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On theories of enhanced CP violation in $B_{s,d}$ meson mixing

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Abstract: The DØ collaboration has measured a deviation from the standard model (SM) prediction in the like sign dimuon asymmetry in semileptonic b decay with a significance of $3.2\,\sigma$. We discuss how minimal flavour violating (MFV) models with multiple scalar representations can lead to this deviation through tree level exchanges of new MFV scalars. We review how the two scalar doublet model can accommodate this result and discuss some of its phenomenology. Limits on electric dipole moments suggest that in this model the coupling of the charged scalar to the right handed u-type quarks is suppressed while its coupling to the d-type right handed quarks must be enhanced. We construct an extension of the MFV two scalar doublet model where this occurs naturally.

Keywords: Beyond Standard Model, CP violation, B-Physics, Higgs Physics

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Contents

1 Introduction

The DØ collaboration has reported a $3.2\,\sigma$ deviation from the standard model (SM) prediction of the like sign dimuon asymmetry in semileptonic b decay [1]. This observation joins past anomalous measurements of $B_s \to J/\psi\,\phi$ and $B^- \to \tau\,\nu$ decays that can be interpreted as a pattern of deviations consistent with new physics contributing a new phase in $B_{s,d}$ mixing (for a recent global fit and discussion see [2]).[1]

If the explanation of the like sign dimuon asymmetry measurement and these correlated deviations are not statistical fluctuations, then new physics interpretations of this pattern are of interest. General operator analyses have been carried out [10–12] and indicate that operators induced by scalar exchange with unenhanced Yukawa couplings and order one parameters in the potential (i.e. order one Wilson coefficients) the mass scale suppressing the operators of interest is a few hundred GeV.

Such a low mass scale is challenging to reconcile with known constraints from flavour physics unless minimal flavour violation (MFV) [13–15] is imposed. New physics (NP) models with MFV have the quark flavor group $\mathrm{SU}(3)_{U_R} \times \mathrm{SU}(3)_{D_R} \times \mathrm{SU}(3)_{Q_L}$ only broken by the Yukawa couplings. However this scenario does allow new phases and so provides

[1]The observed $2.6\,\sigma$ deviation from the standard model (SM) expectation [2] in the averaged measurements of $B^- \to \tau\,\nu$ performed at Belle and Babar [3–6] correlates correctly with a new physics (NP) contribution of a phase to B_d with a sign consistent with the NP phase implied by the DØ dimuon measurement. Such a NP phase also correlates with the expectation of a shift in $\sin 2\,\beta$ extracted from $B_s \to J/\psi\,\phi$ compared to the SM expectation [7, 8] and extractions from measurements in $B_s \to J/\psi\,K_s$. Such a consistent deviation is also observed, its statistical significance is $2.1\,\sigma$. Also see [9] for a discussion on the evidence for a NP phase in B_d and B_s meson mixing.

a framework for explaining the anomalies mentioned above without giving rise to flavor changing neutral current effects that are in conflict with experiment.

In this paper we discuss scalar models with MFV that can explain these anomalies. We first review how tree level exchanges of a neutral complex scalar in a simple two scalar doublet model can lead to enhanced CP violation in the B_q meson system and discuss the phenomenology of this model. We then show that limits on electric dipole moments suggest that the coupling of the charged scalar to the right handed u-type quarks is suppressed while its coupling to the d-type right handed quarks must be enhanced to be consistent with the data. We construct an extension of the MFV two scalar doublet model where this occurs naturally.[2]

2 Set up

We will utilize the recent fit of [2] to determine the new contribution to $B_q - \bar{B}_q$ mixing (here $q = s, d$). This fit is consistent in its conclusions with an earlier analysis [10]. The DØ result (a^b_{SL}) and the SM prediction [2] (A^b_{SL}) are given by

$$a^b_{SL} = \frac{N_b^{++} - N_b^{--}}{N_b^{++} + N_b^{--}},$$
$$= -(9.57 \pm 2.51 \pm 1.46) \times 10^{-3}, \tag{2.1}$$
$$A^b_{SL} = (-3.10^{+0.83}_{-0.98}) \times 10^{-4}. \tag{2.2}$$

where the number of $X\, b\, \bar{b} \to \mu^+\, \mu^+\, Y$ events is given by N_b^{++} for example. The quoted a^b_{SL} is a combination of the the asymmetry in each B_q, denoted a^{bq}_{SL}. Each of these contributions to a^b_{SL} can be expressed in terms of the mass and width differences (M_{12}, Γ_{12}) of the B_q meson eigenstates and the CP phase difference between these quantities ϕ_q as

$$a^{bq}_{SL} = \frac{|\Gamma^q_{12}|}{|M^q_{12}|} \sin\phi_q. \tag{2.3}$$

Naively one can effect the SM prediction through modifying M_{12} or Γ_{12} and both approaches have been explored in the literature. Modifying the decay width significantly [16, 17] as an explanation is problematic[3] and we will focus on MFV NP explanations that involve a NP contribution (that includes a new CP violating phase) to M^q_{12}.

The effect of NP on B_s and B_d mass mixing can be parametrized by two real parameters, $h_q > 0$ and σ_q by writing

$$M^q_{12} = (M^q_{12})^{\mathrm{SM}} + (M^q_{12})^{\mathrm{NP}}\,, \tag{2.4}$$

[2] Of course models with scalar doublets that are not supersymmetric suffer from the well know naturalness problem of keeping the doublets light compared to the Planck scale.

[3] The decay width can be removed in the relation between measured quantities under the assumption of small CP violation in NP induced tree level decays of B_q [10, 18] and the anomalous measurements can still be fit to finding $\sim 3\,\sigma$ evidence for a deviation from the SM [2, 10]. Also, an explanation of the like sign dimuon asymmetry through a NP contribution to $|\Gamma^q_{12}|$ would not necessarily explain the anomalies in $B_s \to J/\Psi\,\phi$ and $B \to \tau\,\nu$ with the correct correlation.

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where the new physics contribution to the mass mixing is related to the standard model value of the mass mixing by,

$$(M_{12}^q)^{\rm NP} = (M_{12}^q)^{\rm SM}\, h_q\, e^{2\, i\, \sigma_q}. \tag{2.5}$$

The models we discuss have $h_s = h_d$ and $\sigma_s = \sigma_d$ which is generally expected in NP models that obey MFV.[4] This scenario is argued to be a better fit to the current data then the SM in [2], which is disfavoured with a p-value of 3.1σ. In this case, the best fit values are $h_q = 0.255$ and $2\sigma_q = 180^o + 63.4^o$. The best fit magnitude of the correction h_q is small but its phase is large.

