

***B*-physics constraints on baryon number violating couplings: Grand unification, supersymmetry, or compositeness**

Debrupa Chakraverty* and Debajyoti Choudhury†

Mehta Research Institute, Chhatnag Road, Jhusi, Allahabad 211 019, India‡

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We investigate the role that baryon number violating interactions may play in B phenomenology. Present in various grand unified theories, supersymmetric theories with R -parity violation and composite models, a fundamental diquark state could be quite light. Using the data on B decays as well as B - \bar{B} mixing, we find strong constraints on the couplings that such a light diquark state may have with the standard model quarks. Our analysis is most general in nature, and the particular models mentioned above are only specific examples thereof.

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I. INTRODUCTION

It is now widely accepted that the standard model (SM), despite its great success, is only an effective theory. The ills plaguing it may be cured only in the context of a more fundamental theory operative at higher energies. The quest to find such a theory has, over the years, inspired many a model going beyond the SM. Three of the most attractive classes of such models are those incorporating grand unification [1], supersymmetry [2] or compositeness [3]. The exact nature of such a theory, however, is a matter of intense debate, occasioned, not in the least, by the absence of any experimental signature yet. It is not surprising, thus, that the search for new physics effects constitutes a major component of research in high energy physics. Such efforts can be broadly classified into two categories. On the one hand there are the direct searches typified by high energy collider experiments where new particles are sought to be produced on-shell and detected through their subsequent decays. The other approach concentrates on indirect effects as can be deduced from possible deviations from the SM predictions for low-energy observables. In this article we shall focus on one such set of low and intermediate energy experiments.

The next decade will see a blossoming of experimental facilities planning to explore B - \bar{B} mixing as well as B -meson decays with greater accuracy and for an increasing number of different final states. In light of these upcoming experiments (CLEO, BaBar, BELLE, HERAB, BTeV and LHCb), it is of importance to examine their sensitivity to new physics beyond the SM.

In this paper, we investigate the possible influence that a baryon number violating interaction may have on B phenomenology. Within the SM, baryon (\hat{B}) and lepton (\hat{L}) number conservations come about due to accidental symmetries. In other words, such conservations are not guaranteed by any principle, but are rather the consequences of the choice of the

particle content.¹ In extensions of the SM, such an accidental occurrence is obviously not guaranteed. For example, even in the simplest grand unified theories (GUTs), both the gauge and the scalar sector interactions violate each of \hat{B} and \hat{L} . The corresponding particles, namely the diquarks [4] and leptiquarks [5], have been studied in the literature to a considerable extent.

Simultaneous breaking of both \hat{B} and \hat{L} symmetry is obviously a recipe for disaster as this combination is more than likely to lead to rapid proton decay. Within GUTs, gauge boson-mediated proton decay could be naturally suppressed by postulating the symmetry breaking scale to be very large. However, there do exist a class of GUTs [6], where the next set of thresholds need not be very high and B -violating gauge particles can be relatively light. Proton decay, however, remains suppressed on account additional symmetries in the theory. Suppression of the scalar mediated contribution to proton decay in a generic GUT, on the other hand, is easier to obtain: the particle content can be so chosen that there is no diquark-leptoquark mixing, at least as far as the light sector is concerned.

In the case of the minimal supersymmetric standard model (MSSM), though, we do not have the option of demanding the ‘‘offending’’ fields (the supersymmetric partners of the SM fermions) to be superheavy. Ruling out the undesirable terms necessitates the introduction of a discrete symmetry, $R \equiv (-1)^{3(\hat{B}-\hat{L})+2\hat{S}}$ (with \hat{S} denoting the spin of the field) [7]. Apart from ruling out both \hat{B} and \hat{L} violating terms in the superpotential, this symmetry has the additional consequence of rendering the lightest supersymmetric partner absolutely stable. However, such a symmetry is *ad hoc*. Hence, it is of interest to consider possible violations of this symmetry especially since it has rather important experimental consequences, not the least of which concerns the detection of the supersymmetric partners.

It can thus be argued that, in such models as well as in models of compositeness [3], it is quite likely that baryon number violating interactions may not be suppressed too se-

*Email address: rupa@mri.ernet.in

†Email address: debchou@mri.ernet.in

‡Presently known as Harish-Chandra Research Institute.

¹Indeed, nonperturbative effects within the SM itself do break $\hat{B}+\hat{L}$ symmetry.

TABLE I. Gauge quantum numbers and Yukawa couplings of diquarks ($Q_{em}=T_3+Y/2$).

Diquark type	Coupling	$SU(3)_c \times SU(2)_L \times U(1)_Y$
Φ_1	$h_{ij}^{(1)}(\bar{Q}_{Li})^c Q_{Lj} \Phi_1$	$(\bar{6}, 3, -\frac{2}{3})$
Φ_2	$h_{ij}^{(2)}(\bar{Q}_{Li})^c Q_{Lj} \Phi_2$	$(3, 3, -\frac{2}{3})$
Φ_3	$[h_{ij}^{(3)}(\bar{Q}_{Li})^c Q_{Lj} + \tilde{h}_{ij}^{(3)}(\bar{u}_{Ri})^c d_{Rj}] \Phi_3$	$(\bar{6}, 1, -\frac{2}{3})$
Φ_4	$[h_{ij}^{(4)}(\bar{Q}_{Li})^c Q_{Lj} + \tilde{h}_{ij}^{(4)}(\bar{u}_{Ri})^c d_{Rj}] \Phi_4$	$(3, 1, -\frac{2}{3})$
Φ_5	$h_{ij}^{(5)}(\bar{u}_{Ri})^c u_{Rj} \Phi_5$	$(\bar{6}, 1, -\frac{8}{3})$
Φ_6	$h_{ij}^{(6)}(\bar{u}_{Ri})^c u_{Rj} \Phi_6$	$(3, 1, -\frac{8}{3})$
Φ_7	$h_{ij}^{(7)}(\bar{d}_{Ri})^c d_{Rj} \Phi_7$	$(\bar{6}, 1, \frac{4}{3})$
Φ_8	$h_{ij}^{(8)}(\bar{d}_{Ri})^c d_{Rj} \Phi_8$	$(3, 1, \frac{4}{3})$
V_1^μ	$\vartheta_{ij}^{(1)}(\bar{Q}_{Li})^c \gamma_\mu d_{Rj} V_1^\mu$	$(\bar{6}, 2, \frac{1}{3})$
V_2^μ	$\vartheta_{ij}^{(2)}(\bar{Q}_{Li})^c \gamma_\mu d_{Rj} V_2^\mu$	$(3, 2, \frac{1}{3})$
V_3^μ	$\vartheta_{ij}^{(3)}(\bar{Q}_{Li})^c \gamma_\mu u_{Rj} V_3^\mu$	$(\bar{6}, 2, -\frac{5}{3})$
V_4^μ	$\vartheta_{ij}^{(4)}(\bar{Q}_{Li})^c \gamma_\mu u_{Rj} V_4^\mu$	$(3, 2, -\frac{5}{3})$

verely. Even more interestingly, such processes may be mediated by relatively low-lying states, generically called diquarks.

This paper is organized as follows: Section II constitutes a brief review on diquarks. Section III deals with hadronic B decays. In Sec. IV, we concentrate on B - \bar{B} mixing. Section V contains the numerical results. We conclude in Sec. VI with a summary and outlook.

II. DIQUARKS: A BRIEF REVIEW

In this section we shall briefly examine all possible tree-level \hat{B} -violating couplings involving the SM quarks. We shall adopt a purely phenomenological standpoint without any particular reference or prejudice to the origin of such couplings or states. A generic diquark² is a scalar or vector particle that couples to a quark current with a net baryon number $\hat{B} = \pm 2/3$. Clearly, under $SU(3)_c$, it may transform as either a triplet or a sextet. For scalars, the Yukawa term in the Lagrangian can be expressed as

$$\mathcal{L}_{SD} = h_{ij}^{(A)} \bar{q}_i^c P_{L,R} q_j \Phi_A + \text{H.c.}, \quad (1a)$$

where i, j denote quark flavors, A denotes the diquark type and $P_{L,R}$ reflect the quark chirality. Standard model gauge invariance demands that a scalar diquark transforms either as a triplet or as a singlet under $SU(2)_L$. For a vector diquark, on the other hand, the relevant term in the Lagrangian can be parametrized as

$$\mathcal{L}_{VD} = \vartheta_{ij}^{(A)} \bar{q}_i^c \gamma_\mu P_{L,R} q_j V_A^\mu + \text{H.c.} \quad (1b)$$

with V_A transforming as a $SU(2)_L$ doublet. The full list of quantum numbers, for either case, is presented in Table I.

