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Vector-like compositeness meets B-physics $R_{K^{(*)}}$ anomaly

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We have discussed the possibility where the observed deviations in the $b \to s\mu^+\mu^-$ data including the variables $R_{K^{(*)}}$ are addressed, in the context of vectorlike compositeness (or so called hypercolor). Major virtues of this type of scenario are as follows: (i) The theory is manifestly gauge anomaly free; (ii) A desirable form in flavor-changing interactions is naturally realized by exchanges of the hypercolor vector ρ mesons, which are composite spin-1 particles being analogous to the vector ρ mesons in the quantum chromodynamics. We can address the $R_{K^{(*)}}$ anomaly consistently, while negligible effects happen in the observables $R_{D^{(*)}}$ due to the structure of a global symmetry in the scenario.

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1. Introduction

Recently, deviations from the Standard Model (SM) have been reported in several B meson decays such as $B \to K^{(*)}\mu^+\mu^-$ and $B \to D^{(*)}\tau\bar{\nu}$ which have been receiving increased attention. A tremendous amount of studies have been made for addressing the anomalies in the data set of $b \to s\mu^+\mu^-$, including the observables $R_{K^{(*)}} = \mathscr{B}(B \to K^{(*)}\mu^+\mu^-)/\mathscr{B}(B \to K^{(*)}e^+e^-)$ (see *e.g.,* Refs. [1, 2, 3, 4, 5] for earlier works) and in the observables $R_{D^{(*)}} = \mathscr{B}(\bar{B} \to D^{(*)}\tau\bar{\nu})/\mathscr{B}(\bar{B} \to D^{(*)}\ell\bar{\nu})$ (see *e.g.,* Refs. [6, 7, 8, 9, 10] for earlier works) in various types of models.

We briefly look at situations. Results of global fits of the $b \to s\mu^+\mu^-$ data tell us that a new vector boson (or a vector leptoquark) which couples with corresponding left-handed components of quarks and leptons is a good candidate for explanation. Here, consistency in gauge anomaly cancellation puts a restriction on the choice of quantum numbers of fermions (see *e.g.,* [11]). Also we point out that, in explanations with fundamental vector boson and/or vector leptoquark, we need additional (vector-like) quarks and scalars (see *e.g.,* [12]). On the discrepancy in $R_{D^{(*)}}$, explanations by the exchange of scalars were shown to be disfavored by the B_c meson data, while vector bosons and leptoquarks would still have a chance (see [13] based on the discussions in [14, 15]).

To shed light on the idea that vector-like compositeness or hypercolor (HC) addresses such anomalies provides us a new point of view [16]. In the scenario where vector-boson and vector-leptoquark candidates are naturally realized as 'vector-ρ mesons' of a hidden strongly-coupled vector-like gauge theory, where the theory is manifestly anomaly-free. Here, the composite vector particles can contain couplings with the SM fermions in a gauge-invariant way, which is described by the language of hidden-local symmetry (HLS) (see a review, *e.g.,* [17]).

2. Model Description

2.1 One-family model of HC

We introduce the one-family model of the HC and provide outline the HC model scenario. The gauge group of the strongly-coupled HC interaction (at around a TeV scale) is assumed to be $SU(N_{\rm HC})$ with $N_{\rm HC} \geq 3$, where the eight HC fermions (F) belong to the fundamental representation. Among the HC fermions, namely HC quarks and HC leptons, their representations under the $SU(3)_c$, $SU(2)_W$ and $U(1)_Y$ gauge groups are identical with those of the quark doublets and lepton doublets, respectively. The summary of the charge assignment is listed in Table 1. The HC sector possesses the approximate global "chiral" $U(8)_{F_L} \times U(8)_{F_R}$ symmetry, which is explicitly broken in part by the SM gauging and possible vectorlike fermion masses.

At the scale $\Lambda_{\rm HC}$ ($\Lambda_{\rm HC} = \mathscr{O}(1-10)$ TeV) the HC gauge interaction gets strong to develop the nonzero "chiral" condensate $\langle \bar{F}^A F^B \rangle \sim \Lambda_{\rm HC}^3 \cdot \delta^{AB}$, where A and B represent indices for $SU(8)$ fundamental representations. The emergence of the condensation generates breakdown of the "chiral" symmetry in the eight HC fermions down to the vectorial one: $SU(8)_{F_L} \times SU(8)_{F_R} \to SU(8)_{F_V}$. As a result of the spontaneous breaking, the 63 Nambu-Goldstone (NG) bosons emerge, which will be pseudoscalars by the explicit breaking terms including the SM gauge interactions and possibly present vectorlike fermion mass terms like $m_F^0 \bar{F}F$ like in the quantum chromodynamics (QCD).