For simplicity in this paper we treat perturbative QCD in the leading logarithmic approximation and evaluate the needed matrix elements of four quark operators using the vacuum insertion approximation at the bottom mass scale. At the t-quark mass scale, in the SM, the effective Hamiltonian for $B_q - \bar{B}_q$ mixing is,

$$\mathcal{H}_q^{\rm SM} = (V_{tq}^\star\, V_{tb})^2 C^{\rm SM}(m_t) \bar{b}_L^\alpha \gamma^\mu q_L^\alpha \bar{b}_L^\beta \gamma_\mu q_L^\beta, \tag{2.6}$$

where α and β are color indices and

$$C^{\rm SM}(m_t) = \frac{G_F^2}{4\, \pi^2}\, M_W^2\, S(m_t^2/M_W^2). \tag{2.7}$$

Here $S(m_t^2/M_W^2) \simeq 2.35$ is a function of m_t^2/M_W^2 that results from integrating out the top quark and W-bosons. Using

$$(M_{12}^q)^{\rm SM} = \frac{\langle B_q | \mathcal{H}_q^{\rm SM} | \bar{B}_q \rangle}{2 m_{B_q}}, \tag{2.8}$$

we have after running down to the b-quark mass scale that,

$$(M_{12}^q)^{\rm SM} = (V_{tq}^\star\, V_{tb})^2 C^{\rm SM}(m_t) \left(\frac{1}{3}\right) \eta f_{B_q}^2 m_{B_q}. \tag{2.9}$$

Here $\eta \simeq 0.84$ is a QCD correction factor, $C_q^{\rm SM}(m_b) = \eta\, C_q^{\rm SM}(m_t)$.

The models for new physics we discuss generate the effective Hamiltonian at the top scale[5]

$$\mathcal{H}_q^{\rm NP} \simeq (V_{tq}^\star\, V_{tb})^2 C^{\rm NP}(m_t) \bar{b}_R^\alpha q_L^\alpha \bar{b}_R^\beta q_L^\beta. \tag{2.10}$$

[4]It has been proven in [27] that new CP violating effects can be larger in B_s than in B_d in nonlinear MFV. This observation has recently been explored in a general operator analysis [12] which showed that enhancements of CP violation in B_s mixing over B_d mixing by m_s/m_d requires contributions in the MFV expansion out to forth order in both the up and down Yukawas for operators induced by scalar exchange. The results on the neutron EDM using naive dimensional analysis (NDA) on page 7 disfavour order one down and up Yukawas with CP violating phases for these operators, so such contributions are expected to be very small. For alternative estimates of the relevant matrix element not using NDA see [19].

[5]For QCD running we don't distinguish between the top scale, weak scale and the mass scale of the new scalars we shall add.

Running down from m_b operator mixing induces the analogous operator with color indices rearranged. However, its coefficient is very small and we neglect it resulting in the relation, $C^{\rm NP}(m_b) \simeq \eta' C^{\rm NP}(m_t)$, where $\eta' \simeq 1.45$ [20]. Again using the vacuum insertion approximation at the b-quark mass scale we arrive at,

$$(M_{12}^q)^{\rm NP} \simeq (V_{tq}^\star V_{tb})^2 C^{\rm NP}(m_t) \left(-\frac{5}{24}\right) \eta' f_{B_q}^2 m_{B_q}. \tag{2.11}$$

Comparing with eq. (2.5)

$$h_q \, e^{2\,i\,\sigma_q} \simeq -\frac{5}{8} \left(\frac{C^{\rm NP}(m_t)}{C^{\rm SM}(m_t)}\right) \frac{\eta'}{\eta}. \tag{2.12}$$

3 Minimal two scalar doublet model

We now discuss how the minimal two scalar doublet model with MFV can have enhanced CP violation in B_q mixing due to tree level exchange of neutral scalars.[6] We denote by H the doublet that gets a vacuum expectation value and by S the doublet that does not. The Lagrangian in the Yukawa sector is

$$\mathcal{L}_Y = \bar{u}_R^i \, {g_U}^j{}_i \, Q_{Lj} H + \bar{d}_R^i \, {g_D}^j{}_i \, Q_{Lj} H^\dagger + \bar{u}_R^i \, {Y_U}{}^j_i \, Q_{Lj} S + \bar{d}_R^i \, {Y_D}{}^j_i \, Q_{Lj} S^\dagger + \text{h.c.} \tag{3.1}$$

where flavour indicies i, j are shown and color and $\mathrm{SU}(2)_L$ indices have been suppressed. MFV asserts that any NP also has the quark flavour symmetry group only broken by insertions proportional to Yukawa matrices so that ${Y_U}{}^j_i, {Y_D}{}^j_i$ are proportional to ${g_U}^j{}_i, {g_D}^j{}_i$. One can construct allowed NP terms by treating the Yukawa matrices as spurion fields that transform under flavour rotations as,

$$g_U \to V_U \, g_U \, V_Q^\dagger, \qquad g_D \to V_D \, g_D \, V_Q^\dagger, \tag{3.2}$$

where V_U is an element of $\mathrm{SU}(3)_{U_R}$, V_D is an element of $\mathrm{SU}(3)_{D_R}$, and V_Q is an element of $\mathrm{SU}(3)_{Q_L}$, i.e., the Yukawa matrices transform as $g_U \sim (\mathbf{3}, \mathbf{1}, \bar{\mathbf{3}})$ and $g_D \sim (\mathbf{1}, \mathbf{3}, \bar{\mathbf{3}})$ under the flavour group. MFV can be formulated up to linear order in top Yukawa insertions, or extended to a nonlinear representation of the symmetry [26, 27]. For enhanced CP violation in B_q mixing we are interested in a nonlinear realization of MFV. It is sufficient to only expand to next order in insertions of g_U so that

$$\begin{aligned} {Y_U}{}^j_i &= \eta_U \, {g_U}^j{}_i + \eta'_U \, {g_U}^j{}_k [(g_U^\dagger)^k{}_l \, (g_U)^l{}_i] + \cdots, \\ {Y_D}{}^j_i &= \eta_D \, {g_D}^j{}_i + \eta'_D \, {g_D}^j{}_k [(g_U^\dagger)^k{}_l \, (g_U)^l{}_i] + \cdots. \end{aligned} \tag{3.3}$$

We decompose the second scalar doublet as

$$S = \begin{pmatrix} S^+ \\ S^0 \end{pmatrix}, \tag{3.4}$$

where $S^0 = (S_R^0 + i S_I^0)/\sqrt{2}$.

[6]Previous analyses focused on scalar exchange to explain the like sign dimuon asymmetry include [21–23]. Also see [24, 25] for some phenomenological studies of models of this form.

