If the underlying flavor symmetry is Abelian, quark mixings in $d_R$ sector are the most prominent. Such flavor violating effects can reveal itself through $d_R$ quark mixings if supersymmetry is realized in Nature. Quark-squark alignment is necessary to deal with $\Delta m_K$ and $\epsilon_K$ constraints, but interestingly, with $m_1, m_2 \sim$ TeV, the $d_R$ mixing effects are comparable to $B_L$ and $B_s$ mixings in the Standard Model, while $D^0$ mixing is tantalizingly close to some hints from data. $CP$ phases in these mixings would therefore be deviant, and $|V_{td}|$ and $V_{ub}^{\ast}$ may be larger than allowed by unitarity constraints, which can be checked by the BaBar and Belle experiments. Mixing induced $CP$ violation in $b \rightarrow s\gamma$ and $d\gamma$ transitions can be obtained, in particular, by sizable enhancement with an extra tan $\beta$ factor from non-standard soft breaking terms. Heavy superparticles can escape present flavor changing neutral current (FCNC) bounds and direct searches at colliders, but reveal themselves in the $B$ system.

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I. INTRODUCTION

Despite its great success, the Standard Model (SM) is widely regarded as a weak scale effective theory. We expect to encounter new phenomena that arise from a more complete theory as we increase the luminosity and energy of our probes. Although we have not observed any convincing deviation of experimental results from SM predictions (except neutrino oscillations [1]), there are a few hints on the existence of New Physics.

It is well known that the size of the $CP$ violating phase, $\phi_1$ (or $\beta$) $\equiv \arg V_{td}$ (PDG phase convention [2]), can be measured via the time dependent asymmetry,

$$a_{J/\psi K^0} = \frac{\Gamma(B^0(t) \rightarrow J/\psi K^0) - \Gamma(B^0(t) \rightarrow \bar{J}/\psi K^0)}{\Gamma(B^0(t) \rightarrow J/\psi K^0) + \Gamma(B^0(t) \rightarrow \bar{J}/\psi K^0)} = \sin 2\phi_1 \sin \Delta m_{B_d} \tau.$$

Following earlier measurements [2, 3], the BaBar and Belle Collaborations have recently firmly established [2, 3] $\sin 2\phi_1$ to be nonzero. Compared to earlier low results of $\sin 2\phi_1 = 0.58^{+0.32}_{-0.34} \pm 0.09$ (Belle) [3] and $0.34 \pm 0.20 \pm 0.05$ (BaBar) [2], the corresponding numbers are now $0.99 \pm 0.14 \pm 0.06$ (Belle) [3] and $0.59 \pm 0.14 \pm 0.05$ (BaBar) [2], respectively. Combining the two most recent values without systematic errors, one gets the average value

$$\sin 2\phi_1 = 0.79 \pm 0.10. \quad (2)$$

While this is still consistent with the Cabibbo–Kobayashi–Maskawa unitarity (CKM) fit value of $\sin 2\phi_1 = 0.698 \pm 0.066$ [4, 5], the central value is now somewhat on the high side, especially for the Belle number. If this trend persists — which we would know by summer 2002 — it would imply the presence of New Physics. With this in mind, it is clearly a good time to study other related $CP$ violation processes, especially those that are suppressed in the SM.

It has recently been shown that charmless rare $B$ decays favor [1] a value for $\phi_3$ (or $\gamma$) $\equiv \arg V_{ub}^{\ast}$ that is larger than the one obtained from the CKM unitarity fit [6]. The latter is dominated by recent improved bounds on $\Delta m_{B_s}/\Delta m_{B_d}$ [7],

$$\Delta m_{B_s} = 0.484 \pm 0.010 \, \text{ps}^{-1}, \quad (3)$$
$$\Delta m_{B_d} > 15.0 \, \text{ps}^{-1} \text{ at } 95\% \, \text{C.L.}, \quad (4)$$

which tends to squeeze out the $\phi_3$ possibility. One also has $\text{Br}(K^+ \rightarrow \pi^+\nu\bar{\nu}) = 4.2^{+3.7}_{-3.5} \times 10^{-10}$ from the E787 Collaboration [8], where the central value is several times above the SM expectation, implying a rather large $|V_{td}|$. The last branching ratio, of course can be viewed as decreasing with time since no new events have been found. Combining the above, however, perhaps a more consistent picture would be if $B_d$ or $B_s$ mixings have additional New Physics sources. This may already be indicated by the measurement of $\sin 2\phi_1$ in Eq. (2) as we have discussed. It can alternatively be tested in the $CP$ phase of $B_s$ mixing, which can be studied at the Tevatron collider in the next few years. If a non-vanishing value is found, it would definitely imply New Physics since the SM prediction is zero.

The E791, CLEO, FOCUS, Belle and BaBar Collaborations have reported search results for $CP$ asymmetries in the neutral $D$-meson system [14, 15]. The search for $D^0-\bar{D^0}$ mixing using CLEO II. V data suggests that $x_D \equiv \Delta m_D/\Gamma$ is less than 2.9% at 95% C.L. [16], which is far below previous results [17]. The actual numbers, however, are $x_D^2/2 < 0.041$ and $-5.8\% < y_D < 1.0\%$, in terms of

$$x_D' = x_D \cos \delta_D + y_D \sin \delta_D$$
$$y_D = y_D \cos \delta_D - x_D \sin \delta_D, \quad (5)$$

where $\delta_D$ is the relative strong phase between the doubly Cabibbo suppressed $D^0 \rightarrow K^+\pi^-$ and the Cabibbo favored $D^0 \rightarrow K^+\pi^-$ decay amplitudes. The SM predictions of these mixing parameters are small, $x_D \sim 10^{-5} - 10^{-4}$ and $y_D \sim 10^{-4} - 10^{-2}$ [17] (a recent discussion suggests $y_D \sim 1\%$ [18]). The CLEO Collaboration
been observed \cite{27}. However, the Belle Collaboration has arrived at $|x_D| < 2.9\%$ by assuming $\delta_D$ is as small as suggested by SU(3) and other arguments \cite{23,24}. If one removes \cite{25} the prejudice that $D^0 \to K^+\pi^-$ and $K^-\pi^+$ amplitudes have the same strong phase, the result on $y_D'$ may actually be hinting at a sizable $x_D$, which would strongly suggest short distance New Physics interactions. It is therefore important to compare $x_D$, $y_D'$ with other $D^0$-$\bar{D}^0$ mixing related measurements.

The CP asymmetry parameter $y_{CP}$ is related to the lifetime difference between $D^0 \to K^+\pi^-$ and $D^0 \to K^-\pi^+$, where the former is flavor specific and the latter a CP eigenstate. The FOCUS Collaboration had found the intriguing value of $y_{CP} = (3.42 \pm 1.39 \pm 0.74)\%$ \cite{14}, which has a significance of more than $2\sigma$. However, recently, the CLEO, Belle and BaBar Collaborations find much lower values. The current world average is $(1.1 \pm 0.9)\%$ \cite{22}. Still, the point remains that one could have large $x_D$ if $\sin\delta_D$ is sizable \cite{26}.

The processes $b \to s\gamma$, $d\gamma$ occur only at loop level in the SM, therefore they are naturally sensitive to New Physics. The processes $B \to K^{*}\gamma$ and $b \to s\gamma$ have long been observed \cite{27}. However, the Belle Collaboration observes a $3.5 \pm 2.1$ event excess for $B \to \rho\gamma$, giving $\text{Br}(B \to \rho\gamma) < 10.6 \times 10^{-6}$ \cite{28} at 90\% C.L. The quark level $b\gamma$ coupling is usually parametrized as

$$H_{SM} = -\frac{G_F}{\sqrt{2}} \frac{e}{4\pi^2} V_{tb} V_{ts}^* m_b \bar{q}(C_{7\gamma}, R)$$
$$+C_{7\gamma}'L|\sigma_{\mu\nu} F^{\mu\nu} b + h.c.,$$

(6)

where in the SM, $C_{7\gamma} \approx -0.3$ at the typical B decay energy scale $\mu \approx 5$ GeV. Due to the left-handed nature of weak interactions, $C_{7\gamma}$ dominates while $C_{7\gamma}'$ is suppressed by $m_q/m_b$. This may not be the case, however, in models beyond the SM, and interesting CP violating asymmetries in mixing induced radiative B decay can occur \cite{29}. In a previous work \cite{30}, we have discussed the $b \to s\gamma$ process in the context of New Physics. We found that large direct CP violations are possible, in contrast to SM expectations which are very small. For mixing induced CP violation in B decay, we also find large asymmetries due to enhanced $C_{7\gamma}$. Furthermore, the mixing dependent asymmetry in $B \to \rho\gamma$ is more accessible than in $B \to K^{*\gamma}$, because $\rho^0 \to \pi^+\pi^-$ can give the vertex information needed for time dependence, while $K^{*0} \to K\pi^0$ unfortunately does not \cite{24}. Given that SM predictions on direct CP violation in $b \to d\gamma$ are in general not small \cite{31}, mixing induced CP violation in $b \to d\gamma$ is a more sensitive probe of New Physics. Of course, mixing dependent CP violation in $b \to s\gamma$ processes can be more readily probed in $B_s$ system such as in $B_s \to \phi\gamma$.