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	$SU(N_{\rm HC})$	$SU(3)_c$	$SU(2)_W$	$U(1)_Y$
$Q_{L/R} = \begin{pmatrix} U \\ D \end{pmatrix}_{L/R}$	$N_{\rm HC}$	**3**	**2**	$1/6$
$L_{L/R} = \begin{pmatrix} N \\ E \end{pmatrix}_{L/R}$	$N_{\rm HC}$	**1**	**2**	$-1/2$

Table 1: The SM charge assignment for eight HC fermions $F_{L/R} = (Q,L)^T_{L/R}$ in the one-family model.

By naively scaling the hadron spectroscopy in QCD, we may find 63 composite vectors (HC ρ mesons) as the next-to-lightest HC hadrons. Thus the low-energy effective theory of the HC sector would be constructed from the 63 HC pions ($\sim \bar{F}^A i\gamma_5 F^B$) and also 63 HC rho mesons ($\sim \bar{F}^A \gamma_\mu F^B$).

2.2 HLS formulation

The effective Lagrangian for those vectors (and pions) can be formulated based on the HLS formalism, which has succeeded in QCD rho meson physics. Based on the nonlinear realization of the HLS and the "chiral" $SU(8)_{F_L} \times SU(8)_{F_R}$ symmetry, the Lagrangian is written as

$$\mathscr{L} = -\frac{1}{2}{\rm tr}[\rho_{\mu\nu}^2] + f_\pi^2 {\rm tr}[\hat{\alpha}_{\perp\mu}^2] + \frac{m_\rho^2}{g_\rho^2}{\rm tr}[\hat{\alpha}_{||\mu}^2] + \cdots, \tag{2.1}$$

in a manner invariant under the $SU(8)_{F_L} \times SU(8)_{F_R} \times [SU(8)_{F_V}]_{\rm HLS}$ symmetries, including the terms of the lowest derivative order. Here we define

$$\rho_{\mu\nu} = \partial_\mu \rho_\nu - \partial_\nu \rho_\mu - i g_\rho [\rho_\mu, \rho_\nu], \tag{2.2}$$

$$\hat{\alpha}_{\perp\mu} = \frac{D_\mu \xi_R \cdot \xi_R^\dagger - D_\mu \xi_L \cdot \xi_L^\dagger}{2i}, \qquad \hat{\alpha}_{||\mu} = \frac{D_\mu \xi_R \cdot \xi_R^\dagger + D_\mu \xi_L \cdot \xi_L^\dagger}{2i}, \tag{2.3}$$

$$D_\mu \xi_{R(L)} = \partial_\mu \xi_{R(L)} - i g_\rho \rho_\mu \xi_{R(L)} + i \xi_{R(L)} \mathscr{R}_\mu(\mathscr{L}_\mu), \tag{2.4}$$

with the HLS gauge coupling g_ρ, the HC pion decay constant f_π, the HC ρ meson mass scale m_ρ, and the external gauge fields $\mathscr{R}_\mu$ and $\mathscr{L}_\mu$ that are associated by (external) gauging the "chiral" symmetry (see Table 1). Under the HLS and the "chiral" symmetry, the transformation properties for basic variables – $\xi_{L,R}$ (nonlinear bases), ρ_μ (HLS field), and $\hat{\alpha}_{\perp\mu}$, $\hat{\alpha}_{||\mu}$ (covariantized Maurer–Cartan one forms) – are described as

$$\xi_L \to h(x)\cdot \xi_L \cdot g_L^\dagger(x), \quad \xi_R \to h(x)\cdot \xi_R \cdot g_R^\dagger(x), \tag{2.5}$$

$$\rho_\mu \to h(x)\cdot \rho_\mu \cdot h^\dagger(x) + \frac{i}{g_\rho} h(x)\cdot \partial_\mu h^\dagger(x), \quad \rho_{\mu\nu} \to h(x)\cdot \rho_{\mu\nu}\cdot h^\dagger(x), \tag{2.6}$$

$$\hat{\alpha}_{\perp\mu} \to h(x)\cdot \hat{\alpha}_{\perp\mu}\cdot h^\dagger(x), \quad \hat{\alpha}_{||\mu} \to h(x)\cdot \hat{\alpha}_{||\mu}\cdot h^\dagger(x), \tag{2.7}$$

where $h(x) \in [SU(8)_{F_V}]_{\rm HLS}$ and $g_{R,L}(x) \in [SU(8)_{F_{R,L}}]_{\rm gauged}$. The nonlinear bases ξ_L and ξ_R can be parametrized by the NG bosons π for the "chiral" symmetry and $\mathscr{P}$ for the HLS. Hence, they are parametrized as

$$\xi_{\substack{R \\ L}} = e^{i\mathscr{P}/f_{\mathscr{P}}} \cdot e^{\pm i\pi/f_\pi}, \tag{2.8}$$