²These diquarks should not be confused with the two quark ‘‘bound’’ states that appear in discussion of baryons. Rather, to the scale that B -physics probes, these are essentially fundamental scalars or vectors.

Clearly, the couplings $h_{ij}^{(1)}$, $h_{ij}^{(4)}$, $h_{ij}^{(5)}$ and $h_{ij}^{(7)}$ must be symmetric under the exchange of i and j while $h_{ij}^{(2)}$, $h_{ij}^{(3)}$, $h_{ij}^{(6)}$ and $h_{ij}^{(8)}$ must be antisymmetric. For the other couplings, viz. $\tilde{h}_{ij}^{(3)}$, $\tilde{h}_{ij}^{(4)}$, and $\vartheta_{ij}^{(A)}$, no such symmetry property exists. Hereafter, we assume these couplings to be real.³ Note that the quantum numbers of $\Phi_{2,4,6}$ as well as those of $V_{2,4}^\mu$ allow them to couple to a leptoquark (i.e., a quark-lepton) current as well. Clearly, the non-observance of proton decay implies that such L -violating couplings must be suppressed severely.

It should be noted that we are not demanding that the vector diquarks correspond to some gauge theory. While it might be rightly argued that a theory with non-gauged vector particles is non-renormalizable, one should keep in mind that such states may well be there in an effective theory. Since we would be studying the phenomenological implications only at the lowest order of perturbation theory, renormalizability is not an issue here.

We now turn to the MSSM, where both \hat{B} - and \hat{L} -violating terms are allowed, in general, by supersymmetry as well as gauge invariance. As stated earlier, catastrophic rates for proton decay can be avoided by imposing a global Z_2 symmetry [7] under which the quark and lepton superfields change by a sign, while the Higgs superfields remain invariant. However, since such a symmetry is entirely *ad hoc* within the purview of the MSSM, it is conceivable that this R -parity may be broken while keeping either of \hat{B} or \hat{L} intact. In our study, we shall restrict ourselves to the case where only the \hat{B} -violating terms are non-zero. Such scenarios can be motivated from a class of supersymmetric GUTs as well [8]. The corresponding terms in the superpotential can be parametrized as

$$W_k = \lambda_{ijk}'' \bar{U}_R^i \bar{D}_R^j \bar{D}_R^k, \quad (2)$$

³The extension to complex couplings is straightforward. The imaginary parts, however, can be better constrained from an analysis of the CP violating decay modes.

TABLE II. The effective four quark operator for various diquarks. The operators for Φ_2 , Φ_4 , Φ_8 , V_2 and V_4 mirror those for Φ_1 , Φ_3 , Φ_7 , V_1 and V_3 respectively, albeit with a different color factor (see text). Here $L(R) = 1 \mp \gamma_5$.

Diquark type	Effective four quark operator
Φ_1	$\frac{h_{ij}^{(1)} h_{kl}^{(1)}}{16m_{\Phi_1}^2} [(\bar{u}_k \gamma_\mu L u_i)(\bar{d}_l \gamma^\mu L d_j) + (\bar{d}_k \gamma_\mu L d_i)(\bar{u}_l \gamma^\mu L u_j)$ $+ (\bar{d}_k \gamma_\mu L u_i)(\bar{u}_l \gamma^\mu L d_j) + (\bar{u}_k \gamma_\mu L d_i)(\bar{d}_l \gamma^\mu L u_j)$ $+ 2(\bar{d}_k \gamma_\mu L d_i)(\bar{d}_l \gamma^\mu L d_j)] + \text{H.c.}$
Φ_3	$\frac{h_{ij}^{(3)} h_{kl}^{(3)}}{8m_{\Phi_3}^2} [(\bar{u}_k \gamma_\mu L u_i)(\bar{d}_l \gamma^\mu L d_j) + (\bar{d}_k \gamma_\mu L d_i)(\bar{u}_l \gamma^\mu L u_j)$ $- (\bar{d}_k \gamma_\mu L u_i)(\bar{u}_l \gamma^\mu L d_j) - (\bar{u}_k \gamma_\mu L d_i)(\bar{d}_l \gamma^\mu L u_j)] + \text{H.c.}$ $\frac{\tilde{h}_{ij}^{(3)} \tilde{h}_{kl}^{(3)}}{8m_{\Phi_3}^2} (\bar{u}_k \gamma_\mu R u_i)(\bar{d}_l \gamma^\mu R d_j) + \text{H.c.}$ $\frac{\tilde{h}_{ij}^{(3)} h_{kl}^{(3)}}{8m_{\Phi_3}^2} [(\bar{d}_k R u_i)(\bar{u}_l R d_j) - (\bar{u}_k R u_i)(\bar{d}_l R d_j)$ $+ \frac{1}{4}\{(\bar{d}_k \sigma^{\mu\nu} R u_i)(\bar{u}_l \sigma_{\mu\nu} R d_j) - (\bar{u}_k \sigma^{\mu\nu} R u_i)(\bar{d}_l \sigma_{\mu\nu} R d_j)\}] + \text{H.c.}$
Φ_7	$\frac{h_{ij}^{(7)} h_{kl}^{(7)}}{8m_{\Phi_7}^2} (\bar{d}_k \gamma_\mu R d_i)(\bar{d}_l \gamma^\mu R d_j) + \text{H.c.}$
V_1	$\frac{-\vartheta_{ij}^{(1)} \vartheta_{kl}^{(1)}}{4m_{V_1}^2} [(\bar{u}_k \gamma_\mu L u_i) + (\bar{d}_k \gamma_\mu L d_i)](\bar{d}_l \gamma^\mu R d_j) + \text{H.c.}$
V_3	$\frac{-\vartheta_{ij}^{(3)} \vartheta_{kl}^{(3)}}{4m_{V_3}^2} (\bar{d}_k \gamma_\mu L d_i)(\bar{u}_l \gamma^\mu R u_j) + \text{H.c.}$

where \bar{U}_R^i and \bar{D}_R^i denote the right-handed up-quark and down-quark superfields, respectively. The couplings λ_{ijk}'' are antisymmetric under the exchange of the last two indices. The corresponding Lagrangian can then be written in terms of the component fields as

$$\mathcal{L}_k = \lambda_{ijk}'' (u_i^c d_j^c \tilde{d}_k^* + u_i^c \tilde{d}_j^* d_k^c + \tilde{u}_i^* d_j^c d_k^c) + \text{H.c.} \quad (3)$$

Thus, a single term in the superpotential corresponds to two of type $\tilde{h}_{ij}^{(4)}$ and one of type $h_{ij}^{(8)}$ diquark interactions.

The best direct bound on diquark type couplings is derived from an analysis of dijet events at the Tevatron [9]. Considering the process $q_i q_j \rightarrow \Phi_A \rightarrow q_i q_j$, an exclusion curve in the $(m_{\Phi_A}, h_{ij}^{(A)})$ plane can be obtained from this data. A similar statement holds for the vector particles as well. Two points need to be noted though. At a $p\bar{p}$ collider like the Tevatron, the uu and dd fluxes are small and hence the bounds are relatively weak. This is even more true for quarks of the second or third generation (which are relevant for the couplings that we are interested in). Secondly, such an analysis needs to make assumptions regarding the branching fraction of Φ_A (V_A) into quark pairs, a point that is of particular importance in the context of R -parity violating supersymmetric models.

There also exist some constraints derived from low energy processes. Third generation couplings, for example, can be constrained from the precision electroweak data at CERN e^+e^- collider LEP [10] or, to an extent, by demanding perturbative unitarity to a high scale [11]. Couplings involving the first two generations, on the other hand, are constrained⁴ by the non-observance of neutron-antineutron oscillations or from an analysis of rare nucleon and meson decays [12,13]. While many of these individual bounds are weak, certain of their *products* are much more severely constrained by the data on neutral meson mixing and CP violation in the K sector [14] as well as from radiative $b \rightarrow s \gamma$ decay [15]. It is our aim, in this article, to derive analogous but stronger bounds.