The scalar potential is

$$
\begin{aligned}
V = &\frac{\lambda}{4}\left(H^{\dagger\, i}\, H_i - \frac{v^2}{2}\right)^2 + m_1^2\,(S^{\dagger i}\, S_i), \\
&+(m_2^2\, H^{\dagger\, i} S_i + \text{h.c.}) + \lambda_1\,(H^{\dagger\, i} H_i)\,(S^{\dagger\, j} S_j), \\
&+\lambda_2\,(H^{\dagger i}\, H_j)\,(S^{\dagger j}\, S_i) + \left[\lambda_3 H^{\dagger i}\, H^{\dagger j}\, S_i\, S_j + \text{h.c.}\right], \\
&+\left[\lambda_4 H^{\dagger i}\, S^{\dagger j}\, S_i\, S_j + \lambda_5 S^{\dagger i}\, H^{\dagger j}\, H_i\, H_j + \text{h.c.}\right] + \lambda_6 (S^{\dagger i} S_i)^2.
\end{aligned}
\tag{3.5}
$$

where i, j are SU(2) indices. Here $v \simeq 246\text{GeV}$ is the vacuum expectation value (vev) of the Higgs. Since we adopted the convention that the doublet S does not get a vev the parameters m_2^2 and λ_5 are related by,

$$
m_2^2 + \lambda_5^\star \frac{v^2}{2} = 0. \tag{3.6}
$$

The spectrum of neutral real scalars consists of the Higgs scalar h and S_R^0 and S_I^0. However, these are not mass eigenstates. In the (h, S_R^0, S_I^0) basis, the neutral mass squared matrix $\mathcal{M}^2$, where λ_3 is chosen real and positive, is

$$
\mathcal{M}^2 = \begin{pmatrix} m_h^2 & \lambda_5^R v^2 & \lambda_5^I v^2 \\ \lambda_5^R v^2 & m_S^2 + \lambda_3 v^2 & 0 \\ \lambda_5^I v^2 & 0 & m_S^2 - \lambda_3^2 \end{pmatrix}. \tag{3.7}
$$

Where $m_S^2 = m_1^2 + (\lambda_1 + \lambda_2) v^2/2$. Within the convention that λ_3 is real, the couplings η_U, η_U', η_D and η_D' and $\lambda_5 = \lambda_R^R + i\lambda_5^I$ are in general complex. The mass eigenstate scalars N_j^0 with mass m_j are related to h, S_R^0, S_I^0 by the orthogonal transformations

$$
\begin{aligned}
h &= \sum_j O_{hj}\, N_j, \\
S_R &= \sum_j O_{Rj}\, N_j, \qquad S_I = \sum_j O_{Ij}\, N_j.
\end{aligned}
\tag{3.8}
$$

We find the CP violating NP contribution to $B_q - \bar{B}_q$ mixing from neutral scalar exchange is

$$
C^{\text{NP}}(m_t) = \left(\sqrt{2}\,\eta_D'\, m_b/v\right)^2 \left(F\left(\sqrt{2} m_t/v\right)\right)^2 \frac{\Delta}{2}, \tag{3.9}
$$

where $F(x) = x^2 + \ \ldots$, and

$$
\Delta = \sum_j \frac{(O_{Rj} + i\, O_{Ij})^2}{m_j^2}. \tag{3.10}
$$

For the rest of this paper we truncate the expansions in $\sqrt{2} m_t/v$ at the leading non trivial term. So, for example, in eq. (3.11) we use $F(x) = x^2$.

For simplicity we now focus on the case where $\lambda_5 = 0$. Then h, S_R^0 and S_I^0 are mass eigenstates and the term in the potential proportional to λ_3 is of interest as it leads to the

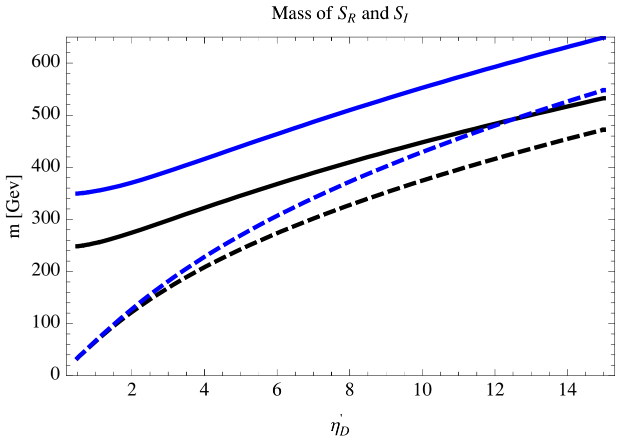

Figure 1. Mass of S_R (solid) and S_I (dashed) as a function of η'_D for fixed λ_3. The upper (blue curves) are for $\lambda_3 = 1$ the lower (black) curves are for $\lambda_3 = 0.5$.

mass splitting between the real neutral fields given by $m_R^2 - m_I^2 = 2\lambda_3\, v^2$. When this term is non vanishing,[7] the tree level exchange of $S_{R/I}$ generates a CP violating NP contribution to $B_q - \bar{B}_q$ mixing and

$$C^{\rm NP}(m_t) = (\eta'_D)^2 \left(\frac{\sqrt{2}\, m_t}{v}\right)^4 \left(\frac{\lambda_3\, m_b^2}{m_S^4 - \lambda_3^2\, v^4}\right). \tag{3.11}$$

In the above equation, the bottom quark mass $m_b \simeq 2.93$ GeV is evaluated at the top quark mass scale and $m^2_{S_{R/I}} = m_S^2 \pm \lambda_3 v^2$.

Using eq. (2.12) the mass scale of the new scalars is given by

$$m_S^4 \simeq \frac{20\,\pi^2\,\lambda_3\,|\eta'_D|^2 \eta' m_b^2\, m_t^4}{h_q\,\eta\, M_W^2\, S(m_t^2/m_W^2)} + \lambda_3^2 v^4. \tag{3.12}$$

Using the best fit value $h_q = 0.255$ [2] we find that

$$m_S^4 \simeq (154\,\text{GeV})^4\,|\eta'_D|^2\lambda_3 + (246\,\text{GeV})^4\,\lambda_3^2. \tag{3.13}$$

Then for example with a value $|\eta'_D| = 5$ and $\lambda_3 = 1$ the scalar mass scale is $m_S \simeq 360\,$GeV. As the mass splitting is significant we show in figure 1 the masses of the neutral scalars S_R, S_I as a function of η'_D. Moderate enhancements of η'_D avoid a light neutral state.

We have checked that the mass scale m_S required for $B_q - \mathrm{B}_q$ mixing is compatible with the constraints from $K - \bar{K}$ mixing. This compatibility is due to MFV, which causes the ratio of the relevant Wilson coefficients to scale as $m_s^2/m_b^2\,|V_{td}/V_{tb}|^2$.

Next we derive constraints on $\eta_U\eta_D$ that come from one loop Feynman diagrams with charged S scalar exchange. We will show that limits on electric dipole moments imply that $|{\rm Im}[\eta_U\eta_D]| \lesssim 10^{-1}$. Note that writing this as a constraint just on $\eta_U\eta_D$ depends on truncating a function of $\sqrt{2}m_t/v$ at leading order. We also examine the constraint on ${\rm Re}[\eta_U\eta_D]$ coming from experimental data on weak radiative B decay.

[7] Note that imposing custodial symmetry on the potential does not force $\lambda_3 \to 0$. Custodial symmetry violation is a measure of the total mass splitting $(m_R^2 - m_\pm^2)(m_I^2 - m_\pm^2) \propto (\lambda_2^2 - (2\lambda_3)^2)\, v^4$ in terms of the potential given in eq. (3.5).