composite vector	constituent	color	isospin
$\rho^{\alpha}_{(8)a}$	$\frac{1}{\sqrt{2}}\bar{Q}\gamma_\mu\lambda^a\tau^\alpha Q$	octet	triplet
$\rho^{0}_{(8)a}$	$\frac{1}{2\sqrt{2}}\bar{Q}\gamma_\mu\lambda^a Q$	octet	singlet
$\rho^{\alpha}_{(3)c}$ $\left(\bar{\rho}^{\alpha}_{(3)c}\right)$	$\frac{1}{\sqrt{2}}\bar{Q}_c\gamma_\mu\tau^\alpha L$ (h.c.)	triplet	triplet
$\rho^{0}_{(3)c}$ $\left(\bar{\rho}^{0}_{(3)c}\right)$	$\frac{1}{2\sqrt{2}}\bar{Q}_c\gamma_\mu L$ (h.c.)	triplet	singlet
$\rho^{\alpha}_{(1)'}$	$\frac{1}{2\sqrt{3}}(\bar{Q}\gamma_\mu\tau^\alpha Q-3\bar{L}\gamma_\mu\tau^\alpha L)$	singlet	triplet
$\rho^{0}_{(1)'}$	$\frac{1}{4\sqrt{3}}(\bar{Q}\gamma_\mu Q-3\bar{L}\gamma_\mu L)$	singlet	singlet
$\rho^{\alpha}_{(1)}$	$\frac{1}{2}(\bar{Q}\gamma_\mu\tau^\alpha Q+\bar{L}\gamma_\mu\tau^\alpha L)$	singlet	triplet

Table 2: The HC rho mesons and their associated constituent HC quarks $Q_c = (U,D)_c$ and leptons $L = (N,E)$. Here λ^a $(a = 1,\cdots,8)$ are the Gell-Mann matrices, τ^α $SU(2)$ generators defined as $\tau^\alpha = \sigma^\alpha/2$ $(\alpha = 1,2,3)$ with the Pauli matrices σ^α, and the label c stands for the QCD-three colors, $c = r,g,b$. The numbers attached in lower scripts $(1,3,8)$ correspond to the representations under the QCD color, i.e., singlet, triplet and octet for $(1,3,8)$.

where the HLS decay constant $f_{\mathscr{P}}$ is related to the HC rho mass as $m_\rho = g_\rho f_{\mathscr{P}}$ and then the $\mathscr{P}$s are eaten by the HLS gauge boson ρ_μ due to the Higgs mechanism. Hereafter, we take the unitary gauge ($\mathscr{P}\equiv 0$). We easily decompose the 63 degrees of freedom into individual particles, where the detail for the ρ mesons is shown in Table 2. The similar classification is found in the π mesons.

2.3 Direct couplings to SM particles

To discuss gauge-invariant effective interactions between the ρ mesons and the SM fermions, we write down the fermions as an eight-dimensional vector on the base of the fundamental representation of $SU(8)$,

$$f_L = \begin{pmatrix} q \\ l \end{pmatrix}_L, \qquad f_R = \begin{pmatrix} q \\ l \end{pmatrix}_R, \tag{2.9}$$

where q and l are $SU(2)_{F_L,F_R}$ doublets for the quark and lepton fields. The SM-covariant derivatives that act on the f-fermion multiplets are then expressed as the 8×8 matrix forms:

$$\begin{aligned} D_\mu f_L &= \mathbf{1}_{8\times8}\cdot(\partial_\mu f_L) - i\,[\mathscr{L}^f_\mu]_{8\times8}\cdot f_L\,, \\ D_\mu f_R &= \mathbf{1}_{8\times8}\cdot(\partial_\mu f_R) - i\,[\mathscr{R}^f_\mu]_{8\times8}\cdot f_R\,, \end{aligned} \tag{2.10}$$

with

$$\begin{aligned} \left[\mathscr{L}^f_\mu\right]_{8\times8} &= \left(\begin{array}{c|c} \mathbf{1}_{2\times2}\otimes g_s G^a_\mu \frac{\lambda^a}{2} + \left(g_W W_\mu\tau^\alpha + \frac{1}{6}g_Y B_\mu\right)\otimes\mathbf{1}_{3\times3} & \mathbf{0}_{6\times2} \\ \hline \mathbf{0}_{2\times6} & g_W W^\alpha_\mu\tau^\alpha - \frac{1}{2}g_Y B_\mu\cdot\mathbf{1}_{2\times2} \end{array}\right) \\ &= \sqrt{2}g_s G^a_\mu T_{(8)a} + \frac{2}{\sqrt{3}}g_Y B_\mu T_{(1)'} + 2g_W W^\alpha_\mu T^\alpha_{(1)}, \\ \left[\mathscr{R}^f_\mu\right]_{8\times8} &= \left(\begin{array}{c|c} \mathbf{1}_{2\times2}\otimes g_s G^a_\mu\frac{\lambda^a}{2} + g_Y\,Q^q_{\rm em}B_\mu\otimes\mathbf{1}_{3\times3} & \mathbf{0}_{6\times2} \\ \hline \mathbf{0}_{2\times6} & g_Y\,Q^l_{\rm em}B_\mu \end{array}\right), \end{aligned} \tag{2.11}$$

where G_μ, W_μ and B_μ are the $SU(3)_c\times SU(2)_W\times U(1)_Y$ gauge fields along with the gauge couplings g_s, g_W and g_Y, respectively (see the caption of Table 2 for notations). $Q^{q,l}_{\rm em}$ are the electromagnetic charges defined as

$$Q^q_{\rm em} = \begin{pmatrix} 2/3 & 0 \\ 0 & -1/3 \end{pmatrix}, \qquad Q^l_{\rm em} = \begin{pmatrix} 0 & 0 \\ 0 & -1 \end{pmatrix}. \tag{2.12}$$