At energy scales well below the mass of the diquark, the latter can be integrated out and effective four quark operators obtained. In Table II, we list these for each diquark type. A few points should be noted here:

we have not displayed the operators resulting from Φ_5

⁴Although many of these analyses have been done for the case of R -parity violating models, clearly similar bounds would also apply to nonsupersymmetric diquark couplings as well.

and Φ_6 as these do not contribute (at the lowest order) to either B decays or B - \bar{B} mixing;

for convenience's sake, we have Fierz-rearranged the operators and, in the process, exchanged the charge-conjugated fermion fields (which come in naturally) for their non-conjugated counterparts;

within a diquark multiplet, we have assumed all the fields to be mass degenerate since large splittings within a multiplet are anyway disfavored by LEP data;

we have neglected the evolution of the diquark mediated effective four-quark interactions from the electroweak scale down to B meson scale through renormalization group equations;

we have not displayed the extra color factors that appear on account of the diquark states being colored objects. The said factors can be determined by reexpressing four quark operators of the forms $(\bar{3}_c \otimes 3_c)_1$ and $(6_c \otimes \bar{6}_c)_1$ in terms of the corresponding $(1_c \otimes 1_c)_1$ and $(8_c \otimes 8_c)_1$ current structures. Thus, transforming $(\bar{q}_i^c \Gamma q_j)(\bar{q}_l \Gamma' q_k^c)$ to the form $(\bar{q}_k \Gamma'' q_i)(\bar{q}_l \Gamma''' q_j)$ implies that we are dealing with linear combinations of the form

$$\begin{aligned} (\bar{3}_c \otimes 3_c)_1 &= \frac{2}{3}(1_c \otimes 1_c)_1 - (8_c \otimes 8_c)_1 \\ (6_c \otimes \bar{6}_c)_1 &= \frac{2}{3}(1_c \otimes 1_c)_1 + \frac{1}{2}(8_c \otimes 8_c)_1. \end{aligned} \quad (4)$$

These extra color factors need to be included while calculating the hadronic matrix elements.

Looking at Table II, it is obvious that the effective Hamiltonian for the full theory can be parametrized as

$$\mathcal{H}_{\text{eff}} = \sum_{i=0}^9 b_i \mathcal{H}_i \quad (5a)$$

with

$$\begin{aligned} \mathcal{H}_0 &= (\bar{q}_1 \gamma_\mu R b)(\bar{q}_2 \gamma^\mu R q_3) & \mathcal{H}_1 &= (\bar{q}_1 \gamma_\mu L b)(\bar{q}_2 \gamma^\mu L q_3) \\ \mathcal{H}_2 &= (\bar{q}_1 \gamma_\mu R b)(\bar{q}_2 \gamma^\mu L q_3) & \mathcal{H}_3 &= (\bar{q}_1 \gamma_\mu L b)(\bar{q}_2 \gamma^\mu R q_3) \\ \mathcal{H}_4 &= (\bar{q}_1 L b)(\bar{q}_2 R q_3) & \mathcal{H}_5 &= (\bar{q}_1 R b)(\bar{q}_2 L q_3) \\ \mathcal{H}_6 &= (\bar{q}_1 L b)(\bar{q}_2 L q_3) & \mathcal{H}_7 &= (\bar{q}_1 R b)(\bar{q}_2 R q_3) \\ \mathcal{H}_8 &= (\bar{q}_1 \sigma_{\mu\nu} L b)(\bar{q}_2 \sigma^{\mu\nu} L q_3) \\ \mathcal{H}_9 &= (\bar{q}_1 \sigma_{\mu\nu} R b)(\bar{q}_2 \sigma^{\mu\nu} R q_3). \end{aligned} \quad (5b)$$

The strengths b_i include both the SM contributions as well as diquark contributions (as in Table II) wherever applicable.

It now remains to calculate the hadronic matrix elements for \mathcal{H}_i , a task that is rendered very difficult by the associated strong interaction dynamics. Hence, instead of attempting an exact calculation, one normally takes recourse to some appropriate approximation. In the ‘naive factorization’ approach [16], the matrix elements of a four-quark operator are

TABLE III. Values of decay constants in MeV.

f_π	f_ρ	f_K	f_{K^*}	f_D	f_{D^*}	f_{D_s}	$f_{D_s^*}$	$f_{J/\Psi}$
131	207	158	214	200	230	250	275	405

approximated by products of matrix elements of the associated quark bilinears. As an example, the amplitude for the decay $B \rightarrow X_1 + X_2$ (where $X_{1,2}$ are arbitrary mesons) can be expressed as

$$\begin{aligned} &\langle X_1 X_2 | (\bar{q}_1 \Gamma b)(\bar{q}_2 \Gamma' q_3) | B \rangle \\ &\approx \langle X_1 | \bar{q}_1 \Gamma b | B \rangle \langle X_2 | \bar{q}_2 \Gamma' q_3 | 0 \rangle \\ &\quad + \langle X_2 | \bar{q}_1 \Gamma b | B \rangle \langle X_1 | \bar{q}_2 \Gamma' q_3 | 0 \rangle \\ &\quad + \frac{1}{N_c} \langle X_1 | \bar{q}_1 \Gamma'' b | B \rangle \langle X_2 | \bar{q}_2 \Gamma''' q_3 | 0 \rangle \\ &\quad + \frac{1}{N_c} \langle X_2 | \bar{q}_1 \Gamma'' b | B \rangle \langle X_1 | \bar{q}_2 \Gamma''' q_3 | 0 \rangle \end{aligned} \quad (6)$$

where the last two terms refer to Fierz rearranged currents. Of course, only some of the matrix elements are non-vanishing. For one, within this approximation, the contributions of the tensor operators in Eq. (5b) vanish identically. It should be noted that Eqs. (5b) and (6) contain only color-singlet currents. For color-octet currents to contribute, one would need to consider additional gluon exchanges. Such effects are clearly not factorizable. Within this approximation then, the color octet parts of Eq. (4) can be neglected, or, in other words, the contribution of a color-sextet diquark is almost indistinguishable from that of the corresponding color-triplet one.

III. HADRONIC B DECAYS

Within the SM, hadronic B decays may proceed through either tree level W boson exchange diagrams and/or through penguin diagrams (both QCD and electroweak). The corresponding effective Hamiltonian, including the QCD corrections have been presented in Refs. [16,17]. For brevity's sake, we do not repeat the entire list here. It suffices to remember that the SM amplitudes are proportional to the Fermi constant G_F , the relevant product of two CKM matrix elements $V_{ib} V_{jk}^*$ and/or $V_{tb} V_{tk}^*$ (with i and j as generic up type quarks and k as down type quark) and the combination of the Wilson coefficients that incorporates the short distance QCD corrections at the B mass scale. The considerable suppression due to the smallness of the CKM mixing is what makes B -decays sensitive to new physics effects.

Reverting to the calculation of the hadronic matrix elements, the decay constant f_i for a generic (pseudoscalar or vector) meson is defined through the relations

TABLE IV. Form factors at zero momentum transfer in the BSW model.

Decay mode	$F_1(0)$	$F_0(0)$	$V(0)$	$A_1(0)$	$A_2(0)$	$A_0(0)$
$B \rightarrow \pi$	0.33	0.33				
$B \rightarrow K$	0.38	0.38				
$B \rightarrow D$	0.69	0.69				
$B \rightarrow \rho$			0.33	0.28	0.28	0.28
$B \rightarrow K^*$			0.37	0.33	0.33	0.32
$B \rightarrow D^*$			0.71	0.65	0.69	0.62

$$\begin{aligned} \langle P(p_P) | \bar{q}_j \gamma_\mu \gamma_5 q_i | 0 \rangle &= -i f_P p_P^\mu \\ \langle V(p_V) | \bar{q}_j \gamma_\mu q_i | 0 \rangle &= f_V m_V \epsilon^\mu. \end{aligned} \quad (7)$$

Here it is assumed that the meson is composed of a $q_j \bar{q}_i$ pair. The decay constants are best determined from an analysis of the respective leptonic decay modes and the relevant ones

TABLE V. Values of pole masses in GeV.