3.0.1 Neutron electric dipole moment

The large CP violating phases needed in this two scalar doublet model contribute to other CP violating observables. Notable among them are electric dipole moments (EDM's). We will restrict our discussion here to the dominant contribution that is not suppressed by small quark masses when naive dimensional analysis (NDA) [28] is used. It comes about through the colour electric dipole moment of the b quark [29, 30] due to the effective Hamiltonian

$$\delta\mathcal{H}_{bg} = C_{gb}\, g_3\, m_b\, \bar{b}\, \sigma_{\mu\nu}\, T_a\, G^a_{\lambda\sigma}\, \epsilon^{\mu\nu\lambda\sigma}\, b, \tag{3.14}$$

inducing the dimension six CP violating operator

$$O_G = g_3^3\, f_{\alpha\beta\gamma}\, \epsilon^{\mu\nu\lambda\sigma}\, G_{\alpha\mu\rho}\, G^{\rho}_{\beta\nu}\, G_{\gamma\lambda\sigma}, \tag{3.15}$$

of Weinberg [31] when the b quark is integrated out. Our discussion will largely parallel the discussion of [32]. As S couples both to the up and down type quarks it induces a one loop contribution to the effective Hamiltonian above with

$$C_{gb}(m_S) = \frac{-\mathrm{Im}[\eta_U^\star\, \eta_D^\star]}{64\pi^2\, m_{S^\pm}^2} \left(\frac{\sqrt{2}\, m_t}{v}\right)^2 f(m_t^2/m_{S^\pm}^2), \tag{3.16}$$

where

$$f(x) = \frac{\log x}{(x-1)^3} + \frac{x-3}{2\,(x-1)^2}. \tag{3.17}$$

Running to $\mu \sim m_b$ using [29, 30] and estimating the matrix element of the operator with NDA[8] gives in e-cm units

$$d_n \sim 2\,\mathrm{Im}[\eta_U^\star\, \eta_D^\star]\, f(m_t^2/m_{S^\pm}^2) \left(\frac{1\,\mathrm{TeV}}{m_{S^\pm}}\right)^2 10^{-26}. \tag{3.18}$$

This is a significantly larger effect on EDM's than quoted in the general operator analysis [11] examining the effects of four Fermi operators on EDM's as this contribution is not suppressed by small mixing angles or light quark masses. For $m_{S^\pm} = 360$ GeV the neutron EDM experimental bound of $d_n < 2.9 \times 10^{-26}$ e-cm implies that $|\mathrm{Im}[\eta_U^* \eta_D^*]| < 0.26$ We plot the allowed $|\mathrm{Im}[\eta_U^* \eta_D^*]|$ as a function of mass for this NDA estimate in figure 2.

This suggests that $|\mathrm{Im}[\eta_U'^* \eta_D'^*]|$ (and the sum of the effect of all other cross terms such as $|\mathrm{Im}[\eta_U''^* \eta_D''^*]|$ etc.) is also small. However, given the uncertainties from hadronic matrix elements and given the fact that the parameters that enter the contribution to EDM's are not identical to those in $B_q - \bar{B}_q$ mixing it is difficult to draw precise conclusions on the parameters in the model that are important for mixing.

[8] We use method (a) of [32] with $\alpha_s(\mu = 1\,\mathrm{GeV}) \sim 4\,\pi$.

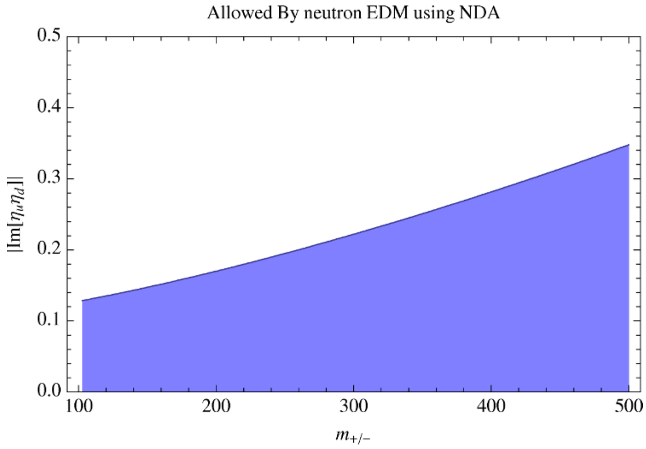

Figure 2. Allowed $|\mathrm{Im}[\eta_U^* \eta_D^*]|$ as a function of the charged scalar mass.

3.0.2 $B \to X_s\,\gamma$ constraints

Of course the two scalar doublet model also gives new contributions to quantities that are not CP violating. Here we briefly review the constraints on this model from $B \to X_s\,\gamma$ with these assumptions. The extra term in the effective Hamiltonian arises from charged scalar exchange and has the form,

$$\delta\,\mathcal{H}_{\bar{B}\to X_s\,\gamma} = [V_{ts}^\star\,V_{tb}]C_\gamma\,\left(\frac{e\,m_b}{16\,\pi^2}\,\bar{s}_L\,\sigma_{\mu\,\nu}\,F^{\mu\,\nu}\,b_R\right), \tag{3.19}$$

where $e<0$ is the electric charge. The Wilson coefficient is given by

$$C_\gamma = \eta_U^\star\,\eta_D^\star\,\left(\frac{2\,m_t^2}{v^2}\right)\,\frac{f_\gamma(m_t^2/m_{S^\pm}^2)}{3\,m_{S^\pm}^2}, \tag{3.20}$$

with

$$f_\gamma(x) = \frac{1}{4}\left(\frac{1+2\,x\,\mathrm{log}x-x^2}{(1-x)^3}\right)-\left(\frac{1+\mathrm{log}x-x}{(1-x)^2}\right). \tag{3.21}$$

This operator's contribution to the measured branching fraction $\mathrm{BR}(\bar{B}\to X_s\,\gamma)_{E_\gamma>1.6\,\mathrm{GeV}}$ is known [33]

$$\frac{\mathrm{BR}(\bar{B}\to X_s\,\gamma)_{E_\gamma>1.6\,\mathrm{GeV}}}{10^{-4}} = 3.15\pm 0.23-4.0v^2\,C_\gamma.$$

This constraint includes the effect of running this operator down to the scale m_b. Comparing to the world experimental average [34] we obtain a 1σ bound on the parameters of the form

$$-\,0.17 < Re[\eta_U^\star\,\eta_D^\star]\,f_\gamma(m_t^2/m_{S^\pm}^2)\,\frac{m_t^2}{3\,m_{S^\pm}^2} < 0.07. \tag{3.22}$$

For $m_{S^\pm} = 360$ GeV we find $-1.7 < \mathrm{Re}[\eta_U^\star\eta_D^\star] < 0.7$. Although this constraint is weak it is interesting that EDM's constrain $\mathrm{Im}[\eta_U^\star\,\eta_D^\star]$ while $B\to X_s\,\gamma$ constrains $\mathrm{Re}[\eta_U^\star\,\eta_D^\star]$.