Mixing induced CP violation in radiative B decay is similar to the golden $J/\psi K_S$ mode. The hadronic uncertainties factor out for self-conjugate CP eigenstate(s). Furthermore, one needs both $B$ and $\bar{B}$ to decay to the same final state to allow interference to take place. Therefore, the formula of the asymmetry $a_{M^{\pm\gamma}}$, where $M^0$ is a CP self-conjugate hadron such as $\rho^0$, $\omega$, $K^{*0}(in K_S\pi^0)$, resembles that of $a_{J/\psi K_S}$:

$$a_{M^{\pm\gamma}} = \xi \sin 2\theta \sin(2\phi_1 - \phi_{7\gamma} - \phi_{7\gamma}') \sin \Delta m t$$

(7)

where $\xi$ is the CP eigenvalue of $M^0$,

$$\sin 2\theta = \frac{2|C_{7\gamma}C_{7\gamma}'|}{|C_{7\gamma}|^2 + |C_{7\gamma}'|^2},$$

(8)

is the mean strength of the two chiral amplitudes, and $\phi_{7\gamma}'$ is the phase of $C_{7\gamma}'$. While the “$\sin 2\theta$” in $B \to J/\psi K_S$ case is equal to one, an important feature for $B \to M^{0\gamma}$ is that one needs both $C_{7\gamma}$ and $C_{7\gamma}'$, for the interference to occur. This is the reason why these asymmetries are suppressed in the SM since $|C_{7\gamma}^{\text{SM}}/C_{7\gamma}'^{\text{SM}}| \ll 1$, hence precisely why they are sensitive to New Physics. The Belle Collaboration may be able to test the asymmetry to an accuracy of 10\% in about a decade \cite{22}. Hadron machines may also be able to study this if they can observe the radiative B decay modes, since they produce many more $B\bar{B}$s than the $e^+e^-$ machines.

Another hint for New Physics may come from the rather large value of the newly observed $\epsilon'/\epsilon$ \cite{32}. It has prompted New Physics considerations \cite{34}, even though the experimental value could be accommodated within the SM. We do not consider New Physics hints from muon anomalous magnetic moment.

We are interested in New Physics models that can lead to deviations in the above mentioned processes. It is interesting that supersymmetry (SUSY) models with Abelian horizontal symmetry (AHS) can provide a unified framework for all such New Physics effects. This will be presented in Section II, where we will explain the implications of AHS on quark mixing, in particular on the origin of large right-handed down quark mixings. The effect is carried over to the squark sector in SUSY models, and squark mixings will impact on flavor changing neutral currents (FCNC). In face of stringent constraints from kaon mixings, quark-squark alignment (QSA) is invoked to produce texture zeros, where we will give an explicit example of horizontal charge assignments. The subsequent sections are devoted to various FCNCs induced by squark mixings. In Section III we study the effect on $B_d$-$\bar{B}_d$ and $B_s$-$\bar{B}_s$, which sets the scale of superparticle masses. In Section IV, we show that the chargino contribution on kaon mixings gives a similar superparticle mass scale. A generic feature of QSA is the possibility of sizable $D^0-\bar{D}^0$ mixing. It is interesting that in AHS models with SUSY, with sparticle scale fixed by $B-B$ mixing, $x_D$ could be right in the ball-park of the CLEO range, as we show in Section V. Section VI is devoted to radiative B decay in SUSY models, and discussions and conclusions are given in Sections VII and VIII, respectively.
II. ABELIAN FLAVOR SYMMETRY WITH SUSY

A. Abelian Flavor Symmetry and Large $d_R$ Mixings

Fermion masses and mixings in the Cabibbo-Kobayashi-Maskawa matrix $V_{\text{CKM}}$ exhibit an intriguing hierarchical pattern in powers of $\lambda \equiv |V_{\text{us}}|$:  

$$m_u/m_c \sim \lambda^3, \quad m_c/m_t \sim \lambda^4, \quad m_t/\sqrt{G_F} \sim 1, \quad m_d/m_s \sim \lambda^2, \quad m_s/m_b \sim \lambda^2, \quad m_b/m_t \sim \lambda^3, \quad (9)$$

$$|V_{cb}| \sim \lambda^2, \quad |V_{ub}| \sim \lambda^3.$$  

The structure could be due to an underlying symmetry \cite{33,34}, the breaking of which gives an expansion in $\lambda \sim (S)/M$, with $S$ a scalar field and $M$ a high scale. In these models, after integrating out some massive fields of mass $M$, one obtains non-renormalizable terms \cite{33,34}.

$$L_{\text{mass}} = \lambda_q^2 H_d (S/M)^{\alpha_q ij} Q_i \bar{q}_{Rj} + \text{h.c.},$$  

where $q$ are summed over up and down type quarks, $\lambda^q$ are $O(1)$ numbers, and $i, j$ are generation indices. $L_{\text{mass}}$ is made a horizontal symmetry singlet by choosing appropriate powers of $S$, i.e. $\alpha_{q ij}$. For models with Abelian horizontal symmetry, without loss of generality \cite{33}, we can define the horizontal charges of the scalar fields as

$$H(H_d) = H(H_u) = 0, \quad H(S) = -1.$$  

The breaking of the horizontal symmetry as well as electroweak symmetry lead to quark mass elements,

$$M_q^{ij} = \lambda_q^2 (H_d)(S/M)^{\alpha_q ij} Q_i \bar{q}_{Rj} + H(Q_i) + H(\bar{q}_{Ri}).$$  

It is now easy to see that

$$M_{ij} M_{ji} \sim M_i M_{jj}, \quad (i, j \text{ not summed}),$$  

which follows as a consequence of the commuting nature of horizontal charges.

By assuming small mixing angles, as one tries to understand the hierarchical pattern in $\lambda$ by AHS, the quark mass ratios fix the order of magnitude of the diagonal elements of quark mass matrices. The upper right part of the mass matrix $M_q$ corresponds to $U_{qL}$ rotation, which is related to $V_{\text{CKM}} = U_{UL} U_{dL}^\dagger$. Small mixing angle and naturalness imply $U_{qL} \sim V_{\text{CKM}}$. By using Eq. (13), one can work out the lower left part and hence the whole mass matrix \cite{35,36}.

$$\frac{M_u}{m_t} \sim \begin{bmatrix} \lambda^7 & \lambda^5 & \lambda^3 \\ \lambda^5 & \lambda^4 & \lambda^2 \\ \lambda^4 & \lambda^3 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{bmatrix}, \quad \frac{M_d}{m_b} \sim \begin{bmatrix} \lambda^4 & \lambda^3 & \lambda^2 \\ \lambda^3 & \lambda^2 & \lambda^2 \\ \lambda^2 & \lambda^2 & \lambda \end{bmatrix}. \quad (14)$$

Since $U_{qL}$ are restricted by $V_{\text{CKM}}$, mixing angles in $U_{qR}$ are in general greater. In particular, we find that $M_{32}^2/m_b$ and $M_{31}^2/m_b$ are the most prominent off-diagonal elements.

To summarize, we have,

$$U_{qL} \sim \begin{bmatrix} 1 & \lambda & \lambda^3 \\ \lambda & 1 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{bmatrix}, \quad U_{dR} \sim \begin{bmatrix} 1 & \lambda & \lambda \\ \lambda & 1 & \lambda \\ \lambda & \lambda & 1 \end{bmatrix},$$

$$U_{uR} \sim \begin{bmatrix} 1 & \lambda & \lambda^4 \\ \lambda & 1 & \lambda^2 \\ \lambda^4 & \lambda^2 & 1 \end{bmatrix}. \quad (15)$$

It is clear that mixing angles in $U_{dR}$ are in general greater than those in $U_{UL}$ and $U_{uR}$, especially when the $b$ flavor is involved. Although large mixings in the right handed down quark sector are useless or well hidden within SM, $B_d$ and $B_s$ mixings are naturally susceptible to New Physics involving further dynamics related to the right-handed down flavor sector.