Current	$m(0^-)$	$m(1^-)$	$m(1^+)$	$m(0^+)$
$\bar{u}b$	5.28	5.32	5.37	5.73
$\bar{d}b$	5.28	5.32	5.37	5.73
$\bar{s}b$	5.37	5.41	5.82	5.89
$\bar{c}b$	6.30	6.34	6.73	6.80

are listed in Table III. The matrix elements for the associated density operators may then be evaluated using the Dirac equation:

$$\begin{aligned} \partial^\alpha (\bar{q}_i \gamma_\alpha q_j) &= i(m_j - m_i) \bar{q}_i q_j \\ \partial^\alpha (\bar{q}_i \gamma_\alpha \gamma_5 q_j) &= i(m_j + m_i) \bar{q}_i \gamma_5 q_j. \end{aligned}$$

The matrix elements for quark bilinears between a B meson and a pseudoscalar-vector meson can be parametrized in terms of form factors:

$$\begin{aligned} \langle P(p_P) | \bar{q}_j \gamma_\mu (1 - \gamma_5) b | B(p_B) \rangle &= \left[(p_B + p_P)_\mu - \frac{m_B^2 - m_P^2}{q^2} q_\mu \right] F_1(q^2) + \frac{m_B^2 - m_P^2}{q^2} q_\mu F_0(q^2) \\ \langle V(p_V) | \bar{q}_j \gamma_\mu (1 - \gamma_5) b | B(p_B) \rangle &= -\epsilon_{\mu\nu\alpha\beta} \epsilon^{\nu*} p_B^\alpha p_V^\beta \frac{2V(q^2)}{(m_B + m_V)} - i \left(\epsilon_\mu^* - \frac{\epsilon^* \cdot q}{q^2} q_\mu \right) (m_B + m_V) A_1(q^2) \\ &\quad + i \left((p_B + p_V)_\mu - \frac{(m_B^2 - M_V^2)}{q^2} q_\mu \right) \frac{(\epsilon^* \cdot q) A_2(q^2)}{(m_B + m_V)} \\ &\quad - i \frac{2m_V (\epsilon^* \cdot q)}{q^2} q_\mu A_0(q^2), \end{aligned} \quad (8)$$

where $q = p_B - p_{P(V)}$ and ϵ is the polarization vector of V . The apparent poles at $q^2 = 0$ are fictitious since

$$F_1(0) = F_0(0)$$

$$2m_V A_0(0) = (m_B + m_V) A_1(0) - (m_B - m_V) A_2(0).$$

The numerical values of the form factors can be calculated within a given model. For our analysis, we adopt the BSW model [19,18], and the relevant form factors, at zero momentum transfer, are given in Table IV [16,19]. It can easily be checked that choosing a different model for the calculation of hadronic matrix elements would not change our results appreciably.

For the q^2 dependence of these form factors we assume a simple pole formula [16,19] $F(q^2) = F(0)/(1 - q^2/m_{pole}^2)$ with the pole mass m_{pole} the same as that of the lowest lying meson with the appropriate quantum numbers ($J^P = 0^+$ for F_0 ; 1^- for F_1 and V ; 1^+ for A_1 and A_2 ; 0^- for A_0). The values of these pole masses are presented in Table V [16,19].

With Eqs. (7) and (8) in place, calculation of the full matrix elements, within the factorization approximation, is now a straightforward task. Consider the decay $B(b\bar{q}_4) \rightarrow M_1(q_1\bar{q}_4) M_2(q_2\bar{q}_3)$ where M_i are generic mesons (pseudoscalar or vector). For simplicity's sake, assume that no two quarks are identical so that the quark bilinears [see Eq. (6)] cannot relate the B to M_2 . In this case, the amplitudes are given by

$$\begin{aligned}
\mathcal{A}[B \rightarrow P_1 P_2] &= i f_{P_2} (m_B^2 - m_{P_1}^2) F_0^{B \rightarrow P_1} (m_{P_2}^2) \left[-b_0 + b_1 + b_2 - b_3 - \frac{(b_4 - b_5 - b_6 + b_7) m_{P_2}^2}{(m_b - m_{q_1})(m_{q_2} + m_{q_3})} \right] \\
\mathcal{A}[B \rightarrow P_1 V_2] &= 2 f_{V_2} m_{V_2} F_1^{B \rightarrow P} (m_{V_2}^2) [b_0 + b_1 + b_2 + b_3] (\epsilon^* \cdot p_{P_1}) \\
\mathcal{A}[B \rightarrow V_1 P_2] &= 2 m_{V_1} f_{P_2} (\epsilon^* \cdot p_{P_2}) A_0^{B \rightarrow V_1} (m_{P_2}^2) \left[b_0 + b_1 - b_2 - b_3 + \frac{(b_4 + b_5 - b_6 - b_7) m_{P_2}^2}{(m_b + m_{q_1})(m_{q_2} + m_{q_3})} \right] \\
\mathcal{A}[B \rightarrow V_1 V_2] &= f_{V_2} m_{V_2} \left[-\epsilon_{\mu\nu\alpha\beta} \epsilon_2^{\mu*} \epsilon_1^\nu p_B^\alpha p_{V_1}^\beta \frac{V(m_{V_2}^2)}{(m_B + m_{V_1})} (b_0 + b_1 + b_2 + b_3) - i (\epsilon_1^* \cdot \epsilon_2) (m_B + m_{V_1}) \right. \\
&\quad \left. \times A_1(m_{V_2}^2) (-b_0 + b_1 - b_2 + b_3) + 2i (p_B \cdot \epsilon_2) (p_B \cdot \epsilon_1^*) \frac{A_2(m_{V_2}^2)}{(m_B + m_{V_1})} (-b_0 + b_1 - b_2 + b_3) \right]. \quad (9)
\end{aligned}$$

For decay modes wherein $q_3 = q_4$, the second set amplitudes in Eq. (6) contribute too. These additional pieces, however, can be easily read off from Eq. (9).

IV. $B^0 - \bar{B}^0$ MIXING

The main motivation for considering $B_d^0 - \bar{B}_d^0$ mixing to constrain diquark couplings is that this mixing is mediated by flavor changing neutral current, which is forbidden at the tree level in SM. The mixing is characterized by the experimentally measurable mass difference

$$\Delta M_d = m_{B_d^H} - m_{B_d^L} = \frac{|\langle \bar{B}_d^0 | \mathcal{H}_{\text{eff}} | B_d^0 \rangle|}{m_{B_d^0}} \quad (10)$$

with H and L denoting heavy and light mass eigenstates. The recent world average value of ΔM_d at 1σ limit is [20]

$$\Delta M_d = (0.472 \pm 0.017) \times 10^{-12} \text{ s}^{-1}. \quad (11)$$

In SM, this mixing proceeds through the box diagrams with internal top quark and W boson exchanges [21,22]. The diagrams in which one or both top quarks are replaced by up or charm quarks are negligible on account of: (i) the small mixing angles and (ii) the corresponding loop integrals being suppressed to a great extent due to the smallness of the light quark masses. Integrating out the internal particles, one thus gets an effective four quark interaction, with a $(V-A) \otimes (V-A)$ current structure, and scaling as to $m_W^2 G_F^2 |V_{tb} V_{td}^*|^2$. The short distance QCD corrections are well determined [21,22], while the long-distance corrections are estimated to be small, unlike in the case of $K - \bar{K}$ mixing.

In presence of diquarks, two different types of contributions may appear. If there exist $\Delta b = \Delta d = 2$ operators, then such a mixing can occur at the tree level itself. Else, additional contributions may appear in the form of new diquark-mediated box diagrams. However, as we are interested in small diquark couplings, we shall confine ourselves to tree level (in diquarks) processes only.

In the calculation of the hadronic matrix element $\langle \bar{B}_d^0 | \mathcal{H}_i | B_d^0 \rangle$, the vacuum saturation approximation is a convenient one. Herein, one inserts a complete set of states between the two currents and assumes that the sum is dominated by the vacuum and thus the hadronic matrix elements are proportional to f_B^2 by virtue of Eq. (7). The bag factor, B_B , introduced to parametrize all possible deviations from the vacuum saturation approximation, can be evaluated in various nonperturbative approaches. We use here the values of $B_B(\mu_b)$ and f_B as obtained by UKQCD Collaboration in a quenched lattice calculation [23].

