/\gamma^* \to \ell^+\ell^-$, and therefore there is no lepton flavor violation at all. The constraints on leptonic decays are mostly on leptons directly decaying from B mesons. If there exists a τ in the final state, τ decay will complicate the search due to $D^{\pm}$ decays. Therefore, the leading constraints are:

$$\mathrm{Br}(B_s \to \mu^+\mu^-) < 4.7 \times 10^{-8},$$
$$\mathrm{Br}(B_s \to e^+e^-) < 5.4 \times 10^{-5},$$
$$\mathrm{Br}(B_s \to e^\pm\mu^\mp) < 6.1 \times 10^{-6}.$$

In our model, the horizontal gauge boson has maximal coupling to b, s and μ, τ, and the leading leptonic decay constraint is from $B_s \to \mu^\pm\tau^\mp$. But as we mentioned, $B_s \to \mu^\pm\tau^\mp$ does not exist in SM physics at leading order and it only arises from new physics contributions.

Using $\langle 0|\bar{s}\gamma^\mu\gamma_5 b|\bar{B}^0\rangle = i f_B p_B^\mu$, one can compute the decay branching ratio as:

$$\mathrm{Br}(B_s \to \mu^\pm\tau^\mp) = \tau_B \Gamma(B^0 \to \mu^\pm\tau^\mp) = \frac{f_{B_s}^2 m_{B_s}^3 \tau_B}{8\pi} \left(\frac{g_H^2}{M_{Z'}^2}\right)^2 \left(1 - \frac{m_\tau^2}{m_{B_s}^2}\right)^2,$$

where the Wilson coefficients $C_{9,10} = g_H^2/M_{Z'}^2$. To estimate the decay branching ratio, we take $M_{Z'} \approx 10^3$ GeV, $g_H \approx 0.02$ and substitute $m_b = 4.7$ GeV, $m_\tau = 1.777$ GeV, $f_{B_s} = 230$ MeV, $m_B = 5.3$ GeV, $\tau_B = 1.6$ ps, obtaining $\mathrm{Br}(B_s \to \mu^\pm\tau^\mp) \approx 1.2 \times 10^{-9}$.

The only relevant search for $B_d \to \mu^\pm\tau^\mp$ is from CLEO [?], giving $\mathrm{Br}(B_d \to \mu^\pm\tau^\mp) < 3.8 \times 10^{-5}$. Due to the horizontal gauge boson coupling, the B_d decay partial width has an additional λ^2 suppression, so our prediction is well below the experimental bound. One can also estimate $B_s \to \mu^+\mu^-$ using the above result. The $B_s \to \mu^+\mu^-$ channel has a factor of λ^4 suppression, making the prediction about two orders of magnitude lower than the current experimental bound.

Other possible rare decays that can be induced by the horizontal gauge boson include FCNC decays in the top quark, for instance $t \to c/u + \mu^\pm\tau^\mp$. However, given the large $M_{Z'}$, the three-body decay is highly suppressed:

$$\Gamma(t \to c/u + \mu^\pm\tau^\mp) = \frac{m_t^5}{192\pi^3} \left(\frac{g_H^2}{M_{Z'}^2}\right)^2 \sim 1.7 \times 10^{-13} \text{ GeV}.$$

Even at a top factory like the Large Hadron Collider, it is impossible to observe such rare decay events.

IV. Conclusion

The dimuon asymmetry reported by the D∅ Collaboration is much larger than the SM prediction, which suggests new sources for CP violation. In this paper, we propose the possibility to explain such an anomaly through tree-level exchange of a gauged $U(1)_H$ horizontal symmetry in B meson mixing. The

$U(1)_H$ horizontal symmetry is a remnant symmetry of $SU(3)_H$ broken at $M_R \sim \mathcal{O}(10^{14})$ GeV through a sextet scalar which gives neutrino mass. Such a $U(1)_H$ gauge boson only couples to $b-s$ and $\mu-\tau$ in the gauge eigenstate, which suppresses all other dangerous meson mixings and B meson decays after flavor rotation. For a general flavor rotation matrix, we find there is a parameter region around the phase $\text{Arg}(G'_{bs}) \subset (\pi/2, 3\pi/4)$ which fits the data at 90% C.L. For B decays, the dominant enhanced channel is $B_s \to \mu^\pm\tau^\mp$. Nevertheless, such an enhanced decay channel is still one order of magnitude smaller than the current experimental bound.

V. Acknowledgements

We would like to thank Tao Liu, Hitoshi Murayama, Bai Yang, and Guohuai Zhu for useful discussions. The work is partially supported by the World Premier International Research Center Initiative (WPI Initiative) MEXT, Japan. S.C.P., J.S., and K.W. are also supported by the Grant-in-Aid for Scientific Research (Young Scientists (B) 21740172), (Young Scientists (B) 21740169), and (Young Scientists (B) 22740143) from JSPS, respectively. J.S. would like to acknowledge the hospitality of the Institute of Nuclear, Particle, Astronomy and Cosmology (INPAC), Shanghai Jiao Tong University, while the work was initiated.