We may relate the charges of the HC fermions with those of the SM quark and lepton charges, involving the HC-quark and -lepton numbers. Then the nonlinear bases $\xi_{L,R}$ in Eq. (2.5) transform under the HLS and the SM gauge group $\mathscr{G} = SU(3)_c \times SU(2)_W \times U(1)_Y$ as

$$\xi_L \to h(x) \cdot \xi_L \cdot [g_L^\dagger(x)]_{\mathscr{G}}, \qquad \xi_R \to h(x) \cdot \xi_R \cdot [g_R^\dagger(x)]_{\mathscr{G}}. \tag{2.13}$$

From Table 1, one thus finds that the external gauge fields $\mathscr{L}_\mu$ and $\mathscr{R}_\mu$, coupled to the nonlinear bases $\xi_{L,R}$ as in Eq. (2.8), are identified with those coupled to the SM fermions as described in Eq.(2.10). When we focus on the vector ρ mesons and the external SM gauge bosons, the 1-forms in Eqs. (2.3) and (2.4) are represented as

$$\hat{\alpha}_{||\mu} = \mathscr{L}^f_\mu - g_\rho \rho_\mu + \cdots, \qquad \hat{\alpha}_{\perp\mu} = 0 + \cdots. \tag{2.14}$$

We may define the dressed fields for the left-handed SM fermions,

$$\Psi_L \equiv \xi_L \cdot f_L, \qquad \psi_L \equiv \xi_R \cdot f_L, \tag{2.15}$$

which transform as

$$\Psi_L \to h(x) \cdot \Psi_L, \qquad \psi_L \to h(x) \cdot \psi_L. \tag{2.16}$$

These transformations allow us to write down the HC ρ couplings to the left-handed SM fermions in the HLS-invariant way as

$$\mathscr{L}_{\rho ff} = g^{ij}_{1L} \left(\bar{\Psi}^i_L \gamma^\mu \hat{\alpha}_{||\mu} \Psi^j_L \right) + g^{ij}_{2L} \left(\bar{\Psi}^i_L \gamma^\mu \hat{\alpha}_{||\mu} \psi^j_L + \text{h.c.} \right) + g^{ij}_{3L} \left(\bar{\psi}^i_L \gamma^\mu \hat{\alpha}_{||\mu} \psi^j_L \right), \tag{2.17}$$

where i and j label the generations of the SM fermions $(i, j = 1, 2, 3)$.

Using Eqs. (2.10) and (2.14), one can thus extract the HC ρ and $V_{\rm SM}$ (SM gauge boson) couplings to the left-handed SM fermions. As a result, we have

$$\begin{aligned} \mathscr{L}^{\rm direct}_{V f_L f_L} = \; & g^{ij}_L \cdot \bar{q}^i_L \gamma_\mu \left[g_s G^\mu_a \left(\mathbf{1}_{2\times2} \otimes \frac{\lambda_a}{2} \right) + \left(g_W W^{\alpha\mu} \frac{\sigma^\alpha}{2} + \frac{g_Y}{6} B^\mu \mathbf{1}_{2\times2} \right) \otimes \mathbf{1}_{3\times3} - g_\rho \rho^\mu_{QQ} \right] q^j_L \\ & + g^{ij}_L \cdot \bar{l}^i_L \gamma_\mu \left[g_W W^{\alpha\mu} \frac{\sigma^\alpha}{2} - \frac{g_Y}{2} B^\mu \mathbf{1}_{2\times2} - g_\rho \rho^\mu_{LL} \right] l^j_L \\ & - g^{ij}_L g_\rho \cdot \left[\bar{q}^i_L \gamma_\mu \rho^\mu_{QL} l^j_L + \text{h.c.} \right], \end{aligned} \tag{2.18}$$

where $g^{ij}_L = (g_{1L} + 2g_{2L} + g_{3L})^{ij}$; ρ^μ_{QQ}, ρ^μ_{LL}, and ρ^μ_{QL} are combinations of the HC ρ mesons as defined as

$$\rho = \begin{pmatrix} (\rho_{QQ})_{6\times6} & (\rho_{QL})_{6\times2} \\ (\rho_{LQ})_{2\times6} & (\rho_{LL})_{2\times2} \end{pmatrix}, \tag{2.19}$$