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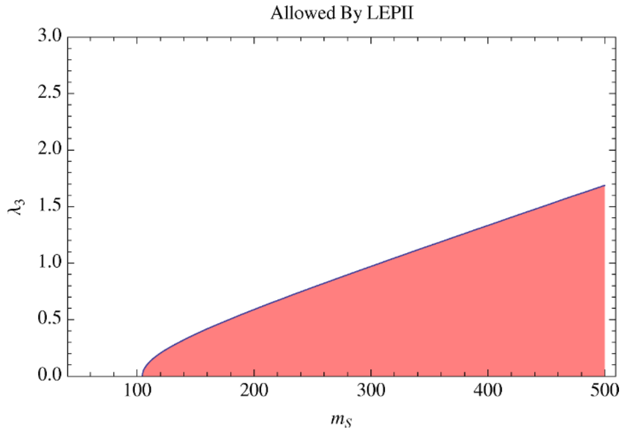

Figure 3. The allowed range of parameters m_S, λ_3 in the two scalar doublet model considering LEPII direct production bounds.

3.0.3 Collider physics: two scalar doublet model

Pairs of S particles can be produced through the tree level exchange of vector bosons produced through $q\,\bar{q}$ initial states in the case of the Tevatron and LHC and $e^+\,e^-$ in the case of LEPII.

From LEPII a bound on the mass scale of the new scalar doublet is obtained as no anomalous two and four jet events were seen when operating at $\sqrt{s} = 209\,\mathrm{GeV}$ where $0.1\ fb^{-1}$ of integrated luminosity was collected. The relevant cross sections in this case are given in [35] and the masses are bound to be

$$m_{S^\pm} \gtrsim 105\,\mathrm{GeV}, \quad m_{S^0_R} + m_{S^0_I} \gtrsim 209\,\mathrm{GeV}. \tag{3.23}$$

We plot the allowed m_S, λ_3 that satisfy this second bound for the minimal two scalar doublet model in figure 3. We have also performed an electroweak precision data fit. For scalar masses $\sim 100\,\mathrm{GeV}$ the constraints are weak. The allowed mass splitting in this model is $|m_I - m_\pm| \lesssim 200\,\mathrm{GeV}$ using the 95%CL region.

A light mass of S^0_I is allowed as these states must be produced in pairs through vector boson exchange and S^0_R can be heavy. However, a single neutral scalar particle can be produced at the Tevatron in association with a charged scalar though $W^\pm$ exchange. The partonic production cross section for producing $S^\pm\,S^0_I$ or $S^\pm\,S^0_R$ (when the width is neglected) is

$$\sigma = \frac{(p^2/s)^{3/2}}{s_W^4} \left(\frac{\pi\,\alpha_e^2(M_Z)}{6\,s}\right) \left|1 - \frac{M_W^2}{s}\right|^{-2}, \tag{3.24}$$

where p is the center of mass momentum of one of the produced particles and s is the partonic center of mass energy squared. We scan over the parameter space allowed by LEPII using this formula for the Tevatron production cross section (with MSTW 2008 PDF's [36]) where the $W^\pm$ is produced off the valence u, d quarks. The renormalization scale in what follows is always varied between $m_S/2$ and $2\,m_S$. The cross sections as a

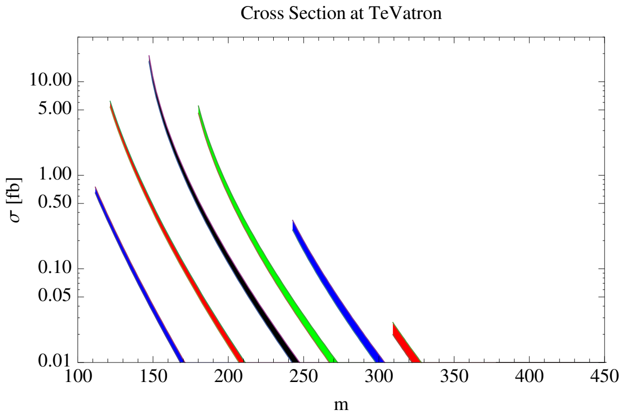

Figure 4. The cross section $\sigma(S^\pm S_I) + \sigma(S^\pm S_R)$ as a function of m_S for $\lambda_3 = (0.1, 0.2, 0.35, 0.5, 0.75, 1)$ going left to right. Here we have also imposed custodial symmetry on the potential $\lambda_2 = \pm 2\lambda_3$ for simplicity in the parameter scans.

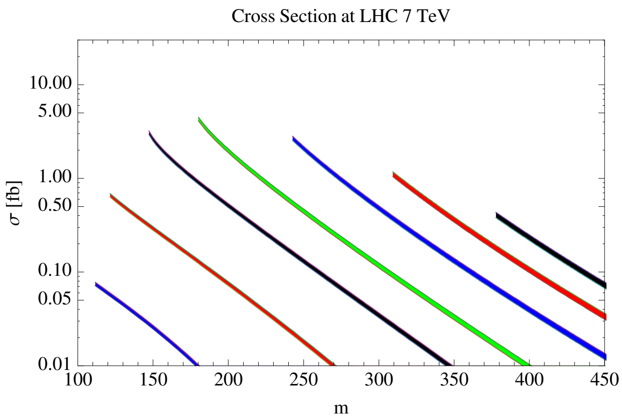

Figure 5. The cross section $\sigma(S^\pm S_I) + \sigma(S^\pm S_R)$ as a function of m_S for $\lambda_3 = (0.1, 0.2, 0.35, 0.5, 0.75, 1, 1.25)$ going left to right. Here we have also imposed custodial symmetry on the potential $\lambda_2 = \pm 2\lambda_3$ for simplicity in the parameter scans. The other production cross sections through $Z^\star, \gamma^\star$ are similar.

function of mass are shown in figure 4. The Tevatron can potentially constraint some of the allowed parameter space. Search strategies for pair production through weak boson fusion of charged $S^\pm$ particles that decay into $\bar{t}\, b\, \bar{b}\, t$ are also somewhat promising. In this case the production cross section for $m_S \sim 200\,\mathrm{GeV}$ is $\sigma \sim 1$ fb with a signal of two b jets and two t jets is produced in association with tagging light quark jets at large $\mathrm{p_T}$.

At the LHC, production through the tree level exchange of a vector boson is no longer dominated by $W^\pm$ exchange. The cross sections for the pair production of scalars are all similar in their dependence on λ_3 and as a function of m_S. We show $\sigma(pp \to W^\pm \to S^\pm S_{R/I})$ for $\sqrt{s} = 7\,\mathrm{TeV}$ in figure 5.