Incorporating the contributions of all such current structures, in addition to that of the SM, we obtain

$$\begin{aligned}
\Delta M_d &= f_B^2 B_B m_{B_d^0} \left| b_0 + b_1 - b_2 - b_3 \right. \\
&\quad \left. + \frac{m_{B_d^0}^2}{(m_b + m_d)^2} (b_4 + b_5 - b_6 - b_7) \right|. \quad (12)
\end{aligned}$$

This may then be compared with the experimental value to obtain the required constraints.

V. RESULTS

Before we determine the bounds obtainable from B -phenomenology, it is worthwhile to reexamine the SM predictions for the decay modes of interest; this helps in selecting the channels likely to result in stronger constraints. As mentioned earlier, within the SM, the hadronic B decays are mediated by one or more of tree (W -mediated), electroweak penguin and QCD penguin diagrams. The branching fractions are determined primarily by the CKM mixings operative in the particular decay and, in case of one-loop processes, by the corresponding loop integral. For example, the decays $B^- \rightarrow D_s^- \pi^0$ and $\bar{B}_d^0 \rightarrow \pi^- \pi^+$ are suppressed in comparison to $\bar{B}_d^0 \rightarrow D^+ \pi^-, D^+ D_s^-, B^- \rightarrow D^0 D_s^-$ and $\bar{B}_s^0 \rightarrow D_s^+ \pi^-$ on account of V_{ub} being much smaller than V_{cb} .

Similar statements obviously hold for the decays into the corresponding excited states. In spite of such a suppression, the tree diagram far outweighs the one-loop contributions for any of these decays. For the decays $B^- \rightarrow K^- \pi^0$, $\bar{B}^0 \rightarrow \bar{K}^0 \pi^0$ (and the corresponding PV and VV modes) though, the tree level contributions are double Cabibbo suppressed with the consequence that these decays are dominated by penguin diagrams.⁵ Decays like $B^- \rightarrow K^- K^0$, $B^- \rightarrow \pi^- \bar{K}^0$, on the other hand, are governed solely by electroweak and/or QCD penguins.

A different suppression occurs for the processes $\bar{B}^0 \rightarrow \pi^0 \pi^0$, $\bar{B}^0 \rightarrow D^0 \pi^0$. Compared to the analogous modes $\bar{B}^0 \rightarrow \pi^+ \pi^-$, $\bar{B}_d^0 \rightarrow D^+ \pi^-$ wherein the charged mesons are created from the vacuum by a color-singlet current, these decays are obviously color-suppressed. In fact, the respective short distance coefficients differ by as much as a factor of 20 [16]. And finally, there are the annihilation diagrams in the decays like $B(b\bar{q}) \rightarrow X_1(q_1\bar{q})X_2(q_2\bar{q}_1)$. In these decays, b and \bar{q} in B meson annihilate to produce \bar{q} and q_2 quarks, which, in turn form final state mesons with $q_1\bar{q}_1$ pair, created from vacuum. As these contributions are proportional to the wave function at zero, they are typically much smaller than either of the tree or penguin mediated spectator contributions. For example, Ref. [16] argues that the annihilation amplitude for $B \rightarrow P_1 P_2$ modes is proportional to the mass difference of the mesons in the final state and hence there is essentially no annihilation contribution to $\bar{B} \rightarrow \pi^0 \pi^0$, $\bar{B}^0 \rightarrow K^+ K^-$ etc.

It is tempting to assume that the modes suppressed within the SM would be the ones most sensitive to effects from new physics. While this is largely so, a few points should be remembered. For one, color suppression and/or suppression of annihilation diagrams are essentially independent of weak matrix elements, and hence equally applicable to either the SM or a theory with diquarks. Secondly, even for decays wherein the SM amplitudes are Cabibbo suppressed, the experimental data may not be precise enough for it to be a very sensitive probe. Rather, it could well turn out that an unsuppressed mode may turn out to be one of the most sensitive on account of the observations matching very well with the SM predictions.

In obtaining numerical results, we assume that only one pair of diquark couplings are nonzero. While this restriction may seem unwarranted, it is an useful approximation that allows one a quantitative appreciation of the various experimental constraints. Furthermore, we assume a common mass of 100 GeV for all the diquarks. As explained earlier, mass splittings between states in a single multiplet is disfavored by LEP data. And since the effective four-Fermi operator goes as $m_{\Phi(V)}^{-2}$, for a general diquark mass, all our bounds on the products need only be rescaled by a factor of $(m_{\Phi(V)}/100 \text{ GeV})^2$.

⁵In our numerical calculations, we have used the Wilson coefficients as listed in Ref. [16] for $N_c=3$.

The bounds, as obtained from a given decay mode, can be broadly classified into two sets. An experimentally observed channel (with an associated error bar) would, in general, allow the diquark coupling pairs to lie in one of two non-contiguous windows, with the separation between the windows determined by the agreement of the SM contribution. On the other hand, decay modes that are yet to be experimentally seen, can only lead to a single window. For a specific combination of diquark couplings, we look at all such individual bounds and then delineate the range satisfied by each. As an illustration, let us consider the product $h_{13}^{(1)} h_{12}^{(1)}$. At 90% C.L., the ranges allowed by individual decays are as follows:

- (a): $\bar{B}^0 \rightarrow D^+ \rho^-$: $[-6.0 \times 10^{-2}, -5.3 \times 10^{-2}]$,
 $[-7.3 \times 10^{-3}, 5.7 \times 10^{-4}]$
- (b): $B^- \rightarrow D^0 \rho^-$: $[-4.7 \times 10^{-2}, -4.2 \times 10^{-2}]$,
 $[-4.7 \times 10^{-4}, 4.8 \times 10^{-3}]$
- (c): $B^- \rightarrow \bar{K}^{0*} \pi^-$: $[-8.9 \times 10^{-4}, 2.2 \times 10^{-3}]$
- (d): $\bar{B}^0 \rightarrow \bar{K}^0 \rho^0$: $[-2.1 \times 10^{-3}, 1.6 \times 10^{-3}]$.

Clearly the first window allowed by (a) is completely ruled out by the the bounds from (c) and (d). The same is true for the first window from (b). Progressively eliminating parts of the domains allowed by the individual decays, we find that the actual allowed range for this particular combination is only $[-4.7 \times 10^{-4}, 5.7 \times 10^{-4}]$. An identical strategy is adopted for all other combinations, and we list the best bounds in Tables VI–XIII.

A very important point is to be noted here. In the preceding analysis, while we have selected the range of parameters common to each constraint, we have not really used the entire information available to us. Such an analysis would involve the use of a statistical discriminator such as a χ^2 test or a likelihood test. While such an exercise is a straightforward one and would have led to bounds stricter than those we list, the decision to forego it was a conscious one. For, in the absence of higher order corrections and a more precise calculation of the hadronic matrix elements, the bounds derived here are only indicative. Hence, further refinement using statistical methods is not really called for.

From Table II, it is easy to see that $h_{ij}^{(1)} h_{kl}^{(1)}$, $h_{ij}^{(2)} h_{kl}^{(2)}$, $h_{ij}^{(3)} h_{kl}^{(3)}$, $h_{ij}^{(4)} h_{kl}^{(4)}$, $h_{ij}^{(7)} h_{kl}^{(7)}$ and $h_{ij}^{(8)} h_{kl}^{(8)}$ result in both neutral and charged current structures. In a given hadronic decay, both the operators contribute with one of them being color suppressed. Furthermore, the flavor structure determines whether two contributions interfere constructively or destructively. As a consequence, the bounds on $h_{ij}^{(2)} h_{kl}^{(2)}$, $h_{ij}^{(4)} h_{kl}^{(4)}$ and $h_{ij}^{(8)} h_{kl}^{(8)}$ are weaker by a factor of $(N_c + 1)/(N_c - 1)$ compared to those for $h_{ij}^{(1)} h_{kl}^{(1)}$, $h_{ij}^{(3)} h_{kl}^{(3)}$ and $h_{ij}^{(7)} h_{kl}^{(7)}$. A color-unsuppressed operator associated with combinations $\tilde{h}_{ij}^{(3)} \tilde{h}_{kl}^{(3)}$ and $\tilde{h}_{ij}^{(4)} \tilde{h}_{kl}^{(4)}$ are neutral current ones, their contributions to charged current decays are naturally