with

$$\begin{aligned}
\rho_{QQ} &= \left[\sqrt{2}\rho^{\alpha}_{(8)a}\left(\tau^{\alpha}\otimes\frac{\lambda^a}{2}\right)+\frac{1}{\sqrt{2}}\rho^{0}_{(8)a}\left(\mathbf{1}_{2\times2}\otimes\frac{\lambda^a}{2}\right)\right] \\
&\quad+\left[\frac{1}{2}\rho^{\alpha}_{(1)}\left(\tau^{\alpha}\otimes\mathbf{1}_{3\times3}\right)+\frac{1}{2\sqrt{3}}\rho^{\alpha}_{(1)'}\left(\tau^{\alpha}\otimes\mathbf{1}_{3\times3}\right)+\frac{1}{4\sqrt{3}}\rho^{0}_{(1)'}\left(\mathbf{1}_{2\times2}\otimes\mathbf{1}_{3\times3}\right)\right], \\
\rho_{LL} &= \frac{1}{2}\rho^{\alpha}_{(1)}\left(\tau^{\alpha}\right)-\frac{\sqrt{3}}{2}\rho^{\alpha}_{(1)'}\left(\tau^{\alpha}\right)-\frac{\sqrt{3}}{4}\rho^{0}_{(1)'}\left(\mathbf{1}_{2\times2}\right), \\
\rho_{QL} &= \rho^{\alpha}_{(3)c}\left(\tau^{\alpha}\otimes\mathbf{e}_c\right)+\frac{1}{2}\rho^{0}_{(3)c}\left(\mathbf{1}_{2\times2}\otimes\mathbf{e}_c\right), \\
\rho_{LQ} &= \left(\rho_{QL}\right)^{\dagger}.
\end{aligned} \tag{2.20}$$

Note that the V_{SM}-f_L-f_L term in Eq. (2.18) is not the normal SM interactions but additional contributions in this model. $\mathbf{e}_c$ represents the three-dimensional unit vector in color space (see also the caption of Table 2). This type of vector interactions where only the left-handed part is active is one of the desirable cases pointed out, *e.g.,* in Ref. [18].

3. Addressing flavor anomalies

As can be seen in Eq. (2.18), our model involves lots of new interactions at the tree level, most of which are obviously already disfavored. In particular, it is easily expected that couplings to the first and second generations are severely constrained. To avoid such matters as well as to address the flavor anomalies in B decays, the reasonable setup may be given as [19]

$$g_L^{ij} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & g_L^{33} \end{pmatrix}^{ij}. \tag{3.1}$$

and flavor-changing effects are assumed to be induced from the mixing effects from the gauge basis to the mass basis

$$(u_L)^i = U^{iI}(u'_L)^I, \quad (d_L)^i = D^{iI}(d'_L)^I, \quad (e_L)^i = L^{iI}(e'_L)^I, \quad (\nu_L)^i = L^{iI}(\nu'_L)^I, \tag{3.2}$$

where U, D, and L are three-by-three unitary matrices and the spinors with the prime symbol denote the fermions in the mass basis (I being the associated index),

$$D = \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos\theta_D & \sin\theta_D \\ 0 & -\sin\theta_D & \cos\theta_D \end{pmatrix}, \qquad L = \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos\theta_L & \sin\theta_L \\ 0 & -\sin\theta_L & \cos\theta_L \end{pmatrix}. \tag{3.3}$$

The Cabibbo–Kobayashi–Maskawa (CKM) matrix element is then given by $V_{\mathrm{CKM}} \equiv U^{\dagger}(\mathbf{1}+\Delta_W)D \simeq U^{\dagger}D$ with $\Delta_W^{33} \leq \mathscr{O}(10^{-3})$ taken into account. Employing Fierz transformations, the operators are simplified into six types:

$$\mathscr{O}^{ijkl}_{4q(1)} = \left(\bar{q}^i_L \gamma_\mu q^j_L\right)\left(\bar{q}^k_L \gamma^\mu q^l_L\right), \quad \mathscr{O}^{ijkl}_{4q(3)} = \left(\bar{q}^i_L \gamma_\mu \sigma^\alpha q^j_L\right)\left(\bar{q}^k_L \gamma^\mu \sigma^\alpha q^l_L\right), \tag{3.4}$$

$$\mathscr{O}^{ijkl}_{4\ell(1)} = \left(\bar{l}^i_L \gamma_\mu l^j_L\right)\left(\bar{l}^k_L \gamma^\mu l^l_L\right), \quad \mathscr{O}^{ijkl}_{4\ell(3)} = \left(\bar{l}^i_L \gamma_\mu \sigma^\alpha l^j_L\right)\left(\bar{l}^k_L \gamma^\mu \sigma^\alpha l^l_L\right), \tag{3.5}$$

$$\mathscr{O}^{ijkl}_{2q2\ell(1)} = \left(\bar{q}^i_L \gamma_\mu q^j_L\right)\left(\bar{l}^k_L \gamma^\mu l^l_L\right), \quad \mathscr{O}^{ijkl}_{2q2\ell(3)} = \left(\bar{q}^i_L \gamma_\mu \sigma^\alpha q^j_L\right)\left(\bar{l}^k_L \gamma^\mu \sigma^\alpha l^l_L\right), \tag{3.6}$$

along with the Wilson coefficients having the form $g^2_\rho g^{ij}_L g^{kl}_L/m^2_\rho$. Here we assume that all of the masses of the ρ vector mesons are degenerated, which is justified with accuracy [16]. The above operators contribute to the following phenomena:

- $\mathscr{O}_{2q2\ell(n)}$: $\bar{B} \to D^{(*)}\tau\bar{\nu}$, $\bar{B} \to K^{(*)}\mu^+\mu^-$, $\bar{B} \to K^{(*)}\nu\bar{\nu}$, and $\tau \to \phi\mu$,
- $\mathscr{O}_{4q(n)}$: B_s–$\bar{B}_s$ mixing,
- $\mathscr{O}_{4\ell(n)}$: $\tau \to 3\mu$.

In Fig. 1, we summarize the situation in the parameter point, $|g_\rho g^{33}_L| = 1$ and $m_\rho = 1\,\mathrm{TeV}$, as a function of the mixing angles θ_D and θ_L, where we find the region where we can realize the target magnitude (where we adopted the value in [18] derived through their global fit) for addressing the anomaly in the $b \to s\mu^+\mu^-$ data consistently. Here, we should swallow the parameter tuning in θ_D as $\lesssim \mathscr{O}(10^{-2})$ to avoid a sizable tree-level contribution to the B_s–$\bar{B}_s$ mixing. The bound from $\tau \to 3\mu$ restricts the valid range of θ_L to the tau-philic region ($\theta_L \lesssim \pi/8$) or the mu-philic region ($\theta_L \sim \pi/2$). We provide two comments. With smaller values of $|g_\rho g^{33}_L|$ with keeping $m_\rho = 1\,\mathrm{TeV}$ and $\theta_D \lesssim \mathscr{O}(10^{-2})$, explanations are still possible since the $\tau \to 3\mu$ constraint is relaxed (where a larger θ_L gets to be allowed). When $\theta_L \sim \pi/2$, (almost) no constraint comes from $\tau \to 3\mu$ and then we can address the $R_{K^{(*)}}$ anomaly in a wide part of such parameter space.

We briefly comment on the contribution to the $R_{D^{(*)}}$ variables. It was pointed out that the vanishing contribution is observed for the degenerated HC ρs due to the global $SU(8)$ structure. Thereby, we cannot address the anomaly in $R_{D^{(*)}}$ in the setup.

4. Constraint from 13 TeV LHC dilepton searches

The latest null results in the new physics searches in the dilepton final states in the Large Hadron Collider (LHC) would put a significant bound on the parameter space of the present scenario. We derived analytical formulas of the cross sections of the processes $b\bar{b} \to \rho'\mathrm{s} \to \tau\bar{\tau}/\mu\bar{\mu}$. For simplicity, we focus on the 'maximal' cross section without suppression from the lepton mixing angle ($\theta_L = 0$ for ditau channel and $\theta_L = \pi/2$ for dimuon channel) and the quark mixing angle ($\theta_D = 0$). For numerical calculations, we use the `CUBA` package [20] with the `Mathlink` protocol in `Mathematica`.

Our results are summarized in Fig. 2. In the tau-philic case, when $m_\rho \gtrsim 1\,\mathrm{TeV}$, the constraint on the $R_{K^{(*)}}$ anomaly explanation is still week, where only possibilities near the 'maximized $g_{\rho L}$' are discarded. The evaluated bound on the dimuon case is much more stringent, where the configuration with $g_\rho g^{33}_L = 1$ is excluded if $m_\rho \lesssim 4\,\mathrm{TeV}$. However, from the right panel of Fig. 1, we recognize that addressing the anomaly is still possible if $|g_\rho g^{33}_L| \times (1\,\mathrm{TeV})/m_\rho \lesssim 0.15$ since