Although these vector boson exchange cross sections for LHC are small, potentially observable signals at LHC do exist when η_D is larger than one, and $S_{R/I}$ is made with large logarithms associated with collinear gluon splitting [37] and small p_T of the spectator

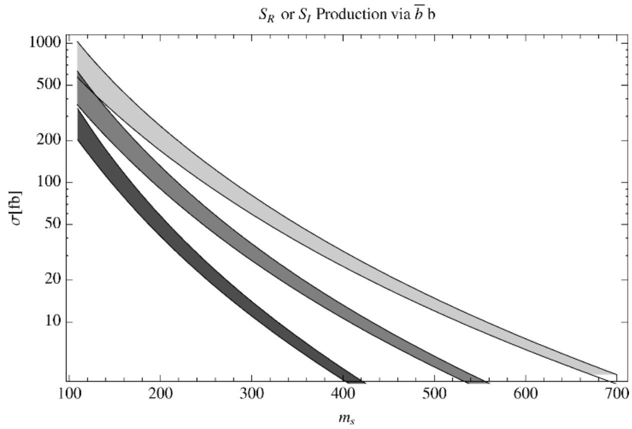

Figure 6. The cross section $\sigma(p\,p \to b\,\bar{b}\,S_I^0)$ for S_I^0 produced through collinear gluon splitting and b quark fusion. Shown are the cross sections for $\sqrt{s} = 7, 10, 14\,\text{TeV}$.

b quarks. The cross-section for the production of the lightest state $b\,\bar{b} \to S_I^0$ at leading log takes the form [38]

$$\sigma(b\bar{b}S_I^0) \simeq \frac{|\eta_D|^2\pi}{3\,s}\left(\frac{m_b^2}{v^2}\right)\int_{\frac{m_{S_I}^2}{s}}^1 \frac{dx}{x} b(x,\mu)\bar{b}\left(\frac{m_{S_I^0}^2}{xs},\mu\right), \tag{3.25}$$

where $b(x,\mu)$ and $\bar{b}(x,\mu)$ are the b quark and antiquark PDFs respectively. The large logs from collinear gluon splitting are summed into the parton distribution functions by choosing $\mu \sim m_S$. When we let $\eta_D = \sqrt{m_S/(154\,\text{GeV})}$ and choose $\lambda_3 = 1$ the production cross sections for the LHC are given by figure 6. This production mechanism must compete with the large b production background from QCD. However, we note that this signal has a distinct feature in its reconstruction of a resonance in the highest p_T b quark pair with a larger percentage of its total number of events at high p_T and small rapidity than the SM background, which has an approximate Rutherford scattering angular dependence in its production of b quarks.

4 A model with η_U naturally small

The charged scalar in the two Higgs doublet model has couplings to the quarks that (at leading order in the Yukawa matrices) are given by,

$$\mathcal{L}_{\text{charged}} = \eta_U \bar{u}_R g_U d_L S^+ + \eta_D \bar{d}_R g_D u_L S^- + \text{h.c.} \tag{4.1}$$

We need large CP violating phases to get the fit value of σ_q. For large phases the limits on electric dipole moments suggest that $|\eta_D\eta_U| \lesssim 10^{-1}$ for charged scalars with mass of a few hundred GeV. Unless the charged scalars are considerably heavier than the neutral ones this bound is expected to hold in the minimal two scalar doublet model if the model is to give the central value for h_q. If the limits on the electric dipole moments improve then this may become a more serious constraint.

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There does not seem to be any acceptable symmetry reason that this product is small and at the same time η_D is enhanced. In order to see that this is the case it is convenient to work in the basis where both H and S get a vevs v_H and v_S respectively. These can be chosen to be real. Then the charged scalar P^+ is the linear combination of the fields h^+ and S^+,

$$P^+ = \frac{v_S h^+ - v_H S^+}{\sqrt{v_H^2 + v_S^2}} \tag{4.2}$$

Because H no longer plays a special role we write the couplings of the scalars to the quarks as,

$$\mathcal{L}_Y = \epsilon_H \bar{u} \tilde{g}_U Q_L H + \epsilon_S \bar{u}_R \tilde{g}_U Q_L S + \epsilon'_H \bar{d}_R \tilde{g}_D Q_L H^\dagger + \epsilon'_S \bar{d}_R \tilde{g}_D Q_L S^\dagger + \text{h.c.} \tag{4.3}$$

Here we are using MFV and taking the quantities that break the flavor symmetry to be $\tilde{g}_{U/D}$. These matrices are proportional to the usual Yukawa matrices,

$$g_U = \left(\frac{\epsilon_H v_H + \epsilon_S v_S}{\sqrt{v_H^2 + v_S^2}} \right) \tilde{g}_U, \qquad g_D = \left(\frac{\epsilon'_H v_H + \epsilon'_S v_S}{\sqrt{v_H^2 + v_S^2}} \right) \tilde{g}_D. \tag{4.4}$$

Writing the charged scalar interaction as

$$\mathcal{L}_{\text{charged}} = \eta_U \bar{u}_R g_U d_L P^+ + \eta_D \bar{d}_R g_D u_L P^- + \text{h.c.} \tag{4.5}$$

we find that,

$$\eta_U = \frac{\epsilon_H v_S - \epsilon_S v_H}{\epsilon_H v_H + \epsilon_S v_S}, \quad \eta_D = \frac{\epsilon'_H v_S - \epsilon'_S v_H}{\epsilon'_H v_H + \epsilon'_S v_S}. \tag{4.6}$$

One way to get η_U small while η_D is large is to have $v_H \gg v_S$ so that $\eta_U \sim -\epsilon_S/\epsilon_H$ and $\eta_D \sim \epsilon'_H/\epsilon'_S$ and take the corresponding ratios of ϵ's to be small and large respectively. For their product to be small we also need, $\epsilon_S \epsilon'_H / \epsilon'_S \epsilon_H$ to be small. This is clearly possible, however there doesn't appear to be any symmetry reason behind these choices.

Note that there is an interchange symmetry where $H \leftrightarrow S$ that forces both $\eta_U = \eta_D = 0$. But in the limit of that symmetry, $P^+ = (h^+ - S^+)/\sqrt{2}$, and the symmetry's action on P^+ is $P^+ \to -P^+$. Hence there is a stable charged scalar.

The Glashow-Weinberg model [39] where H couples to the u-type quarks and S couples to the d-type quarks has, $\epsilon'_H = \epsilon_S = 0$ and so $\eta_U = v_S/v_H$ and $\eta_D = -v_H/v_S$. In this model $\eta_D \eta_U = -1$, so when η_U is small η_D is large, but their product cannot be made small even with a tuning of parameters.

In this section we construct a MFV model that has η_U small for a symmetry reason. New scalars that transform under flavour [40] can naturally have a small η_U. Consider a scalar field S_8 that transforms the same way as the Higgs doublet under the gauge group but as $(\mathbf{1}, \mathbf{8}, \mathbf{1})$ under the flavor group,

$$S_8 \to V_D \, S_8 \, V_D^\dagger. \tag{4.7}$$

We choose to represent the scalar in terms of the Gell-Mann matrices $S_8 = S_8^a\, T^a$ where $a = 1, \ldots, 8$ is a flavour index. The Yukawa couplings are given by

$$\mathcal{L}_Y = \bar{u}_R^i\, \hat{Y}_U{}^l{}_i\, (g_D^\dagger)^o{}_l\, (T^a)^n{}_o\, (g_D)^j{}_n\, Q_{Lj}\, S_8^a + \bar{d}_R^i\, (T^a)^m{}_i\, (\hat{Y}_D)^j{}_m Q_{Lj}\, S_8^{\dagger\, a} + \text{h.c.} \tag{4.8}$$

where we have made the flavour indices explicit. We use hat superscripts to distinguish this model's parameters from the two scalar doublet model. Recall that,

$$\begin{aligned} \hat{Y}_U{}^j{}_i &= \hat{\eta}_U\, g_U{}^j{}_i + \hat{\eta}'_U\, g_U{}^j{}_k [(g_U^\dagger)^k{}_l\; (g_U)^l{}_i] + \cdots\,, \\ \hat{Y}_D{}^j{}_i &= \hat{\eta}_D\, g_D{}^j{}_i + \hat{\eta}'_D\, g_D{}^j{}_k [(g_U^\dagger)^k{}_l\; (g_U)^l{}_i] + \cdots\,. \end{aligned} \tag{4.9}$$