As one of the leading candidates for New Physics, SUSY helps resolve many of the potential problems that emerge when one extends beyond the SM, for example the gauge hierarchy problem, unification of SU(3)×SU(2)×U(1) gauge couplings, and so on \cite{38}. It is interesting to note that large mixing in right-handed down quark sector will be transmitted to right-handed down squarks, if the breakings of flavor symmetry and SUSY are not closely related. The flavor symmetry gives better control on soft breaking parameters resulting in a more predictive SUSY model. We now elevate Eq. (10) to the superpotential as well as similar forms in trilinear terms, the so called A terms. By flavor symmetry, we still have the same power of $S$ to ensure that the whole term remains a horizontal singlet, that is,

$$\tilde{M}_Q^{ij}_{\text{LR}} \sim \tilde{m}_t \tilde{M}_d^{ij}, \quad (\tilde{M}_d^{ij})_{\text{LR}} = (\tilde{M}_d^{ij})_{\text{LR}}.$$  

which are roughly proportional to respective quark mass matrices, hence their effects are suppressed by $m_q/\tilde{m}$ \cite{37}. While the symmetry does not require the new $\lambda_{ij}$ in the A-term to be the same as $\lambda_{ij}$ in Eq. (10), $(\tilde{M}_d^{ij})_{\text{LR}}$ cannot, in general, be diagonal in the quark mass basis. From Eq. (10), one easily gets $(\tilde{M}_d^{ij})_{\text{LL}}/\tilde{m}_t^2 \sim V_{\text{CKM}}$, while

$$\tilde{M}_d^{ij}_{RR} \sim \tilde{m}_t \begin{bmatrix} 1 & \lambda & \lambda \\ \lambda & 1 & \lambda \\ \lambda & \lambda & 1 \end{bmatrix}. \quad (17)$$

The squark mixings will have impact on FCNC because of extra dynamics involving $\tilde{q} \tilde{g}, \tilde{q} \tilde{\chi}^\pm$ and $\tilde{q} \tilde{\chi}^0$ couplings. It is now clear that FCNC processes involving $b_R$ are sensitive probes of this generic class of SUSY models.

* These large mixings in second and third generation $d_R$’s may be related to large mixing in second and third generation neutrinos \cite{39} when considering grand unified theories (GUT) \cite{40}. 

3
The $RR$ sector could contribute significantly to $B_d$ and $B_s$ mixings, to be discussed in the next section, via large mixings in $\tilde{b}_R\tilde{d}_R$ and in $\tilde{b}_R\tilde{s}_R$ as shown in Eq. (17).

B. Quark-squark Alignment

In order to compare squark mixing angles with FCNC constrains, we will use the mass insertion approximation \[10\] in the following discussion. It is customary to take squarks as almost degenerate at scale $\tilde{m}$. In quark mass basis, one defines \[40\].

\[
\delta_{i,j}^{ij} \equiv \frac{[U^\dagger_q A(M^2_q)_{AB} U_q B]^{ij}}{\tilde{m}^2},
\]

which is roughly the squark mixing angle. $M^2_q$ are squark mass matrices, $A, B = L, R$, and $i, j$ are generation indices. Note that $\delta_{13}^{dRR} \sim \lambda$ and $\delta_{23}^{dRR} \sim 1$, while LR and RL mixings are suppressed by $m_q/\tilde{m}$.

It is well known that kaon mixings give stringent constraint on new flavor violating source. The 12, 21 elements in $M^2_{12}$ and $(M^2_q)_{21}$ in Eq. (17) will induce too large a contribution to kaon mixing via gluino box diagrams \[12\]. One has to suppress these squark mixings by enforcing approximate “texture zeros”. This can be done by invoking quark-squark alignment (QSA) \[38\], by using two (or more) singlet fields $S_i$ to break the $U(1) \times U(1)$ (or higher) Abelian horizontal symmetry, and making use of the holomorphic nature of the superpotential in SUSY models. For example, we may have a term with negative power $\sigma_{qij}$ in $S$ in Eq. (10) to satisfy the horizontal symmetry. The term $S^{-1}^{a_{qij}}$ is simply $(S^*)^a_{a_{qij}}$. As we promote Eq. (10) to superpotential, which can only be a function of superfields and not of conjugate superfields at the same time, one can no longer use $S^*$ and hence there is a zero in that particular $ij$-th element. One can have $M^2_{dLRR} = 0$ which imply $U_{dLRR}^{21} = 0$ or are highly suppressed, and likewise $(M^2_{12})_{dLRR}$ are also suppressed. Thus, $\delta_{d}^{21}$ and $\delta_{d}^{12}$ can be suppressed and the kaon mixing constraint is satisfied accordingly.

There is one subtlety involving our choice to retain $(M^2_{12})_{dLRR}$, which arises from $M^2_{31}$. The mass matrix $M_d$ is diagonalized by a bi-unitary transform, hence, it is of the form

\[
\frac{M_d}{m_b} = U^\dagger_{dL} M^\text{diag}_{dL} U_{dR} = \sum_{i,j} \left( \begin{array}{c} 1 \\ \lambda^a \\ \lambda^b \\ \lambda^c \end{array} \right) \left( \begin{array}{c} \lambda^4 \\ \lambda^2 \\ 0 \\ 0 \end{array} \right) \lambda^d \left( \begin{array}{c} \lambda^a \\ \lambda^b \\ \lambda^c \\ 1 \end{array} \right)
\]

where the diagonal matrix in the middle of the right hand side corresponds to the diagonal down quark mass (ratio) matrix, and the one to its left (right) corresponds to $U_{dL}$ ($U_{dR}$). Multiplying out the matrices in Eq. (18), we have

\[
\frac{M^1}{m_b} \sim \lambda^4 + \lambda^{2+a+d} + \lambda^{b+e}, \quad \frac{M^2}{m_b} \sim \lambda^4 + \lambda^{2+a} + \lambda^{b+f}, \quad \frac{M^3}{m_b} \sim \lambda^4 + \lambda^{2+a} + \lambda^{e+c}.
\]

We see that, by retaining $M^3_{dL}/m_b \sim \lambda$, we have $e = 1$ and hence $U^T_{dR} \sim \lambda$. We will have $c = 2$, if also we retain $M^3_{dL}/m_b \sim \lambda^2$. However, as the kaon mixing constraint requires

\[
\frac{M^1_{dL}/m_b \sim \lambda^4 + \lambda^{2+a} + \lambda^{b+e} \sim 0},
\]

which in turn gives $d = 1$ for $\lambda^{2+a}$ to be of the same order as $\lambda^{c+e} = \lambda^3$, to allow cancellation to take place. This will still give squark mixing angle $\sim \lambda$ between $d_{R} - \tilde{b}_{R}$, which is not acceptable. A closer look reveals that $M^3_{d}/m_b \sim 1$ is also unacceptable. It gives $f = 0$ from $M^3_{d}/m_b \sim 1$. With the requirement of $M^3_{d}/m_b = 0$, we again have $c = 2$ and it follows that $d = 1$ by the same argument. Because $(M^2_{12})_{dLRR} \sim \lambda$ is kept, we need to make $M^3_{d}/m_b$ and $M^3_{d}/m_b$ also vanishing in face of stringent $\Delta m_K$ and $\varepsilon_K$ constraints. The decoupling of $s$ flavor from other generations thus follows from imposing QSA in 12 sector and choosing to retain $M^3_{dL} \neq 0$.

In the usual approach of quark-squark alignment, this subtlety does not arise because one aspires to lower $m_3$, $m_4$ for sake of collider and other signatures. Since $\delta_{13}^{dRR} \sim \lambda$ and $\delta_{23}^{dRR} \sim 1$ would then violate $B_d$ mixing and $b \to s\gamma$ constraints already, they are eliminated from the outset. As a result, $M^3_{d}/m_b \sim \lambda^2$, though of little consequence, can be retained.

C. Explicit Examples of QSA

To be specific, we now give an explicit assignment of the horizontal charges of quark superfields and the resulting mass matrices as an illustrative example.

Since $|V_{ub}| \sim 0.002 - 0.005 < \lambda^3$ \[4\], we may use a smaller parameter $\lambda = 0.18$ instead of $\lambda$. We use two $S_i$ fields to break the horizontal symmetry,

\[
\frac{\langle S_I \rangle}{M} \sim \tilde{\lambda}^{0.5}, \quad \frac{\langle S_2 \rangle}{M} \sim \tilde{\lambda}^{0.5}.
\]

The horizontal charges of $S_I$ and $S_2$ are $(-1, 0), (0, -1)$, respectively and those of $Q, \tilde{d}_R$ and $\tilde{u}_R$ are given by

\[
Q_1: (8, -2), \quad Q_2: (1, 3), \quad Q_3: (2, -2), \quad \tilde{d}_{R1}: (-2, 10), \quad \tilde{d}_{R2}: (9, -3), \quad \tilde{d}_{R3}: (-2, 8), \quad \tilde{u}_{R1}: (-3, 11), \quad \tilde{u}_{R2}: (0, 3), \quad \tilde{u}_{R3}: (-1, 2),
\]
for small \( \tan \beta \). For \( \tan \beta \approx 50 \) we change horizontal charges of \( \bar{d}_R \) to,
\[
\bar{d}_{R1} : (-2, 5), \quad \bar{d}_{R2} : (4, -3), \quad \bar{d}_{R3} : (-2, 3). \tag{24}
\]
In this way we get,
\[
\frac{M_u}{m_t} \sim \begin{pmatrix}
\lambda^6 & \lambda^4 & \lambda^3 \\
0 & \lambda^3 & \lambda^2 \\
0 & \lambda & 1 \\
\end{pmatrix}, \quad \frac{M_d}{m_t} \lambda^{2.5} \sim \begin{pmatrix}
\lambda^4 & 0 & \lambda^3 \\
0 & \lambda^2 & 0 \\
0 & \lambda & 1 \\
\end{pmatrix}. \tag{25}
\]