TABLE VI. Bounds on Φ_1 couplings in units of $(m_{\Phi_1}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$h_{13}^{(1)}h_{11}^{(1)}$	$\bar{B}^0 \rightarrow \pi^+ \pi^-, \bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-2.6 \times 10^{-3}, 1.1 \times 10^{-3}]$
$h_{13}^{(1)}h_{12}^{(1)}$	$\bar{B}^0 \rightarrow D^+ \rho^-, B^- \rightarrow \pi^- \bar{K}^{0*},$ $B^- \rightarrow D^0 \rho^-, \bar{B} \rightarrow \bar{K}^0 \rho^0$	$[-4.7 \times 10^{-4}, 5.7 \times 10^{-4}]$
$h_{23}^{(1)}h_{22}^{(1)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, \bar{B}^0 \rightarrow D^{+*} D_s^-$	$[-5.5 \times 10^{-2}, -5.3 \times 10^{-2}],$ $[-6.3 \times 10^{-3}, 2.2 \times 10^{-3}]$
$h_{23}^{(1)}h_{12}^{(1)}$	$\bar{B}^0 \rightarrow K^0 \bar{K}^0$	$[-1.4 \times 10^{-3}, 9.5 \times 10^{-4}]$
$h_{33}^{(1)}h_{11}^{(1)}$	$B_d^0 - \bar{B}_d^0$ mixing	$[-1.4 \times 10^{-7}, -1.3 \times 10^{-7}],$ $[-1.3 \times 10^{-8}, 2.4 \times 10^{-9}]$

TABLE VII. Bounds on Φ_2 couplings in units of $(m_{\Phi_2}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$h_{13}^{(2)}h_{12}^{(2)}$	$B^- \rightarrow K^- \pi^0, B^- \rightarrow \pi^- \bar{K}^0,$ $B^- \rightarrow K^- \rho^0$	$[-1.2 \times 10^{-3}, 7.7 \times 10^{-4}]$
$h_{23}^{(2)}h_{12}^{(2)}$	$\bar{B}^0 \rightarrow K^0 \bar{K}^0$	$[-2.8 \times 10^{-3}, 1.9 \times 10^{-3}]$

TABLE VIII. Bounds on Φ_3 couplings in units of $(m_{\Phi_3}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$h_{13}^{(3)}h_{12}^{(3)}$	$B^- \rightarrow \pi^0 K^-, B^- \rightarrow K^- \rho^0,$ $\bar{B}^0 \rightarrow D^+ \rho^-$	$[-1.9 \times 10^{-4}, 2.9 \times 10^{-4}]$
$h_{23}^{(3)}h_{12}^{(3)}$	$B^- \rightarrow \pi^0 D_s^-, \bar{B}^0 \rightarrow \rho^0 J/\Psi$	$[-1.5 \times 10^{-3}, 1.7 \times 10^{-3}]$
$\tilde{h}_{13}^{(3)}\tilde{h}_{11}^{(3)}$	$\bar{B}^0 \rightarrow \pi^0 \pi^0, B^- \rightarrow \pi^- \pi^0$	$[-3.3 \times 10^{-3}, 2.9 \times 10^{-3}]$
$\tilde{h}_{13}^{(3)}\tilde{h}_{12}^{(3)}$	$\bar{B}^0 \rightarrow \pi^0 \bar{K}^{0*}, B^- \rightarrow \pi^0 K^-$	$[-1.1 \times 10^{-3}, 7.8 \times 10^{-4}]$
$\tilde{h}_{13}^{(3)}\tilde{h}_{21}^{(3)}$	$B^- \rightarrow \pi^- D^0, B^- \rightarrow \rho^- D^0,$ $\bar{B}^0 \rightarrow D^0 \pi^0$	$[-1.2 \times 10^{-3}, 1.4 \times 10^{-3}]$
$\tilde{h}_{23}^{(3)}\tilde{h}_{22}^{(3)}$	$B^- \rightarrow D^0 D_s^{*-}, \bar{B}^0 \rightarrow D^+ D_s^-$	$[-8.7 \times 10^{-3}, 3.1 \times 10^{-2}]$
$\tilde{h}_{23}^{(3)}\tilde{h}_{21}^{(3)}$	$\bar{B}^0 \rightarrow \pi^0 J/\Psi, \bar{B}^0 \rightarrow \rho^0 J/\Psi$	$[-4.2 \times 10^{-3}, 4.2 \times 10^{-3}]$
$\tilde{h}_{23}^{(3)}\tilde{h}_{12}^{(3)}$	$B^- \rightarrow D_s^- \pi^0$	$[-1.4 \times 10^{-2}, 2.0 \times 10^{-2}]$
$h_{13}^{(3)}\tilde{h}_{11}^{(3)}$	$\bar{B}^0 \rightarrow \pi^0 \pi^0, \bar{B}^0 \rightarrow \pi^+ \pi^-$	$[-5.1 \times 10^{-3}, 7.4 \times 10^{-3}]$
$h_{12}^{(3)}\tilde{h}_{13}^{(3)}$	$B^- \rightarrow K^- \pi^0, \bar{B}^0 \rightarrow \pi^+ K^-,$ $B^- \rightarrow D^0 \pi^-$	$[-1.4 \times 10^{-3}, 2.2 \times 10^{-3}]$
$h_{13}^{(3)}\tilde{h}_{12}^{(3)}$	$B^- \rightarrow K^- \pi^0, \bar{B}^0 \rightarrow \pi^+ K^-,$	$[-3.3 \times 10^{-3}, 1.4 \times 10^{-3}],$
$h_{13}^{(3)}\tilde{h}_{21}^{(3)}$	$\bar{B}^0 \rightarrow D^0 \pi^0, B^- \rightarrow D^0 \pi^-,$ $B^- \rightarrow D^0 \rho^-$	$[-5.2 \times 10^{-3}, 2.2 \times 10^{-3}]$
$h_{23}^{(3)}\tilde{h}_{22}^{(3)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, B^- \rightarrow D^0 D_s^-,$ $\bar{B}^0 \rightarrow D^+ D_s^{*-}$	$[-4.9 \times 10^{-3}, 1.5 \times 10^{-2}]$
$h_{12}^{(3)}\tilde{h}_{23}^{(3)}$	$B^- \rightarrow \pi^0 D_s^-$	$[-1.6 \times 10^{-2}, 1.1 \times 10^{-2}]$
$h_{23}^{(3)}\tilde{h}_{12}^{(3)}$	$B^- \rightarrow \pi^0 D_s^-$	$[-1.1 \times 10^{-2}, 1.6 \times 10^{-2}]$