Appendix A: RG Running of the $\Delta B = 2$ Operators

In the previous section, the Wilson coefficients are given at the $M_{Z'}$ scale, while to calculate physics processes involving low-energy mesons, one needs to calculate the relevant Wilson coefficients at the low-energy scale. The running contains two steps: the first step is from the $M_{Z'}$ scale to m_t where six flavors contribute to the running of α_s, and the second step is from m_t to m_b, where only five flavors contribute. We summarize the running effects of the relevant $\Delta F = 2$ operators below [**?**].

The operators belonging to the LL/RR, LR sectors are:

$$\begin{aligned}
Q_{LL} &= (\bar{s}_\alpha \gamma^\mu P_L b_\alpha)(\bar{s}_\beta \gamma^\mu P_L b_\beta), \\
Q_{RR} &= (\bar{s}_\alpha \gamma^\mu P_R b_\alpha)(\bar{s}_\beta \gamma^\mu P_R b_\beta), \\
Q_{LR} &= (\bar{s}_\alpha \gamma^\mu P_L b_\alpha)(\bar{s}_\beta \gamma^\mu P_R b_\beta), \\
\tilde{Q}_{LR} &= (\bar{s}_\alpha P_L b_\alpha)(\bar{s}_\beta P_R b_\beta).
\end{aligned}$$

The running relations are:

$$\begin{aligned}
C_{LL}(\mu_b) &= [\eta(\mu_b)] C_{LL}(\mu_t), \\
C_{RR}(\mu_b) &= [\eta(\mu_b)] C_{RR}(\mu_t), \\
\begin{pmatrix} C_{LR}(\mu_b) \\ \tilde{C}_{LR}(\mu_b) \end{pmatrix} &= \begin{pmatrix} \eta_{11}(\mu_b) & \eta_{12}(\mu_b) \\ \eta_{21}(\mu_b) & \eta_{22}(\mu_b) \end{pmatrix} \begin{pmatrix} C_{LR}(\mu_t) \\ \tilde{C}_{LR}(\mu_t) \end{pmatrix}.
\end{aligned}$$

The variables η_i are defined as ratios between strong coupling constants at different scales. Given the $U(1)_H$ gauge boson is around 10 TeV, we have two different running regimes for α_s taking into account the threshold correction due to the top quark.

For the running between the B physics scale and the top quark:

$$\eta_s \equiv \frac{\alpha_s(\mu_t)}{\alpha_s(\mu_b)}.$$

For the explicit form of η, we list both the LO (with subscript (0)) and NLO (with subscript (1)) expressions:

$$\begin{aligned}
&\eta^{(0)}(\mu_b) = \eta_s^{6/23},\\
&\eta_{11}^{(1)}(\mu_b) = \eta_s^{3/23},\\
&\eta_{12}^{(1)}(\mu_b) = 0,\\
&\eta_{21}^{(1)}(\mu_b) = \eta_s^{-24/23},\\
&\eta_{22}^{(1)}(\mu_b) = \eta_s^{-24/23},\\
&\eta^{(1)}(\mu_b) = 1.6273(1-\eta_s^5)\eta_s^{6/23},\\
&\eta_{11}^{(1)}(\mu_b) = 0.9250\eta_s^{-24/23},\\
&\eta_{12}^{(1)}(\mu_b) = 1.3875(\eta_s^{26/23} - \eta_s^{-24/23}),\\
&\eta_{21}^{(1)}(\mu_b) = (-11.7329 + 0.7829\eta_s^5)\eta_s^{3/23},\\
&\eta_{22}^{(1)}(\mu_b) = (7.9572 - 8.8822\eta_s^5)\eta_s^{-24/23} + (-2.0994 + 1.1744\eta_s^5).
\end{aligned}$$

Similarly, for the running between the top quark mass scale and the horizontal gauge boson scale, one can replace all μ_b, μ_t by $\mu_t, \mu_{M_{Z'}}$ in the above equations, where all the η_s are replaced by:

$$\eta_s \equiv \frac{\alpha_s(\mu_{M_{Z'}})}{\alpha_s(\mu_t)}.$$

The explicit LO and NLO expressions for this running regime are analogous to those above with appropriate modifications for the different number of active flavors.

References

[1] T. Maehara and T. Yanagida, Prog. Theor. Phys. 60, 822 (1978); T. Maehara and T. Yanagida, Prog. Theor. Phys. 61, 1434 (1979).
[2] F. Wilczek and A. Zee, Phys. Rev. Lett. 42, 421 (1979).
[3] T. Yanagida, Phys. Rev. D 20, 2986 (1979).