From Eq. 23, we have effectively decoupled the second generation from the first and third in \( M_d \), which corresponds to suppressed \( U_{12}^{dL, dR} = 0 \). The case is reminiscent of \cite{80} where we decouple \( d \) flavor. The corresponding squark mass matrices from Eq. 23 are \((\tilde{M}_2)^{ij}_{LR} \sim \bar{m}_t M_q^{ij}\) and
\[
(\tilde{M}_Q)^{ij}_{LL} \sim \bar{m}_t \begin{pmatrix}
\lambda^6 & \lambda^4 & \lambda^3 \\
1 & 1 & 1 \\
\lambda^3 & \lambda^3 & 1 \\
\end{pmatrix}, \tag{26}
\]
\[
(\tilde{M}_u)^{ij}_{RR} \sim \bar{m}_t \begin{pmatrix}
\lambda^6 & \lambda^4 & \lambda^3 \\
1 & 1 & 1 \\
\lambda^3 & \lambda^3 & 1 \\
\end{pmatrix}, \tag{27}
\]
\[
(\tilde{M}_d)^{ij}_{RR} \sim \bar{m}_t \begin{pmatrix}
\lambda^6 & \lambda^4 & \lambda^3 \\
1 & 1 & 1 \\
\lambda^3 & \lambda^3 & 1 \\
\end{pmatrix}. \tag{28}
\]

For large \( \tan \beta \), we change \( (\tilde{M}_d)^{12}_{RR} \) and \( (\tilde{M}_d)^{23}_{RR} \) to \( \lambda^7 \) and \( \lambda^6 \), respectively. We summarize all \( \delta \)'s of interest in Table I. We will see that these values are all well below the limits from \( \Delta m_K \) and \( \epsilon \) constraints, even with \( \mathcal{O}(1) \) phases.

At this point one thing needs to be emphasized. In face of severe kaon mixing constraints, we did not choose horizontal charges to create large right handed \( \bar{d} - \bar{b} \) squark mixings. Instead, the choice of horizontal charges were to retain this natural large mixing that follow from the Abelian nature of the underlying flavor symmetry. Both the Abelian flavor symmetry and the mixing pattern originate from the observed mass mixing hierarchy pattern. Phenomenological consequences of these large mixings \cite{45} should be explored \cite{46}.

There is a generic feature \cite{30, 57} of QSA that is worthy of note. Having \( U_{dL}^{12} \approx 0 \) implies \( U_{dL}^{12} \approx \lambda, \) which can also be read off from Eq. 23. One now has \( \delta_{uL}^{12} \approx \lambda, \) as one can see from Eq. 18. This Cabibbo strength \( \delta_{uL}^{12} \) can contribute to kaon mixing via chargino diagrams, and also \( D^0 - \bar{D}^0 \) mixing via gluino diagrams, as we will discuss in Sec. IV and V, respectively. We note that New Physics contributions to \( D^0 - \bar{D}^0 \) mixing are of great interest at present, since the recent CLEO (and FOCUS) results may be a hint for \( D^0 \) mixing in disguise. Note also that the texture zeros of \( M_u^{21,31} \) are generated through QSA. The zero of \( M_u^{21} \) is needed to avoid \( \delta_{uRR}^{12} \approx \lambda, \) for otherwise, together with \( \delta_{uL}^{12} \approx \lambda \) they will induce too large a contribution to \( D \) mixing.

The zero of \( M_u^{31} \) is to avoid a feed back to \( \delta_{uRR}^{12} \), analogous to the discussion in the previous subsection.

We mention another subtlety arising from the Kähler potential \cite{45}. When the horizontal symmetry is spontaneously broken, mixing also occurs in the kinetic terms. By canonical normalization of the kinetic terms, further mixings are introduced. For example, \( M_q \) now becomes \( L_q M_q R_q, \) where \( L_q \approx (\tilde{M}_Q^2/\bar{m}_t^2)^{-1/2} \) and \( R_q \approx (\tilde{M}_u^2/\bar{m}_t^2)^{-1/2} \). The zeros in Eq. 23 are now all lifted, and are called filled zeros \cite{43}, giving rise to the underlined terms in Table I. \( U_{ql} \) also becomes \( L_q U_{ql} \) and similarly for other mixing matrices. Modifications of previous results can be achieved by suitable rotations and are also shown in Table I.

A second possibility of horizontal charge assignment is to retain \( M_d^{12} \) and \( M_d^{23} \) while having vanishing \( M_u^{12} \), i.e.
\[
\frac{M_d}{m_b} \sim \begin{pmatrix}
\lambda^4 & 0 & 0 \\
0 & \lambda^2 & \lambda^2 \\
0 & 1 & 1 \\
\end{pmatrix}. \tag{29}
\]

The assignment of horizontal charges for this case can be found in Ref. 30. The squark mixing can generate sizable contribution in \( B_s \) mixings \cite{47}. In this case one may have large CP phase in \( B_s \) mixing and possible effects in \( b \rightarrow s \gamma \) to be discussed later. The democratic structure of \( \tilde{M}_d^{2RR} \) in Eq. 17 in the 2-3 sub-matrix leads to approximate maximal mixing in \( \bar{s}_R \) and \( \bar{b}_R \). The
large off-diagonal elements \((\tilde{M}_t)^{23,32}\) lead to large level splitting. This allows for a possibly light strange beauty squark with interesting mixing in \(B_s\) mixing and direct search, but leaving \(Br(B \to X_s \gamma)\) largely unaffected [5].

It is interesting that, faced with stringent kaon constraint, AHS models with QSA allow large mixing in either \(b_R - d_R\) or \(b_R - s_R\), but not both at the same time. Thus, a prediction of this model is, if it is responsible for the smallness of the measured \(\sin 2\alpha_1\) (assuming that the low value persists in the future), there will be no large New Physics contribution to \(B_s\) mixing.

We now study the FCNC induced by these squark mixings in the following sections.

### III. \(B^0 - \bar{B}^0\) MIXING

In this section, we first focus on the general formalism of neutral \(B\) meson mixings in the AHS model with SUSY. We will focus on the \(B_d\) system for applications, which is readily extendable to the \(B_s\) system. For the latter system, the intriguing possibility that large right-handed squark mixings could lead to a light “strange-beauty” squark will be discussed briefly in Sec. VII.C.

The effective Hamiltonian for \(B_d^0 - \bar{B}_d^0\) mixings from SUSY contributions, where \(q = d\) or \(s\), is given by

\[
H_{\text{eff}} = \sum_i C_i \mathcal{O}_i,
\]

where,

\[
\mathcal{O}_1 = \tilde{q}_L^0 \gamma_\mu b_L^0 \tilde{q}_L^0 \gamma_\mu b_L^0,
\]

\[
\mathcal{O}_2 = \tilde{q}_L^0 \gamma_\mu W_R^0 \tilde{q}_L^0 \gamma_\mu W_R^0,
\]

\[
\mathcal{O}_3 = \tilde{q}_R^0 \gamma_\mu \tilde{q}_L^0 \gamma_\mu \tilde{w}_L^0,
\]

\[
\mathcal{O}_4 = \tilde{q}_R^0 \gamma_\mu \tilde{q}_L^0 \gamma_\mu \tilde{w}_R^0,
\]

\[
\mathcal{O}_5 = \tilde{q}_R^0 \gamma_\mu \tilde{q}_L^0 \gamma_\mu \tilde{b}_L^0
\]

together with three other operators \(\tilde{O}_{1,2,3}\) that are chiral conjugations \((L \leftrightarrow R)\) of \(O_{1,2,3}\). The Wilson coefficients receive charged Higgs, chargino, gluino, gluino-neutralino, and neutralino box diagram contributions,

\[
C_i = C_i^{H^\pm} + C_i^{\tilde{q}^0} + C_i^{\tilde{g}^0} + C_i^{\tilde{G}^0} + C_i^{\tilde{H}^0},
\]

where the Feynman diagrams are shown in Fig. 1.