The potential is given by

$$\begin{aligned} V = {} & \frac{\lambda}{4}\left(H^{\dagger i} H_i - \frac{v^2}{2}\right)^2 + 2\hat{m}_1^2\, \text{Tr}[S_8^{\dagger\, i}\, S_{8\, i}] \\ & + \hat{\lambda}_1\, H^{\dagger i}\, H_i\, \text{Tr}[S_8^{\dagger j}\, S_{8\, j}] + \hat{\lambda}_2\, H^{\dagger i}\, H_j\, \text{Tr}[S_8^{\dagger j}\, S_{8\, i}] \\ & + \Big[\hat{\lambda}_3\, H^{\dagger i}\, H^{\dagger j}\, \text{Tr}[S_{8 i}^\dagger\, S_{8\, j}] + \hat{\lambda}_4\, H^{\dagger i} \text{Tr}[S_8^{\dagger\, j} S_{8\, j} S_{8\, i}] \\ & + \hat{\lambda}_5\, H^{\dagger i}\, \text{Tr}[S_8^{\dagger j} S_{8\, i} S_{8\, j}] + \text{h.c.}\Big] \\ & + \hat{\lambda}_6\, \text{Tr}[S_8^{\dagger i}\, S_{8\, i}\, S_8^{\dagger j}\, S_{8\, j}] + \hat{\lambda}_7\, \text{Tr}[S_8^{\dagger i}\, S_{8\, j}\, S_8^{\dagger j}\, S_{8\, i}] \\ & + \hat{\lambda}_8 \text{Tr}[S_8^{\dagger i} S_{8\, i}]\, \text{Tr}[S_8^{\dagger j} S_{8\, j}] + \hat{\lambda}_9 \text{Tr}[S_8^{\dagger i} S_{8\, j}] \text{Tr}[S_8^{\dagger j} S_{8\, i}] \\ & + \hat{\lambda}_{10} \text{Tr}[S_{8 i} S_{8\, j}]\, \text{Tr}[S_8^{\dagger i} S_8^{\dagger j}] + \hat{\lambda}_{11} \text{Tr}[S_{8 i} S_{8\, j}] \text{Tr}[S_8^{\dagger j} S_8^{\dagger i}]. \end{aligned} \tag{4.10}$$

In the potential the index is an SU(2) index and the trace is over the down flavour index. We again rotate the phase of S_8 (relative to H) so that the $\hat{\lambda}_3$ term is real, then the couplings and $\hat{\lambda}_{4,5}$ and the η's are in general complex. In the above potential there are no linear terms in S_8 after H gets its vacuum expectation value and so it is natural for it not to have a vev.

In the potential and the $\hat{Y}'s$ one can also insert arbitrary numbers of $g_D\, g_D^\dagger$ matrices between contractions of a down index. We work in the down basis so that $g_D = \text{diag}(\sqrt{2}m_d/v, \sqrt{2}m_s/v, \sqrt{2}m_b/v)$. The interactions in the potential do not change flavour and are suppressed by m_b^2/v^2 so we neglect them.

Keeping just the leading term in, $\sqrt{2}m_t/v$, the Wilson coefficient of the effective Hamiltonian as defined in eq. (2.10) is

$$C^{\text{NP}}(m_t) = \frac{(\hat{\eta}'_D)^2 \left(\sqrt{2}\, m_t/v\right)^4\, \hat{\lambda}_3\, m_b^2/6}{\hat{m}_S^4 - \hat{\lambda}_3^2\, v^4/4}. \tag{4.11}$$

where $\hat{m}_S^2 = \hat{m}_1^2 + \left(\hat{\lambda}_1 + \hat{\lambda}_2\right) v^2/4$. This leads to the mass bound

$$\hat{m}_S^2 \simeq (98\,\text{GeV})^4\, |\hat{\eta}'_D|^2 \hat{\lambda}_3 + (174\,\text{GeV})^4\, \hat{\lambda}_3^2. \tag{4.12}$$

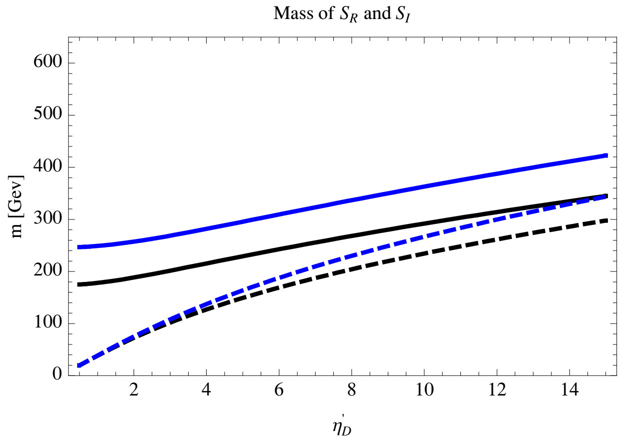

Figure 7. Mass of S_R (solid) and S_I (dashed) as a function of η_D' for fixed λ_3. The upper (blue curves) are for $\lambda_3 = 1$ the lower (black) curves are for $\lambda_3 = 0.5$.

in terms of the parameters defined in the potential. The mass spectrum of the new doublet is given by

$$\begin{aligned} m_{S^\pm}^2 &= \hat{m}_S^2 - \hat{\lambda}_2 \frac{v^2}{4}, \\ m_{S_R^0}^2 &= \hat{m}_S^2 + \hat{\lambda}_3 \frac{v^2}{2}, \\ m_{S_I^0}^2 &= \hat{m}_S^2 - \hat{\lambda}_3 \frac{v^2}{2}, \end{aligned} \tag{4.13}$$

We show the masses of the neutral scalars for this model in figure 7.

This model has eight new scalar doublets. Nevertheless, precision electroweak constraints are satisfied (when the Higgs is fixed to be $m_h = 96^{+29}_{-24}\,\mathrm{GeV}$ for $m_S \gtrsim 100\,\mathrm{GeV}$) when $|m_I - m_\pm| < 50\,\mathrm{GeV}$ [35]. Conversely, custodial SU(2) violation in such a light scalar doublet leading to a positive contribution to ΔT can raise the allowed mass of the Higgs in EWPD [35, 41].

In this model the coupling constant analogous to η_U is naturally of order $(m_b/v)^2 \sim 10^{-3}$ and the $\bar{B} \to X_S\,\gamma$ and neutron EDM effects of the model are suppressed as phenomenologically required due to MFV.