TABLE IX. Bounds on Φ_4 couplings in units of $(m_{\Phi_4}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$h_{13}^{(4)}h_{11}^{(4)}$	$\bar{B}^0 \rightarrow \pi^+ \pi^-, \bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-1.1 \times 10^{-3}, 2.6 \times 10^{-3}]$
$h_{13}^{(4)}h_{12}^{(4)}$	$\bar{B}^0 \rightarrow \pi^+ K^-, \bar{B}_s^0 \rightarrow K^+ K^-,$ $\bar{B}^0 \rightarrow \bar{K}^0 \rho^0$	$[-8.1 \times 10^{-4}, 5.4 \times 10^{-4}]$
$h_{23}^{(4)}h_{22}^{(4)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, \bar{B}^0 \rightarrow D^+ * D_s^-$ $B^- \rightarrow D^0 D_s^-$	$[-2.2 \times 10^{-3}, 6.5 \times 10^{-3}],$ $[5.3 \times 10^{-2}, 5.5 \times 10^{-2}]$
$h_{23}^{(4)}h_{12}^{(4)}$	$B^- \rightarrow \pi^0 D_s^-, \bar{B}^0 \rightarrow \rho^0 J/\Psi$	$[-3.4 \times 10^{-3}, 3.1 \times 10^{-3}]$
$\tilde{h}_{13}^{(4)}\tilde{h}_{11}^{(4)}$	$\bar{B}^0 \rightarrow \pi^0 \pi^0, B^- \rightarrow \pi^- \pi^0$	$[-3.3 \times 10^{-3}, 2.9 \times 10^{-3}]$
$\tilde{h}_{13}^{(4)}\tilde{h}_{12}^{(4)}$	$\bar{B}^0 \rightarrow \pi^0 \bar{K}^{*0}, B^- \rightarrow \pi^0 K^-$	$[-1.1 \times 10^{-3}, 7.8 \times 10^{-4}]$
$\tilde{h}_{13}^{(4)}\tilde{h}_{21}^{(4)}$	$B^- \rightarrow \pi^- D^0, B^- \rightarrow \rho^- D^0,$ $\bar{B}^0 \rightarrow D^0 \pi^0$	$[-1.2 \times 10^{-3}, 1.4 \times 10^{-3}]$
$\tilde{h}_{23}^{(4)}\tilde{h}_{22}^{(4)}$	$B^- \rightarrow D^0 D_s^- * , \bar{B}^0 \rightarrow D^+ D_s^-$	$[-8.7 \times 10^{-3}, 3.1 \times 10^{-2}]$
$\tilde{h}_{23}^{(4)}\tilde{h}_{21}^{(4)}$	$\bar{B}^0 \rightarrow \pi^0 J/\Psi, \bar{B}^0 \rightarrow \rho^0 J/\Psi$	$[-4.2 \times 10^{-3}, 4.2 \times 10^{-3}]$
$\tilde{h}_{23}^{(4)}\tilde{h}_{12}^{(4)}$	$B^- \rightarrow D_s^- \pi^0$	$[-1.4 \times 10^{-2}, 2.0 \times 10^{-2}]$
$h_{11}^{(4)}\tilde{h}_{13}^{(4)}$	$\bar{B}^0 \rightarrow \pi^+ \pi^-, \bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-3.6 \times 10^{-3}, 6.4 \times 10^{-3}]$
$h_{13}^{(4)}\tilde{h}_{11}^{(4)}$	$\bar{B}^0 \rightarrow \pi^+ \pi^-, \bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-6.4 \times 10^{-3}, 3.6 \times 10^{-3}]$
$h_{12}^{(4)}\tilde{h}_{13}^{(4)}$	$\bar{B}_s^0 \rightarrow K^+ K^-, \bar{B}^0 \rightarrow \pi^+ K^-$	$[-1.5 \times 10^{-2}, -1.2 \times 10^{-2}]$ $[-2.3 \times 10^{-3}, 1.5 \times 10^{-3}]$
$h_{13}^{(4)}\tilde{h}_{12}^{(4)}$	$\bar{B}_s^0 \rightarrow K^+ K^-, \bar{B}^0 \rightarrow \pi^+ K^-,$ $\bar{B}^0 \rightarrow \bar{K}^{*0} \pi^0, B^- \rightarrow K^- \rho^0$	$[-1.5 \times 10^{-3}, 2.3 \times 10^{-3}]$
$h_{13}^{(4)}\tilde{h}_{21}^{(4)}$	$\bar{B}^0 \rightarrow D^0 \pi^0, B^- \rightarrow D^0 \rho^-$	$[-3.7 \times 10^{-3}, 2.2 \times 10^{-2}]$
$h_{22}^{(4)}\tilde{h}_{23}^{(4)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, B^- \rightarrow D^0 D_s^-,$	$[-3.5 \times 10^{-3}, 1.0 \times 10^{-2}]$ $[7.4 \times 10^{-2}, 8.8 \times 10^{-2}]$
$h_{23}^{(4)}\tilde{h}_{22}^{(4)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, B^- \rightarrow D^0 D_s^-,$ $\bar{B}^0 \rightarrow D^+ * D_s^-$	$[-1.0 \times 10^{-2}, 3.5 \times 10^{-3}]$
$h_{12}^{(4)}\tilde{h}_{23}^{(4)}$	$B^- \rightarrow \pi^0 D_s^-$	$[-8.0 \times 10^{-3}, 1.2 \times 10^{-2}]$
$h_{23}^{(4)}\tilde{h}_{12}^{(4)}$	$B^- \rightarrow \pi^0 D_s^-$	$[-1.2 \times 10^{-2}, 8.0 \times 10^{-3}]$

color-suppressed. Hence the corresponding bounds are weaker. Similarly, since $h_{ij}^{(3)}\tilde{h}_{kl}^{(3)}$ and $h_{ij}^{(4)}\tilde{h}_{kl}^{(4)}$, are associated only with scalar, pseudoscalar and tensor operators, they cannot contribute to $B \rightarrow VV$ decays. And finally, as the diquarks $\Phi_{7,8}$ couple only to down-type quarks, they can contribute only to those decays that occur in the SM solely through penguin diagrams.

The diquark Φ_3 differs from Φ_4 only by color quantum number. But since the color factors for triplet and sextet diquarks are accidentally equal and $\tilde{h}_{ij}^{(3)}$ and $\tilde{h}_{ij}^{(4)}$'s have no

specific symmetry property under the exchange of i and j , the bounds on the product of these couplings are exactly the same. A similar story obtains for other sets of diquarks, (V_1, V_2) and (V_3, V_4) .

As discussed earlier, $\tilde{h}_{ij}^{(4)}$ is analogous to the trilinear R parity violating coupling λ_{ijk}'' . Thus the constraints on $\tilde{h}_{ij}^{(4)}\tilde{h}_{kl}^{(4)}$ are equivalent to those on $\lambda_{imj}''\lambda_{kml}''$. The upper bound on $\tilde{h}_{23}^{(4)}\tilde{h}_{22}^{(4)}$ is marginally weaker than that, quoted in [13] and [24]. The upper bound on $\tilde{h}_{11}^{(4)}\tilde{h}_{13}^{(4)}$ is much weaker

TABLE X. Bounds on Φ_7 couplings in units of $(m_{\Phi_7}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$h_{13}^{(7)}h_{11}^{(7)}$	$\bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-1.1 \times 10^{-3}, 1.3 \times 10^{-3}]$
$h_{13}^{(7)}h_{12}^{(7)}$	$B^- \rightarrow \pi^- \bar{K}^0, B^- \rightarrow \pi^0 K^-$ $\bar{B}^0 \rightarrow \pi^0 \bar{K}^0$	$[-3.9 \times 10^{-4}, 7.3 \times 10^{-4}]$
$h_{23}^{(7)}h_{12}^{(7)}$	$\bar{B}^0 \rightarrow K^0 \bar{K}^0$	$[-9.5 \times 10^{-4}, 1.4 \times 10^{-3}]$
$h_{33}^{(7)}h_{11}^{(7)}$	$B_d^0 - \bar{B}_d^0$ mixing	$[-1.4 \times 10^{-7}, -1.3 \times 10^{-7}],$ $[-1.3 \times 10^{-8}, 2.4 \times 10^{-9}]$

TABLE XI. Bounds on Φ_8 couplings in units of $(m_{\Phi_8}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$h_{13}^{(8)}h_{12}^{(8)}$	$B^- \rightarrow \pi^- \bar{K}^0, B^- \rightarrow \pi^0 K^-$ $B^- \rightarrow \pi^- \bar{K}^{0*}$	$[-7.9 \times 10^{-4}, 1.2 \times 10^{-3}]$
$h_{23}^{(8)}h_{12}^{(8)}$	$\bar{B}^0 \rightarrow K^0 \bar{K}^0$	$[-1.9 \times 10^{-3}, 2.8 \times 10^{-3}]$