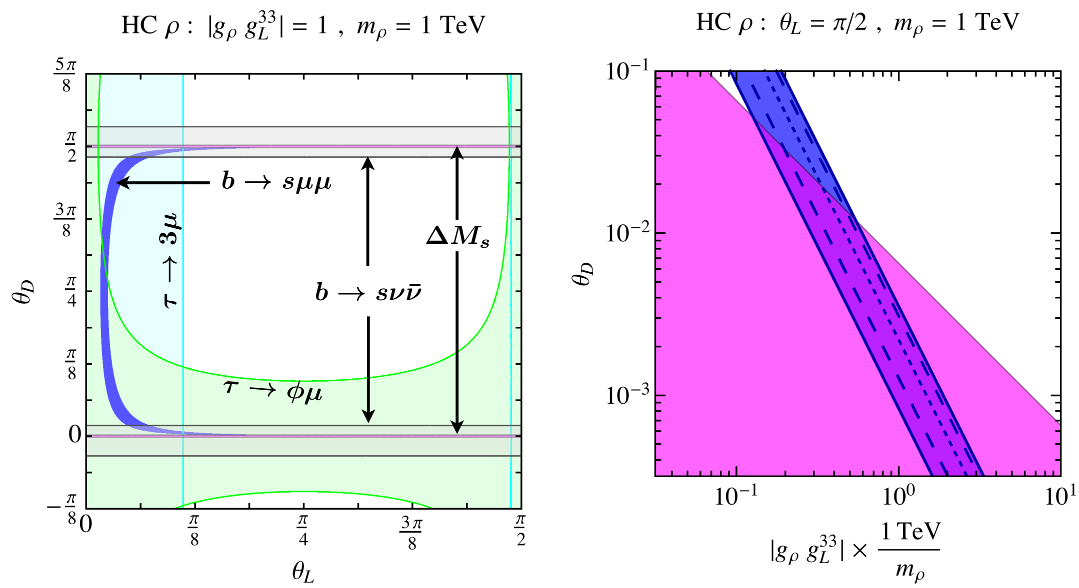

Figure 1: Allowed regions in the (θ_L, θ_D) plane in the HC ρ model for $m_\rho = 1\,\text{TeV}$ and $g_\rho\, g_L^{33} = 1$. The $b \to s\mu^+\mu^-$ anomaly can be explained in the blue region while the constraints from ΔM_s, $\mathscr{B}(\tau \to 3\mu)$, $\mathscr{B}(\tau \to \phi\mu)$, and $\mathscr{B}(B \to K^{(*)}\nu\bar{\nu})$ are satisfied in the magenta, cyan, green, and gray regions, respectively (left panel); The allowed range in terms of $|g_\rho\, g_L^{33}|/m_\rho$ and θ_D for the fixed value $\theta_L = \pi/2$. The color convention is the same as in the left panel (right panel) [16].

$|g_\rho g_L^{33}| \times (1\,\text{TeV})/m_\rho = 0.25$ at the benchmark point $(m_\rho, g_\rho g_L^{33}) = (4\,\text{TeV}, 1)$. Further data accumulation is (also) required for testing the whole region for explaining the anomaly shown in Fig. 1 in the mu-philic regime.

5. Miscellaneous issues

In this part, we provide comments on miscellaneous issues and points in the original paper [16].

- Part of the HC ρ mesons $(\rho^{0\mu}_{(8)a}, \rho^{\alpha\mu}_{(1)}, \rho^{0\mu}_{(1)'})$ mix with the SM gauge bosons $(G^{a\mu}, W^{\alpha\mu}, B^\mu)$, which is apparent from the mass terms in Eq. (2.1),

$$\begin{aligned}
\frac{m_\rho^2}{g_\rho^2}\mathrm{tr}[\hat{\alpha}^2_{||\mu}] \supset \frac{m_\rho^2}{g_\rho^2}\mathrm{tr}\left[(\mathscr{L}^f_\mu - g_\rho\rho_\mu)^2\right] \\
= \frac{1}{2}\frac{m_\rho^2}{g_\rho^2}\Bigg[g_\rho^2\left(\rho^{\alpha\mu}_{(8)a}\right)^2 + g_\rho^2\left(\rho^{\alpha\mu}_{(1)'}\right)^2 + \left(g_\rho\rho^{0\mu}_{(8)a} - \sqrt{2}g_s G^{a\mu}\right)^2 + \left(g_\rho\rho^{\alpha\mu}_{(1)} - 2g_W W^{\alpha\mu}\right)^2 \\
+ \left(g_\rho\rho^{0\mu}_{(1)'} - \frac{2}{\sqrt{3}}g_Y B^\mu\right)^2 + 2g_\rho^2\left(\overline{\rho}^{\alpha\mu}_{(3)c}\rho^{\alpha}_{(3)c\mu} + \overline{\rho}^{0\mu}_{(3)c}\rho^{0}_{(3)c\mu}\right)\Bigg].
\end{aligned} \tag{5.1}$$

This means that the ρ mesons have generation-independent "SM-gauge" interactions through the mass mixing, where typical magnitudes of the mixing angle are estimated as $g_{s,W,Y}/g_\rho$