The collider phenomenology in this model is very similar to the discussion on the two scalar doublet model. The LEPII constraints allow a larger parameter space due to the smaller mass splitting. The main differences for the Tevatron is that the cross sections we have discussed are increased by an order of magnitude due to the larger flavour representation. Slightly smaller production cross sections through b quark fusion with low p_T spectator b quarks are expected at LHC as the normalization of the Gell Mann matrix decreases the cross section by a factor of three.

We consider speculation on the UV origin of such a S_8 doublet, or other accompanying non flavour singlet doublets that transform under the $\mathrm{SU(3)_{U_R}}$ as an **8** to be premature and beyond the scope of this work.

5 Conclusions

In this paper we discussed the new (i.e., beyond the minimal standard model) physics in the region of parameter space for which the two scalar doublet model with MFV gives the additional contributions to $B_q - \bar{B}_q$ mixing that are hinted at by the data on flavor physics in the B-sector. It requires additional light scalars that may be discovered at the Tevatron or LHC. Experimental limits on electric dipole moments suggest a region of parameter space that can occur naturally in some models where the new doublet of scalars transforms non-trivially under the flavour group. We constructed such a model.

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A EWPD calculations

The one loop results for the S_8 model are the same as for the model discussed in [32, 35]. We use the STUVWX parameterization [43] of EWPD as the mass scale of the new scalars is $\sim 100\,\text{GeV}$. The relevant results in terms of Passarino-Veltman functions [42] with standard definitions[9] are

$$\begin{aligned}
\delta\Pi_{WW}(p^2) &= \frac{g_1^2}{2\pi^2}\Bigg[B_{22}(p^2, m_I^2, m_+^2) + B_{22}(p^2, m_R^2, m_+^2), \\
&\quad -\frac{1}{2}A_0(m_+^2) - \frac{1}{4}A_0(m_R^2) - \frac{1}{4}A_0(m_I^2)\Bigg], \\
\delta\Pi_{ZZ}(p^2) &= \frac{g_1^2}{2\pi^2 c^2}\Bigg[(1-2s^2)^2\left(B_{22}(p^2, m_+^2, m_+^2) - \frac{1}{2}A_0(m_+^2)\right), \\
&\quad + B_{22}(p^2, m_R^2, m_I^2) - \frac{1}{4}A_0(m_R^2) - \frac{1}{4}A_0(m_I^2)\Bigg], \\
\delta\Pi_{\gamma\gamma}(p^2) &= \frac{2e^2}{\pi^2}\Bigg[B_{22}(p^2, m_+^2, m_+^2) - \frac{1}{2}A_0(m_+^2)\Bigg], \\
\delta\Pi_{\gamma Z}(p^2) &= \frac{e g_1(1-2s^2)}{\pi^2 c}\Bigg[B_{22}(p^2, m_+^2, m_+^2) - \frac{1}{2}A_0(m_+^2)\Bigg].
\end{aligned}$$

[9] With c, s the cosine and sine of the weak mixing angle.

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For $p^2 = 0$ these expressions become

$$\begin{aligned}
\delta\Pi_{WW}(0) &= \frac{g_1^2}{8\pi^2}\left(\frac{1}{2}f(m_+,m_R) + \frac{1}{2}f(m_+,m_I)\right),\\
\delta\Pi_{ZZ}(0) &= \frac{g_1^2}{8\pi^2 c^2}\left(\frac{1}{2}f(m_R,m_I)\right),
\end{aligned}$$

where

$$f(m_1,m_2) = m_1^2 + m_2^2 - \frac{2m_1^2 m_2^2}{m_1^2 - m_2^2}\log\frac{m_1^2}{m_2^2}.$$

The derivatives of the vacuum polarizations are

$$\begin{aligned}
\delta\Pi'_{\gamma\gamma}(0) &= -\frac{e^2}{6\pi^2}B_0(0,m_+^2,m_+^2),\\
\delta\Pi'_{\gamma Z}(0) &= -\frac{eg_1(1-2s^2)}{12\pi^2 c}B_0(0,m_+^2,m_+^2),\\
\delta\Pi'_{WW}(p^2) &= \frac{g_1^2}{2\pi^2}\left[-\frac{1}{6}\Delta + \frac{\partial b_{22}(p^2,m_I^2,m_+^2)}{\partial p^2} + \frac{\partial b_{22}(p^2,m_R^2,m_+^2)}{\partial p^2}\right]\\
\delta\Pi'_{ZZ}(p^2) &= \frac{g_1^2}{2\pi^2 c^2}\left[-\frac{1}{12}\Delta + \frac{\partial b_{22}(p^2,m_R^2,m_I^2)}{\partial p^2},\right.\\
&\qquad \left.(1-2s^2)^2\left(-\frac{1}{12}\Delta + \frac{\partial b_{22}(p^2,m_+^2,m_+^2)}{\partial p^2}\right)\right].
\end{aligned}$$

Using these results we can construct the STUVWX parameters with the standard definitions [43]

$$\begin{aligned}
\frac{\alpha S}{4s^2\,c^2} &= \left[\frac{\delta\Pi_{ZZ}(M_Z^2) - \delta\Pi_{ZZ}(0)}{M_Z^2}\right] - \frac{(c^2-s^2)}{s\,c}\delta\Pi'_{Z\gamma}(0) - \delta\Pi'_{\gamma\gamma}(0),\\
\alpha T &= \frac{\delta\Pi_{WW}(0)}{M_W^2} - \frac{\delta\Pi_{ZZ}(0)}{M_Z^2},\\
\frac{\alpha U}{4s^2} &= \left[\frac{\delta\Pi_{WW}(M_W^2) - \delta\Pi_{WW}(0)}{M_W^2}\right] - c^2\left[\frac{\delta\Pi_{ZZ}(M_Z^2) - \delta\Pi_{ZZ}(0)}{M_Z^2}\right]\\
&\quad -s^2\,\delta\Pi'_{\gamma\gamma}(0) - 2\,s\,c\,\delta\Pi'_{Z\gamma}(0),\\
\alpha V &= \delta\Pi'_{ZZ}(M_Z^2) - \left[\frac{\delta\Pi_{ZZ}(M_Z^2) - \delta\Pi_{ZZ}(0)}{M_Z^2}\right],\\
\alpha W &= \delta\Pi'_{WW}(M_W^2) - \left[\frac{\delta\Pi_{WW}(M_W^2) - \delta\Pi_{ZZ}(0)}{M_W^2}\right],\\
\alpha X &= -s\,c\left[\frac{\delta\Pi_{Z\gamma}(M_Z^2)}{M_Z^2} - \delta\Pi'_{Z\gamma}(0)\right]
\end{aligned}$$

Here Δ is the divergence that cancels in the pseudo-observables STUVWX but we note we calculate in dimensional regularization and $\overline{\mathrm{MS}}$ in $d = 4-2\epsilon$ dimensions. As the number of degrees of freedom in this S_8 model and in the model [32] are the same, we can directly use the detailed fit results on the allowed masses (determined from these formulas) presented in [35]. These results allow masses for fixed $m_h = 96^{+29}_{-24}\,\mathrm{GeV}$ when $m_S \gtrsim 100\,\mathrm{GeV}$ characterized by $|m_I - m_\pm| < 50\,\mathrm{GeV}$ for S_8.

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