#### A. Formulas

**Charged Higgs box**: [4][5]:

\[
C_i^{H^\pm} = \frac{\alpha^2}{8m_W^2} (V_{tb} V_{ts}^\ast)^2 \left[ x_{W} x_{H} \cot \theta \left( G(x_{W}, x_{H}) \right) + 2x_{W}^2 \cot \theta \left( F(x_{W}, x_{H}) \right) \right],
\]

**Chargino box**:

\[
C_i^{\tilde{G}^0} = \frac{\alpha^2}{8m_W^2} G^{0}(x_{\tilde{G}^0}, x_{\tilde{G}^0}) A_{jk} A_{kj}
\]

\[
C_i^{\tilde{H}^0} = -\frac{\alpha^2}{2m_W^2} \sqrt{x_{\tilde{H}^0}} \tilde{F}(x_{\tilde{H}^0}, x_{\tilde{H}^0}) U_{jk} U_{kj} \tilde{V}_{ij} \tilde{V}_{ij}
\]

where the indices \(j, k\) are summed over 1 to 2, and

\[
A_{jk} \equiv V_{l_1} V_{k_1}^\ast V_{l_2} V_{k_2}^\ast V_{l_3} V_{k_3}^\ast V_{l_4} V_{k_4}^\ast \delta_{u_{R} L} - V_{j_1} V_{j_2} \tilde{Y}_{j_1} \tilde{Y}_{j_2} V_{l_1} V_{l_2} \delta_{u_{R} L}
\]

\[
B_{jk} \equiv V_{l_1} V_{l_2} V_{l_3} V_{l_4} V_{l_5} V_{l_6} \delta_{u_{R} L} - V_{j_1} V_{j_2} \tilde{Y}_{j_1} \tilde{Y}_{j_2} V_{l_1} V_{l_2} \delta_{u_{R} L}.
\]

the indices \(l, m\) are summed over 3 generations of up type squarks, \(\tilde{Y}_{u_{c,t}} = m_{u_{c,t}}(\sqrt{2m_W} \sin \beta)\) and similarly, \(\tilde{Y}_{d_{s,b}} = m_{d_{s,b}}(\sqrt{2m_W} \cos \beta)\). The chargino mix-
ing matrices $\mathcal{U}, \mathcal{V}$ in Eq. (34) diagonalize the chargino mass matrix,

$$M_{\tilde{\chi}^{\pm}} = \mathcal{U}^* \left( \begin{array}{c} \frac{M_2}{\sqrt{2} m_W \cos \beta} & \sqrt{2} m_W \sin \beta \end{array} \right) \mathcal{V},$$

and

$$\left( \tilde{F}'(x, y), \tilde{G}'(x, y) \right) = x^2 \partial_a \partial_b \left( F'(a, b, y), G'(a, b, y) \right)|_{a=b=x},$$

which can also be expressed as

$$\left( \tilde{F}'(x, y), \tilde{G}'(x, y) \right) = \int_0^\infty \frac{dk^2 x^2 (k^2)^{1/2}}{(k^2 + 1)(k^2 + x)^4 (k^2 + y)},$$

which are always positive. There is no chargino contribution to $C_2$ and $C_2$ because of the color structure of the chargino box diagrams, and other terms are suppressed by the smallness of $Y_q$. Since $Y_L$ is large and $Y_R$ can also be large for the case of large tan $\beta$, we keep them in $C_{1,3}$. Note that in the AHS models, $\sum_{l,m} V_{lq} V_{mb} \delta_{llm}$ and $V_{lq} V_{mb} \delta_{llm}$ are roughly of the order of $|V_{lq}|$. As we will show soon, the $\Delta m_{\tilde{g}}$ constraint require $\overline{m} \sim \text{TeV}$ due to large $b_R \to d_R$ mixings. A typical LR stop mixing term contains $\hat{Y}_L V_{lq} \delta_{llm}$ with $V_{lq}$ small in $\text{MeV}$, which will be as small as $0.1 V_{lq}$ for $\overline{m} \sim \text{TeV}$. Furthermore, the flavor scale may not be too far from TeV [30], and there may not be much room for RG running to bring down the stop mass. Unlike the usual approach where one has light stop, the contributions from stop LR mixings are relatively small here.

**Gluino box [1]:**

$$C_1^\tilde{g} = \frac{\alpha_s a^w}{2m_\tilde{g}} \left[ \frac{1}{4} \left( 1 - \frac{1}{N_c} \right)^2 x_{\tilde{g} \tilde{g}} f_0(x_{\tilde{g} \tilde{g}}) \right. \left. + \frac{1}{8} \left( N_c - \frac{2}{N_c} + \frac{1}{N_c^2} \right) \tilde{f}_0(x_{\tilde{g} \tilde{g}}) \right] \delta_{dL}^{\tilde{g}} \delta_{rL}^{\tilde{g}} ;$$

$$C_2^\tilde{g} = \frac{\alpha_s a^w}{m_\tilde{g}^2} \left[ \left( N_c - \frac{2}{N_c} \right) x_{\tilde{g} \tilde{q}} f_0(x_{\tilde{g} \tilde{q}}) \frac{\tilde{f}_0(x_{\tilde{g} \tilde{q}})}{N_c} \right] \delta_{dL}^{\tilde{g}} \delta_{rR}^{\tilde{g}} ;$$

$$C_3^\tilde{g} = \frac{\alpha_s a^w}{m_\tilde{g}^2} \left[ x_{\tilde{g} \tilde{q}} f_0(x_{\tilde{g} \tilde{q}}) \frac{\tilde{f}_0(x_{\tilde{g} \tilde{q}})}{N_c} + \left( 1 + \frac{1}{N_c^2} \right) \right] \delta_{dL}^{\tilde{g}} \delta_{rR}^{\tilde{g}},$$

where $N_c$ is the number of colors, and

$$f_0(x) = \frac{(17 - 9x - 9x^2 + x^3 + 6 \ln x + 18 x \ln x)}{6(x - 1)^5},$$

$$\tilde{f}_0(x) = \frac{(1 + 9x - 9x^2 - x^3 + 6 \ln x + 6x^2 \ln x)}{3(x - 1)^5}.$$  

They are related to $\tilde{F}'(x, y)$ and $\tilde{G}'(x, y)$ by

$$\left( x_{\tilde{g}0}(x), -\tilde{f}_{\tilde{g}}(x) \right) = x^{-1} \left( \tilde{F}'(x^{-1}, 1), \tilde{G}'(x^{-1}, 1) \right).$$

$$\tilde{C}_1^\tilde{g}$$ is obtained by interchanging $L \leftrightarrow R$ in $C_1^\tilde{g}$. We neglect $C_{2,3}$ and $\tilde{C}_{2,3}^\tilde{g}$ due to the smallness of LR and RL mixings. There are usual box and crossed box diagrams. Terms with $N_c$ are from the former while $O(1)$ terms are from the latter, which can be easily checked by 't Hoof double line notation. Terms with $1/N_c, 1/N_c^2$ are sub-leading contributions from these two types of diagrams. We note that $C_4^\tilde{g}$ contains the largest $N_c$ factor hence is the most sensitive to squark mixings. Note also that one has a zero in $C_1^\tilde{g} (\tilde{C}_1^\tilde{g})$ for $x_{\tilde{g} \tilde{g}} \sim 2.43$.

**Gluino-neutralino box:**

$$C_1^\tilde{g} \tilde{\chi}_0^0 = \frac{\alpha_s a^w}{2m_\tilde{g}} \left( 1 - \frac{1}{N_c} \right) \left[ G^1_L G^3_L \tilde{G}'(x_{\tilde{q} \tilde{g}}, x_{\tilde{g} \tilde{q}}) \right. \left. - (G^1_L G^3_L + G^1_L G^3_L) \tilde{F}'(x_{\tilde{g} \tilde{q}}, x_{\tilde{q} \tilde{g}}) \right] (\delta_{dL}^{\tilde{g}})^2 ;$$

$$C_2^\tilde{g} \tilde{\chi}_0^0 = \frac{\alpha_s a^w}{2m_\tilde{g}} \left( 1 - \frac{1}{N_c} \right) H_{\tilde{b}L}^1 H_{\tilde{b}L}^1 \times \tilde{F}'(x_{\tilde{g} \tilde{q}}, x_{\tilde{q} \tilde{g}}) (\delta_{dL}^{\tilde{g}})^2 ;$$

$$C_3^\tilde{g} \tilde{\chi}_0^0 = \frac{\alpha_s a^w}{2m_\tilde{g}} \left( 1 - \frac{1}{N_c} \right) H_{\tilde{b}L}^1 H_{\tilde{b}L}^1 \times \tilde{F}'(x_{\tilde{g} \tilde{q}}, x_{\tilde{q} \tilde{g}}) (\delta_{dL}^{\tilde{g}})^2 ;$$

$$C_4^\tilde{g} \tilde{\chi}_0^0 = \frac{\alpha_s a^w}{m_\tilde{g}^2} \left[ \left( G^1_L G^3_L + G^1_R G^3_L \right) \frac{1}{N_c} H_{\tilde{b}L}^1 H_{\tilde{b}R}^1 \right. \left. \tilde{G}'(x_{\tilde{q} \tilde{g}}, x_{\tilde{g} \tilde{q}}) \right. \left. - 2 (G^1_L G^3_R + G^1_R G^3_L) \right. \right. \left. \times \tilde{F}'(x_{\tilde{g} \tilde{q}}, x_{\tilde{q} \tilde{g}}) \right] (\delta_{dL}^{\tilde{g}})^2 ;$$

$$C_5^\tilde{g} \tilde{\chi}_0^0 = \frac{\alpha_s a^w}{2m_\tilde{g}} \left[ \left( H_{\tilde{b}L}^1 H_{\tilde{b}R}^1 - \frac{2}{N_c} G^1_L G^3_R - \frac{2}{N_c} G^1_R G^3_L \right) \right. \left. \tilde{G}'(x_{\tilde{q} \tilde{g}}, x_{\tilde{g} \tilde{q}}) \right. \left. + \frac{4}{N_c} (G^1_L G^3_R + G^1_R G^3_L) \right. \left. \times \tilde{F}'(x_{\tilde{g} \tilde{q}}, x_{\tilde{q} \tilde{g}}) \right] (\delta_{dL}^{\tilde{g}})^2 ;$$