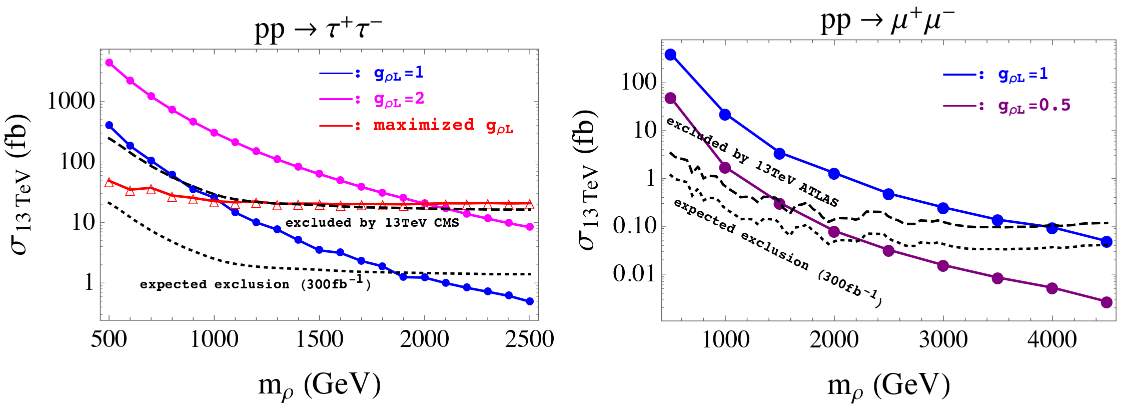

Figure 2: Constraints on $\sigma(pp \to \tau^+\tau^-)$ [left panel] and $\sigma(pp \to \mu^+\mu^-)$ [right panel] in $\theta_D = \theta_L = 0$ are calculated, where the CMS and ATLAS experimental results are from Refs. [21, 22]. The red curves show 'maximized $g_{\rho L}$' being consistent with the requisites from flavor issues [16], where the value of $g_{\rho L}$ is tuned as $|g_{\rho L}| = 1.0 \times (m_\rho/1\,\text{TeV})$. The expected exclusions are estimated simply by scaling the integrated luminosity.

(irrespective of the values of g_L^{ij} in Eq. (2.18)). Under nonzero mixings, small mass splittings emerge among the ρ mesons. In our analysis, we have focused on the case $g_{s,W,Y}/g_\rho \ll 1$ (where all of ρs are degenerated).

- The existence of the above-mentioned V_{SM}-ρ mixing leads to additional contributions to elecroweak precision variables if g_ρ takes a finite value. We showed that such corrections are almost irrelevant as long as g_ρ is ~ 6 (the corresponding value of the QCD) or greater.

- The present HC theory consists of the one-family content with the number of HC fermions $N_F = 8$, where the masses of HC pions having the SM charges could be enhanced by the amplification of the explicit breaking effect, as discussed in [23] and references therein. According to [23], the size of colored HC pion masses from the QCD gluon exchange contribution is evaluated as $M^2_{\pi_{(3),(8)}} \sim C_2\alpha_s(M_\pi)\Lambda^2_{\text{HC}}\ln\left(\Lambda^2_{\text{UV}}/\Lambda^2_{\text{HC}}\right)$, with $C_2 = \frac{4}{3}\,(3)$ for color-triplet (octet) HC pions, where Λ_{UV} denotes some ultraviolet high-energy scale up to which the HC theory is valid. Taking $\alpha_s(M_\pi) \sim 0.1$ and $\Lambda_{\text{UV}} \sim 10^{16}$ GeV as our benchmark, the $\pi_{(3)}$ and $\pi_{(8)}$ masses are estimated as $M_{\pi_{(3)}} \sim 3$ TeV and $M_{\pi_{(8)}} \sim 4$ TeV, respectively, for $\Lambda_{\text{HC}} \sim 1$ TeV. Due to the enhancement in the electroweak gauge interactions, the color-singlet pions $\pi^{\pm,3}_{(1)'}$ and $\pi^{\pm,3}_{(1)}$ are also uplifted as $M_{\pi^{\pm,3}_{(1)'}} \sim 2\,\text{TeV}$ and $M_{\pi^{\pm,3}_{(1)}} \sim 2\,\text{TeV}$, respectively.

- No such enhancement happens in the remaining SM-gauge-sterile HC pion $\pi^0_{(1)'}$, where its mass scale is roughly estimated as $M_{\pi^0_{(1)'}} \sim \mathcal{O}(f_\pi) = \mathcal{O}(100)\,\text{GeV}$ (f_π: HC pion decay constant). Through the Wess–Zumino–Witten terms, the process $GG \to \pi^0_{(1)'} \to \gamma\gamma$ is induced and the constraint from the LHC diphoton searches looks nontrivial. We showed that we can evade the bound by arranging the parameters f_π and $M_{\pi^0_{(1)'}}$.

6. Summary

We analyzed the flavor structure of composite vector bosons arising in a model of vectorlike confinement – often called hypercolor – with eight flavors that form a one-family content of HC fermions, namely HC quark doublets and lepton doublets. This theory is apparently gauge anomaly free, and if the HC quarks and leptons hold the quantum numbers of the SM quarks and leptons, respectively, a desirable interaction pattern (left-handed: active, right-handed: sterile) for addressing the anomaly in the $b \to s\mu^+\mu^-$ data is realized. We showed that the $R_{K^{(*)}}$ anomaly can be suitably addressed in the scenario, while only minuscule contributions to the observables $R_{D^{(*)}}$ occur due to the structure of the $SU(8)$ flavor symmetry.

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