[4] T. Yanagida, in Proc. of the Workshop on "The Unified Theory and Baryon Number in the Universe" , Tsukuba, Japan (1979) p.95; M. Gell-Mann, P. Ramond and R. Slansky, in Supergravity, eds. P. van Nieuwenhuizen et al., (North-Holland, 1979), p. 315; P. Minkowski, Phys. Lett. B 67 (1977) 421.
[5] A. Masiero and T. Yanagida, arXiv:hep-ph/9812225.
[6] C. Amsler et al. [Particle Data Group], Phys. Lett. B667, 1 (2008).
[7] V. M. Abazov et al. [The DØ Collaboration], arXiv:1005.2757 [hep-ex].
[8] A. Lenz and U. Nierste, JHEP 0706, 072 (2007) [arXiv:hep-ph/0612167].
[9] M. Bona et al., arXiv:0906.0953 [hep-ph].
[10] E. Barberio et al. [Heavy Flavor Averaging Group], arXiv:0808.1297.
[11] V. M. Abazov et al. [DØ Collaboration], arXiv:0904.3907.
[12] V. M. Abazov et al. [DØ Collaboration], Phys. Rev. Lett. 101, 241801 (2008) [arXiv:0802.2255]; DØ Note 5933-CONF.
[13] T. Aaltonen et al. [CDF Collaboration], CDF Public Note 9458, August 7, 2008.
[14] CDF and DØ Collaborations, CDF note 9787, DØ Note 5928-CONF.
[15] L. Oakes (CDF Collaboration), talk at FPCP 2010, May 25–29, Torino, Italy.
[16] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49, 652 (1973).
[17] A. Dighe, A. Kundu and S. Nandi, Phys. Rev. D 82, 031502 (2010) [arXiv:1005.4051 [hep-ph]]; B. A. Dobrescu, P. J. Fox and A. Martin, Phys. Rev. Lett. 105, 041801 (2010) [arXiv:1005.4238 [hep-ph]]; C. H. Chen and G. Faisel, arXiv:1005.4582 [hep-ph]; C. W. Bauer and N. D. Dunn, arXiv:1006.1629 [hep-ph]; D. Choudhury and D. K. Ghosh, arXiv:1006.2171 [hep-ph]; J. K. Parry, arXiv:1006.5331 [hep-ph]; S. F. King, arXiv:1006.5895 [hep-ph].
[18] N. G. Deshpande, X. G. He and G. Valencia, arXiv:1006.1682 [hep-ph].
[19] C. H. Chen, C. Q. Geng and W. Wang, arXiv:1006.5216 [hep-ph].
[20] Y. Bai and A. E. Nelson, arXiv:1007.0596 [hep-ph].
[21] CDF Collaboration, Measurement of CP asymmetry in semileptonic B decays, Note 9015, Oct. 2007.
[22] L. Wolfenstein, Phys. Rev. Lett. 51, 1945 (1983).
[23] J. C. Pati and A. Salam, Phys. Rev. D 10, 275 (1974) [Erratum-ibid. D 11, 703 (1975)].
[24] S. L. Glashow and S. Weinberg, Phys. Rev. D 15, 1958 (1977).
[25] Z. Ligeti, M. Papucci, G. Perez and J. Zupan, arXiv:1006.0432 [hep-ph].
[26] V. Barger, L. L. Everett, J. Jiang, P. Langacker, T. Liu and C. E. M. Wagner, JHEP 0912, 048 (2009) [arXiv:0906.3745 [hep-ph]].
[27] G. Cacciapaglia, C. Csaki, J. Galloway, G. Marandella, J. Terning and A. Weiler, JHEP 0804, 006 (2008) [arXiv:0709.1714 [hep-ph]].
[28] M. Bona et al. [UTfit Collaboration], JHEP 0803, 049 (2008) [arXiv:0707.0636 [hep-ph]].
[29] E. Barberio et al. [Heavy Flavor Averaging Group], arXiv:0808.1297 [hep-ex].
[30] J. E. Duboscq, Nucl. Phys. Proc. Suppl. 144, 265 (2005) [arXiv:hep-ex/0412022].
[31] G. Ecker and W. Grimus, Z. Phys. C 30, 293 (1986).

[32] V. Barger, C. W. Chiang, P. Langacker and H. S. Lee, Phys. Lett. B 580, 186 (2004) [arXiv:hep-ph/0310073].
[33] J. Charles et al. [CKMfitter Group], Eur. Phys. J. C 41, 1 (2005) [arXiv:hep-ph/0406184].
[34] E. Barberio et al. [The Heavy Flavor Averaging Group], arXiv:hep-ex/0603003.
[35] G. Buchalla, A. J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. 68, 1125 (1996) [arXiv:hep-ph/9512380].
[36] C. W. Bauer and N. D. Dunn, arXiv:1006.1629 [hep-ph].
[37] J. Laiho, E. Lunghi and R. S. Van de Water, Phys. Rev. D 81, 034503 (2010) [arXiv:0910.2928 [hep-ph]].
[38] A. J. Buras, arXiv:hep-ph/9806471.
[39] A. J. Buras, S. Jager and J. Urban, Nucl. Phys. B 605, 600 (2001) [arXiv:hep-ph/0102316].

Note: Figure translations are in progress. See original paper for figures.