where $j$ is summed over 1 to 4. The mixing matrices $G_{L,R}, H_{L,L,R}$ are given by

$$G^i_L \equiv \tan \theta_W Y_Q N^c_{j1} + T_{3D} N^c_{j1}, \quad G^i_R \equiv \tan \theta_W Q_d N^c_{j1},$$

$$H^j_{d,s,b} \equiv N^c_{j3} \tilde{Y}_{d,s,b}, \quad H^j_{d,s,b} \equiv N^c_{j3} \tilde{Y}_{d,s,b} ;$$

where $Y_Q$ is the usual hypercharge, and the neutralino mixing matrix $\tilde{\cal{N}}$ diagonalizes the mass matrix $M_{\tilde{\chi}_0^0} = N^c \tilde{\cal{N}}^\dagger N^c$,
SUSY contributions are consistent with those in Ref. [49]. Terms with $G^j G^{ij}$ are from the usual box diagram, while terms with $G^j G^i, G^j G^i, H^j H^j$ are from the crossed diagram. If the mass parameters $M_1, M_2, |µ| > m_Z$, we will have simpler forms for $N_i$, where $N_{ij,kl} ≈ δ_{i,j,δ_{k,l}}$, and higgsinos become maximally mixed. This will lead to a large cancellation of higgsino contributions in $C^{θ0}_3, C^{θ0}_3$. The diagonalization of the neutralino mass matrix usually leads to a negative mass eigenvalue, say $m_χ^{θ0}_1$. It is well known that one can deal with it by two equivalent ways. One could either choose the phases in $N_i$ such that $m_χ^{θ0}_i$ are real and positive, or one could absorb the negative sign into $P_L χ^{θ0}_1$, and modify Feynman rules accordingly [31]. However, there is a subtlety when dealing with the crossed diagrams in the latter approach. An additional negative sign is required for crossed box amplitudes when $χ^{θ0}_1$ is in the loop, since $(χ^{θ0}_1)^∗ = -χ^{θ0}_1$ for that particular $i$.

Neutralino box:

$$C^θ_1 = \frac{α_w}{2m^2_{χ^0_1}} \left[ G^1_L G^1_L G^1_L G^1_L G^1_k \bar{G}^i(χ^θ_0, S) \delta^3_{dLL} \right] (\delta^3_{dLL})^2,$$

$$C^θ_2 = \frac{α_w}{m^2_{χ^0_1}} \left[ H^1_L H^1_L G^1_L \bar{G}^1_L \right] \left[ \sqrt{\frac{m_{χ^θ_0}}{m^2_{χ^θ_0}}} F^i(χ^θ_0, S) \delta^3_{dLL} \right] (\delta^3_{dLL})^2,$$

$$C^θ_3 = \frac{α_w}{m^2_{χ^θ_0}} \left[ H^1_L H^1_L G^1_L \bar{G}^1_L \right] \left[ \frac{m^2_{χ^θ_0}}{m_{χ^θ_0}} F^i(χ^θ_0, S) \right] \left[ \sqrt{\frac{m_{χ^θ_0}}{m^2_{χ^θ_0}}} \delta^3_{dLL} \right] (\delta^3_{dLL})^2.$$

B. Impact on $B_d$ Mixing

To obtain $∆ m_{B_d}$, we use $∆ m_{B_d} = 2|m_{12}|$, where

$$M^B_{12} = |M^1_{12}| e^{iφ_4} = |M^SUSY_{12}| e^{iφ_{4SUSY}},$$

and

$$M^{SM}_{12} = 0.33 \left( \frac{f_{B_d} |B_{B_d}|}{230 \text{ MeV}} \right)^2 \left( \frac{|V_{td}|}{8.8 \times 10^{-3}} \right)^2 e^{2iφ_1} \text{ ps}^{-1},$$

where $M^{SM}_{12}$ is the SM contribution, its value is well known [53]. The vacuum insertion matrix elements of $O_i$ are given in Ref. [11]. These matrix elements are modified by bag-factors to include non-factorizable effects. For simplicity, we assume the bag-factors for matrix elements of $O_{2-5}$ are equal to $B_{B_d}$, which is calculated for $O_1$. In the subsequent numerical analysis, we take $f_{B_d} |B_{B_d}| = (230 ± 40) \text{ MeV}$ [14]. For CKM matrix elements, we take $V_{ub}/V_{cb} = 0.41$ and $φ_3 = 65°, 85°$. We use $|V_{td}| × 10^3 = 8.0, 9.2$ to get $∆ m^{SM}_{B_d} ≈ 0.54, 0.72 \text{ ps}^{-1}$, respectively, which are close to the experimental value $∆ m^{SM}_{B_d} = 0.48±0.010 \text{ ps}^{-1}$ [12]. The large uncertainty in $f_B |B_{B_d}|$ makes it possible for $∆ m^{SM}_{B_d}$ to lie within experimental range even for large $φ_3$. However, when one considers $∆ m_{B_d}/∆ m^{SM}_{B_d}$, the hadronic uncertainty is reduced.

by leading order Taylor expansion with respect to squark mixing angles.
significantly, i.e. \( \xi \equiv f_B B^1/2/f_u B^1/2 = 1.16 \pm 0.05 \) from lattice [35], and the SM prediction for large \( \phi_3 \) case is not consistent with experiments, and New Physics would be needed for this case.

With the formulas above, we are ready to discuss the SUSY contributions to \( B \to \bar{B} \) mixing in the AHS models. In Fig. 2 we illustrate the \( m_\tilde{g} \) dependence of \( \Delta m_{Bd}^{SUSY} / (0.504 \text{ ps}^{-1}) \) from gluino box diagrams, where the denominator is the experimental bound at 2\( \sigma \). The solid lines correspond to contributions from RR-RR mixings where each squark propagator has an insertion of \( \delta^{13}_{dRR} \sim 0.18 \). The dashed lines correspond to contributions from LL-RR mixings where one squark propagator has an insertion of \( \delta^{13}_{dLL} \sim 0.18 \) and the other an insertion of \( \delta^{13}_{dRR} \sim 0.18 \).

We can tell from Fig. 2 which mixings, LL-RR or RR-RR, give the dominant contribution in different regions of parameter space. For small \( m_\tilde{g} \) the RR-RR mixings give stronger constraint, but for larger \( m_\tilde{g} \) the LL-RR mixings are more stringent. The parameter space corresponding to \( \Delta m_{Bd}^{SUSY} / (0.504 \text{ ps}^{-1}) \gg 1 \) is excluded. Since contributions from other sparticles are sub-dominant in most of the parameter space, as we will discuss later, they are unlikely to cancel the gluino contribution.

We clearly need TeV range gluino and/or squark masses to satisfy the \( \Delta m_{Bd} \) constraint. This comes as a result of the large mixings in right-handed sector, and can be shown by simple arguments. \( C_1^W \) in the SM is roughly proportional to \( \langle \alpha_W/m_W^2 \rangle^2 \langle V_W V_s \rangle^2 \). For SUSY contribution assuming \( m_\tilde{g} \sim \tilde{m} \), \( \langle \alpha_W/m_W^2 \rangle^2 \langle V_W V_s \rangle^2 \) factor is replaced by \( N_e \langle \alpha_Y \rangle^2/m_\tilde{g}^2 \), and \( \langle V_W V_s \rangle^2 \sim \lambda_0^2 \) is replaced by \( (\delta_{dRR}^{13})^2 \sim \lambda_0^2 \) and \( \delta_{dLL}^{13} \delta_{dRR}^{13} \sim \lambda_0^2 \) for RR-RR and LL-RR mixing contributions, respectively. Since the SM contribution is already close to the experimental observation, one requires

\[
\tilde{m} \gtrsim \left( \frac{\alpha_s}{\alpha_W} \right) \sqrt{\frac{N_e \lambda_0^4}{\alpha_s}} \left( \frac{m_W^2}{m_t} \right) \sim 1, 4 \text{ TeV},
\]

from LL-RR and RR-RR mixings, respectively. Thus the typical scale of superparticles, \( M_{SUSY} \), has to be large due to large squark mixings. Comparing with Fig. 2, we note that the LL-RR case is close to this estimate, while \( \tilde{m} \) for the RR-RR case is weaker than the estimate. This can be traced to the aforementioned cancellation in the RR-RR case where total cancellation in \( C_1^W \) is possible for \( x_{\tilde{g}q} \sim 2.43 \). Considering RR-RR mixings only, there is a valley in the parameter space that even light superparticles with masses less than 250 GeV are allowed.