TABLE XII. Bounds on V_1 couplings in units of $(m_{V_1}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$\vartheta_{13}^{(1)}\vartheta_{11}^{(1)}$	$\bar{B}^0 \rightarrow \pi^+ \pi^-, \bar{B}^0 \rightarrow \rho^0 \rho^0$	$[-6.2 \times 10^{-3}, 3.2 \times 10^{-3}]$
$\vartheta_{31}^{(1)}\vartheta_{11}^{(1)}$	$\bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-2.1 \times 10^{-3}, 1.7 \times 10^{-3}]$
$\vartheta_{13}^{(1)}\vartheta_{12}^{(1)}$	$\bar{B}^0 \rightarrow \pi^+ K^-, B^- \rightarrow K^- \pi^0,$ $\bar{B}^0 \rightarrow \rho^+ K^-$	$[-1.4 \times 10^{-3}, 2.0 \times 10^{-3}],$ $[1.0 \times 10^{-2}, 1.1 \times 10^{-2}]$
$\vartheta_{13}^{(1)}\vartheta_{21}^{(1)}$	$B^- \rightarrow \pi^- \bar{K}^0, B^- \rightarrow \pi^- \bar{K}^{0*}$ $B^- \rightarrow D^0 \rho^-$	$[-1.0 \times 10^{-3}, -2.6 \times 10^{-3}],$ $[-8.1 \times 10^{-4}, 9.8 \times 10^{-4}]$
$\vartheta_{31}^{(1)}\vartheta_{12}^{(1)}$	$\bar{B}^0 \rightarrow \bar{K}^0 \pi^0, B^- \rightarrow \pi^- \bar{K}^0$ $B^- \rightarrow \pi^- \bar{K}^{0*}$	$[-9.8 \times 10^{-4}, 8.1 \times 10^{-4}]$
$\vartheta_{31}^{(1)}\vartheta_{21}^{(1)}$	$B^- \rightarrow \pi^0 K^-, B^- \rightarrow \rho^0 K^-$	$[-1.5 \times 10^{-3}, 5.3 \times 10^{-4}]$
$\vartheta_{23}^{(1)}\vartheta_{22}^{(1)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, B^- \rightarrow D^0 D_s^-$ $B^- \rightarrow D^{0*} D_s^-$	$[-9.2 \times 10^{-3}, 3.1 \times 10^{-3}]$
$\vartheta_{23}^{(1)}\vartheta_{12}^{(1)}$	$\bar{B}^0 \rightarrow K^0 \bar{K}^0$	$[-1.3 \times 10^{-3}, 1.9 \times 10^{-3}]$
$\vartheta_{23}^{(1)}\vartheta_{21}^{(1)}$	$B^0 \rightarrow \pi^0 J/\Psi, B^0 \rightarrow \rho^0 J/\Psi,$	$[-2.1 \times 10^{-3}, 2.1 \times 10^{-3}]$
$\vartheta_{32}^{(1)}\vartheta_{12}^{(1)}$	$B^0 \rightarrow K^0 \bar{K}^0$	$[-4.7 \times 10^{-3}, 7.0 \times 10^{-3}]$
$\vartheta_{32}^{(1)}\vartheta_{21}^{(1)}$	$B^0 \rightarrow K^0 \bar{K}^0$	$[-1.9 \times 10^{-3}, 1.3 \times 10^{-3}]$
$\vartheta_{33}^{(1)}\vartheta_{11}^{(1)}$	$B_d^0 - \bar{B}_d^0$ mixing	$[-5.4 \times 10^{-8}, -4.8 \times 10^{-8}],$ $[-4.9 \times 10^{-9}, 8.9 \times 10^{-10}]$

TABLE XIII. Bounds on V_3 couplings in units of $(m_{V_3}/100 \text{ GeV})$ at 90% C.L.

Product of couplings	Mode	Allowed region
$\vartheta_{13}^{(3)}\vartheta_{11}^{(3)}$	$\bar{B}^0 \rightarrow \pi^+ \pi^-$	$[-3.2 \times 10^{-3}, 1.0 \times 10^{-2}]$
$\vartheta_{31}^{(3)}\vartheta_{11}^{(3)}$	$\bar{B}^0 \rightarrow \pi^0 \pi^0$	$[-1.4 \times 10^{-3}, 1.8 \times 10^{-3}]$
$\vartheta_{13}^{(3)}\vartheta_{12}^{(3)}$	$\bar{B}^0 \rightarrow \pi^+ K^-, \bar{B}^0 \rightarrow \rho^+ K^-$	$[-1.2 \times 10^{-2}, -1.0 \times 10^{-2}],$ $[-2.0 \times 10^{-3}, 1.4 \times 10^{-3}]$
$\vartheta_{13}^{(3)}\vartheta_{21}^{(3)}$	$B^- \rightarrow D^0 \rho^-, \bar{B}^0 \rightarrow D^{0*} \rho^0$	$[-1.0 \times 10^{-3}, 9.0 \times 10^{-3}]$
$\vartheta_{31}^{(3)}\vartheta_{12}^{(3)}$	$\bar{B}^0 \rightarrow D^0 \pi^0, B^- \rightarrow D^0 \pi^-$	$[-1.5 \times 10^{-3}, 4.1 \times 10^{-3}]$
$\vartheta_{31}^{(3)}\vartheta_{21}^{(3)}$	$B^- \rightarrow \pi^0 K^-, B^- \rightarrow K^- \rho^0$	$[-7.3 \times 10^{-4}, 1.7 \times 10^{-3}]$
$\vartheta_{32}^{(3)}\vartheta_{22}^{(3)}$	$\bar{B}^0 \rightarrow D^+ D_s^-, B^- \rightarrow D^0 D_s^-$ $B^- \rightarrow D^{0*} D_s^-$	$[-1.5 \times 10^{-3}, 4.3 \times 10^{-3}]$ $[-2.2 \times 10^{-1}, -1.4 \times 10^{-1}],$ $[-2.0 \times 10^{-2}, 5.7 \times 10^{-2}]$
$\vartheta_{23}^{(3)}\vartheta_{12}^{(3)}$	$\bar{B}^0 \rightarrow D_s^- \pi^+$	$[-1.1 \times 10^{-2}, 1.4 \times 10^{-2}]$
$\vartheta_{23}^{(3)}\vartheta_{21}^{(3)}$	$\bar{B}^0 \rightarrow \rho^0 J/\Psi$	$[-2.5 \times 10^{-3}, 2.1 \times 10^{-3}]$
$\vartheta_{32}^{(3)}\vartheta_{12}^{(3)}$	$\bar{B}^0 \rightarrow \pi^0 J/\Psi$	$[-2.5 \times 10^{-3}, 2.1 \times 10^{-3}]$
$\vartheta_{32}^{(3)}\vartheta_{21}^{(3)}$	$B^- \rightarrow D_s^- \pi^0$	$[-7.0 \times 10^{-3}, 1.0 \times 10^{-2}]$

than that listed in [13] and [24]. For other $\tilde{h}_{ij}^{(4)}\tilde{h}_{kl}^{(4)}$, we obtain more stringent bounds. For most of the nonsupersymmetric product of two diquark couplings, our predicted bounds are much stronger.

In a few hadronic decays ($B^- \rightarrow K^- J/\Psi$, $B^- \rightarrow \pi^- J/\Psi$ etc.), the SM predictions do not agree with the experimental observation even at 2σ level. While this could be construed as an indication of new physics (in our case diquarks), we prefer to tread a more conservative path. Consequently, we have not included such decays in our analysis.

Turning to $B_d^0\text{-}\bar{B}_d^0$ mixing, it is obvious that a tree-level contribution will accrue from any four-quark operator that violates both b and d number by two units each. Furthermore, Φ_2 , Φ_3 and Φ_8 do not contribute to $B_d^0\text{-}\bar{B}_d^0$ mixing by virtue of the antisymmetric nature of their couplings under exchange of the flavor indices. Thus $B_d^0\text{-}\bar{B}_d^0$ mixing only imposes limits on the parameter space of Φ_1 , Φ_7 , V_1 and V_2 diquarks. Unlike in the SM, $B_d^0\text{-}\bar{B}_d^0$ occurs at tree level in the presence of nonzero diquark couplings. Accordingly, we obtain most stringent bounds on $h_{33}^{(A)}h_{11}^{(A)}$ and $\vartheta_{33}^{(A)}\vartheta_{11}^{(A)}$.

VI. CONCLUSION

In summary, we have studied the leading order effects of scalar and vector diquark and/or R parity violating couplings on hadronic B decays and $B\text{-}\bar{B}$ mixing. Clearly a diquark must have more than one non-zero couplings to SM fields to be able to mediate such processes. We take the economic standpoint that only *any two* of such couplings are non-zero. Analyzing the present data on B -decays, we derive constraints on such pairs that are significantly stronger than those derived from other low energy processes. Theoretical improvements on nonfactorization effects and estimates of annihilation form factors as well as precise measurements of the decay modes at the upcoming B factories in the near future will improve our bounds on the parameter space for diquarks and/or R parity violating couplings in the minimal supersymmetric SM.

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