Neutralino box diagrams are induced by the same flavor source as the gluino box diagrams. The neutralino masses could be related to \( m_\tilde{g} \) through a GUT-like relation on the gaugino Majorana masses [36].

\[
m_\tilde{g} = \frac{\alpha_s}{\alpha_W} M_2, \quad M_1 = \frac{5}{3} \frac{\alpha'}{\alpha_W} M_2.
\]

The gluino-neutralino box diagrams dominates over neutralino-neutralino box. The neutralino contribution to \( \Delta m_{Bd}^{SUSY} / (0.504 \text{ ps}^{-1}) \) is less than 10% of gluino contribution for \( m_\tilde{g} \gg 500 \text{ GeV} \) for either RR-RR or LL-RR mixings with \( |\tan \beta| = 2-50 \) and \( |\mu| = 100-1000 \text{ GeV} \). For low \( m_\tilde{g} \), \( |\mu| \) and large \( |\tan \beta| \), \( m_\tilde{g} \), its contribution, through \( C_1^W \), can be comparable with the RR-RR mixing induced gluino contribution, which is suppressed by large \( m_\tilde{g} \). However, as can be seen from Fig. 2, both the RR-RR and LL-RR mixing induced gluino contributions already give \( \Delta m_{Bd}^{SUSY} / (0.504 \text{ ps}^{-1}) \gg 1 \) in that region, and we need to fine-tune the relative phase and size of these mixings to be within experimental bound.

The charged Higgs contributes \( r = |M_{12}^{SUSY}/M_{12}^{SM}| \sim (37, 11, 3)% \) for \( m_{H^+} = (100, 400, 1000) \text{ GeV} \) with \( |\tan \beta| = 2 \), and \( \sim (0.3, 0.2, 0.1)% \) for \( |\tan \beta| = 50 \). The charged Higgs contribution interferes coherently with the SM amplitude. The result is consistent with early studies [40]. Without further interference from other SUSY contributions the \( b \to s \gamma \) branching ratio constrains charged Higgs mass to be quite large [57]. In this case, the charged Higgs contribution to \( B^{00} \) mixing is small. It is known that cancellations from other particle contributions may reduce the bound on the charged Higgs mass [49, 55]. We will return to this point in Section VI. On the other hand the chargino gives even smaller contribution, \( r \sim 0.1, 0.5% \) for TeV range \( m_\tilde{g} \) and \( \mu = \pm 1000, \pm 100 \text{ GeV} \), respectively. The main contribution comes from the term with \( \tilde{Y}_t^T \delta_{dRR}^{13} \) in \( C_1^W \). The ratio is reduced by \( \sim 30\% \) for large \( |\tan \beta| \) mainly due to the reduction of \( \tilde{Y}_t^2 \). The smallness of chargino contributions is in contrast with early studies [40]. In our case there is no large mixing involving third generation in the up type sector, and there is no large splitting due to light stop. In this type of models, the SUSY contribution to \( B^{00} \) mixing is dominated by gluino exchange.
After showing that gluino exchange gives dominant contributions to $M_{SUSY}$, we turn to explore the interference between $M_{12}^{SUSY}$ and $M_{12}^{SM}$. We consider the SUSY phase $\phi \equiv \arg \delta_{dRR}^{13}$. The experimental measurement of $\Phi_{B_d}$ is no longer just $\phi_1$ of SM. For illustration we plot, in Figs. 3(a) and (b), $\Delta m_{B_d}(\equiv 2|M_{12}^B|)$ and $\sin 2\Phi_{B_d}$ vs. $\phi$, respectively, for 1.5 TeV squark mass and 1.5 TeV gluino mass. In Fig. 3(a), the solid (short-dashed) lines correspond to $\phi_3 = 65^\circ$ (85$^\circ$). The upper and lower solid and short-dashed lines denote the 1$\sigma$ boundaries of $f_{B_d}\hat{B}_{B_d}^{1/2} = (230 \pm 40)$ MeV. If RR-RR mixings dominate, the SUSY phase $\phi_{SUSY} \sim 2\phi$, while if LL-RR mixings dominate, then $\phi_{SUSY} \sim \phi$. For $\phi_3 = 65^\circ$ (85$^\circ$), the SUSY model gives $r \sim 22\%$ (16$\%$) from RR-RR mixings, and $r \sim 55\%$ (41$\%$) from LL-RR mixings. These are consistent with Fig. 2, since $r \sim \Delta m_{B_d}^{SUSY}/(0.504 \text{ps}^{-1})$. The two SUSY contributions interfere constructively (destructively) for $\phi \sim \pi (0)$. In the present case, LL-RR mixings dominate over RR-RR mixings and show a $\phi_{SUSY} \sim \phi$ behavior in the graph.

In Fig. 3(a) and (b), we show the same physics measurable but with $\bar{m} = 3$ TeV, $m_3 = 1.5$ TeV. For $\phi_3 = 65^\circ$ (85$^\circ$), the SUSY model contributes $r \sim 29\%$ (22$\%$) from RR-RR mixings and $r \sim 18\%$ (13$\%$) from LL-RR mixings vs $\Delta m_{B_d}^{SM}$. These are again consistent with Fig. 2. The interference pattern is the same as the previous case. But in the present case, the RR-RR mixing contribution dominates over LL-RR mixings hence show a $\phi_{SUSY} \sim 2\phi$ behavior in Fig. 3(a). We can see from Fig. 2 that RR-RR mixings dominate for the point $(m_\tilde{g}, \bar{m}) = (1.5, 3)$ TeV.

We see that $\sin 2\Phi_{B_d}$, as measured from $B_d \rightarrow J/\psi K_S$ can range from 0.1–0.95 and 0.4–0.9 as shown in Figs. 3(b) and 3(b), respectively. These curves are obtained by overplotting various segments from the corresponding curves within 1$\sigma$ range of $f_{B_d}\hat{B}_{B_d}^{1/2}$. For example, in Fig. 3(a), any single line with $\bar{m} = m_\tilde{g} = 1.5$ TeV that corresponds to a value within 1$\sigma$ of $f_{B_d}\hat{B}_{B_d}^{1/2}$ should lie within the two solid lines. Each single line only intersects with the 2$\sigma$ experimental range for $\Delta m_{B_d}$ for some region of $\phi$ and corresponds to some segment in the solid line of Fig. 3(b) (one can compare with Fig. 1(b) of Ref. [46], where $f_{B_d}\hat{B}_{B_d}^{1/2} = 200$ MeV is used). Taking the uncertainty of $f_{B_d}\hat{B}_{B_d}^{1/2}$ into account enlarges the parameter space considerably. The fact that these segments from different value of $f_{B_d}\hat{B}_{B_d}^{1/2}$ lie on a line and not forming a band corresponds to our simplifying assumption of using same bag-factor for all $\mathcal{O}_i$. Thus, they can be factored out without affecting the argument of $M_{12}$ in Eq. (47).

In SM, we have $\sin 2\Phi_1 \sim 0.75$–0.71 for $\phi_3 = 65^\circ$–85$^\circ$. The measurements from Belle and BaBar in year 2000 indicated smaller values for $\sin 2\Phi_1$ vs SM, and would correspond to relatively specific $\phi \equiv \arg \delta_{dRR}^{13}$ phase values (between $\pi/2$ and $\pi$) in Figs. 3(b) and 3(b). More recent definitive measurements from BaBar and Belle, in year 2001, give $\sin 2\Phi_1$ values close to 1! The average of BaBar and Belle 2001 values is given in Eq. (47). It is rather intriguing that this range corresponds to the main parameter space allowed by $B_d$ mixing, suggesting that SUSY contributions could be comparable to SM. In particular, if the Belle 2001 central value of $\sin 2\Phi_1 = 0.99$ holds up, it would imply that $\phi \equiv \arg \delta_{dRR}^{13}$ is between...
3π/2 and 2π. The lighter squark mass case of Figs. 3 is preferred, but the heavier squark mass case of Figs. 4 is also possible.

To conclude this section, we note that one requires heavy squark and gluino masses to satisfy ∆m_{BB} bound. While the direct search of such heavy superparticles become less promising, these particles can show their effect in B_d–B_d mixing phases, even with TeV masses.

C. Brief Discussion on B_s Mixing and CP Phase

As shown in previous section, one could have s flavor decoupled and the above discussion is applicable, and is being tested right now. Alternatively, and mutually exclusive to the above case, it could be the d flavor that is decoupled, and the B_d system would be SM-like, which may still turn out to be the case in 2002. If so, B_s mixing may be the place where SUSY AHS effects show up.

With the CKM like relation \( \delta^{13}_{dRR}/\delta^{23}_{dRR} \sim V_{td}/V_{ts} \sim \lambda \) from Eqs. (17), (18), the B_d mixing case is rather similar to the discussion of B_d mixing, so long that the mass insertion approximation can be used. One can just scale up from previous discussion. The gluino contribution to \( B_s^0–\bar{B}_s^0 \) mixing is discussed in Ref. [17] and [18], where in the latter work the mass insertion approximation is relaxed. Since \( s_{BK}–\bar{B}_R \) mixing \( \sim 1 \) in SUSY AHS model, one in principle could have a relatively light “strange-beauty” squark, which, unlike the heavy SUSY scale that is the focus of this paper, can impact on direct search. We will discuss this case later in Sec. VII.

The B_s mixing phase may not be vanishingly small as in SM, and can be searched for at the Tevatron collider in a matter of years. The most interesting situation would be to find (soon!) \( \Delta m_{B_s} \), not far above SM expectation, but with large sin 2Φ_{B_s}.

IV. CONSTRAINTS FROM \( K^0–\bar{K}^0 \) MIXING

As shown in Eqs. (14) and (17). AHS models not only give large mixing in RR sector involving third generation down squark, they also give large mixing in 1-2 generations. It is well known that \( \Delta m_K \) is much smaller than \( \Delta m_{B_s} \) hence offers a much stronger constraint, while \( \varepsilon \) is even tighter. They make \( \delta^{12}_{dLL,RR} \sim \lambda \) impossible to sustain even with \( \tilde{m}, m_{\tilde{g}} \gtrsim \text{TeV} \). The formulas for koon mixing are similar to that for \( B_q–\bar{B}_q \) mixing, with only some modifications needed. For charged Higgs exchange diagrams, Eq. (33) now becomes,

\[
C_1^H = \frac{\alpha_W^2}{8m_W^2} V_{ts} V_{td} V_{ts}^* \frac{V_{td}^*}{x_i \omega_{x_j} H x_j} \cot^4 \beta \frac{1}{4} G(x_i H x_j H) + 2x_i x_j \omega_{x_i} \cot^2 \beta (F'(x_i W x_j W) + \frac{1}{4} G'(x_i W x_j W x_H W)),
\]

where \( i, j \) are generation indices of up type quarks and summed over. Other terms are neglected due to the smallness of quark masses \( m_{u,d} \). Chargino contributions are modified by changing \( V_{i\gamma} V_{n\omega} \delta_{u,LL}^m \) in Eq. (34) to \( V_{i\gamma} V_{n\omega} \delta_{u,LL}^m \) and neglecting other terms. Gluino and neutralino contributions are modified by changing \( \theta \) in Eqs. (39), (42), (44) to \( \delta \) and neglecting all \( H_{L,R} \) terms. The QCD running formula is also modified accordingly.

The charged Higgs contributions are in general small. For gluino and neutralino contributions we show in Table II the limits on \( \sqrt{\text{Re} \delta_{dLL}^2 \delta_{dRR}^2} \) and \( \sqrt{\text{Im} \delta_{dLL}^2 \delta_{dRR}^2} \) from \( \Delta m_{K}^{\text{SUSY}} < 3.521 \times 10^{-12} \text{MeV for } \tilde{m} = 1.5 \text{ TeV,} \)

where \( A, B, C, D = L, R \). The constraints from \( \varepsilon \) can also be estimated by using

\[
|\varepsilon_K| = \frac{\text{Im} M_{12}}{\sqrt{2}\Delta m_K} < 2.268 \times 10^{-3}.
\]

For different values of \( \tilde{m} \), the limits can be roughly obtained by multiplying a factor \( \tilde{m}/(1.5 \text{ TeV}) \). We take the hadronic scale in Eq. (40) to be \( \sim \text{GeV} \) and \( a_\mu (M_2) = 0.1185 \). We can reproduce the results of Ref. [23] by using \( a_\mu (\mu) \sim 1 \) and \( \tilde{m} = 500 \text{ GeV} \) and by considering gluino contributions only. The QCD effects enhance the SUSY contributions by a few times for contributions arising from \( \delta_{dLL,RR}^\epsilon \), as one can see from \( \eta_4 \sim 6 \). Since the most severe constraint is on \( \delta_{dLL,RR}^\epsilon \), the QCD effects make it more stringent [53]. From Eq. (63), \( \varepsilon \) gives even more stringent constraint than \( \Delta m_K \). However, the constraint becomes less severe if the phases of \( \delta \) are small (of order 0.01). Together with constraint from electric dipole moment of neutron, one may be led to the idea of approximate CP (see for example [53]).

From the above discussion, it is clear that \( \delta_{dLL,RR} \sim \lambda \) cannot be sustained. One has to invoke QSA as discussed in Sec. II to impose appropriate “texture zeros”. We see that the values given in Table I are all well below the limits from \( \Delta m_K \) and \( \varepsilon \) constraints, Table II, even with \( \mathcal{O}(1) \) phases. The approximate CP assumption can be relaxed.
As noted in Sec. II, QSA will induce $\delta_{uLL}^2 \sim \lambda$ by shifting the source of the Cabibbo angle to the up-type sector. The full strength $\delta_{uLL}^2$ can contribute to kaon mixing via chargino diagrams. In Fig. 5, we show the parameter space constrained by $\Delta m_K$. We use the GUT relation on gaugino mass as given in Eq. (50) for sake of simplicity and definiteness. The horizontal axis can be converted to wino mass by multiplying $m_{\tilde{g}}$ by $\sim 0.4$. For the $\varepsilon_K$ constraint, we would need $\arg(\delta_u)$ to be less than 0.1. We see from Fig. 5 that the kaon mixing constraint also points to TeV scale gluino and squarks. We stress that this is a generic feature of QSA and has nothing to do with the choice of retaining $M_3^{\tilde{d}}$ or not. Keeping $M_3^{\tilde{d}}$ leads to interesting low energy physics, even with TeV scale particles as a result of kaon mixing constraints.

We note that it is possible for the chargino diagrams to interfere destructively with LL (or RR) mixing induced gluino and neutralino contributions, and one can satisfy the kaon constraint with a lower mass scale. However, since these correspond to different set of parameters, it is unlikely for the interference to be destructive in general.

**V. IMPLICATIONS FOR $D^0-\bar{D}^0$ MIXING**

The experimental situation for $D^0-\bar{D}^0$ mixing is rather volatile at the present time. A search by the CLEO Collaboration gives $1/2x_D^2 < 0.041\%$ and $-5.8\% < \gamma_D < 1.0\%$ [15], where $x_D'$ and $y_D$ are defined in Eq. (3). CLEO further adopted $\delta_D \simeq 0$ from model arguments to reach a more stringent bound of $x_D \simeq x_D' < 2.9\%$. If $\delta_D \neq 0$ [23], however, the preferred negative value of $y_D'$ may in fact be hinting at $x_D \sim \lambda$.

Another approach is to compare $D^0 \rightarrow K^-\pi^+$ and $K^-K^+$ decays and measure the lifetime difference between $CP$ even and odd final states. The current world average from Belle [13], BaBar [17], CLEO [17], E791 [14] and FOCUS [16] Collaborations is $1.1 \pm 0.99\%$ [22]. Since this is consistent with zero and does not support the nonzero claim by FOCUS, we shall take the more stringent constraint of $x_D < 2.9\%$ from CLEO in the following. What we find intriguing, however, is that $\delta_{uLL}^2 \sim \lambda$ with $\bar{m}$, $m_{\tilde{g}} \sim$ TeV brings $x_D$ right into the ball-park of the % level! Furthermore, this can be probed in detail in the next few years at the B factories, and in the longer run, by hadron collider detectors.

We consider gluino and neutralino exchange diagrams induced by up-squark mixing with $\mu = 1$ TeV and $\tan \beta = 2$ with formulas similar to $B$ mixing. The dependence on $\tan \beta$ is weak. The SUSY contribution from gluino box to $x_D/0.029$ is illustrated in Fig. 6 in the $m_{\tilde{g}}-\bar{m}$ plane. The solid lines correspond to $\delta_{uLL}^2 \sim 0.18$, while the dashed lines correspond to $\delta_{uLL}^2 \delta_{uRR} \sim (0.18)^{5.5}$. It is clear that LL-LL induced gluino box diagrams dominate. In Fig. 6 we illustrate $x_D$ vs. $\bar{m}$ for $m_{\tilde{g}} = 0.8, 1.5, 3$ TeV, respectively. As in the $B$ mixing case, there is a narrow valley from $\delta_{uLL}$ induced gluino contributions around $m_{\tilde{g}}^2/\bar{m}^2 \sim 2.43$ when $C_2^D$ of Eq. (24) vanishes. This could make the parameter space from Fig. 6 too restrictive. However, the actual zeros in Fig. 6, occur at slightly shifted mass ratios, reflecting a cancellation between various contributions from $\delta_{uLL}$ and $\delta_{uLL}\delta_{uRR}$ when they have a common phase. Although $\varepsilon_K$ constrains $\arg(\delta_{uLL}^2)$ to be less than 0.1, the phase of $\delta_{uLL}^2$...
is not constrained since \( \delta_{LR}^{12} \) is by itself small. In general, the SUSY phase \( \delta_{LR}^{12} \) does not have to vanish, and having phase in common with \( \delta_{LL} \) is not likely. Thus, the deep valley would in general be filled, but the figure illustrates the adjustability of \( x_D \). It also gives an explicit example where detectable \( D^0 \) mixing would likely carry a CP violating phase.

To conclude this section, we note that AHS with QSA is known to produce large \( D \) meson mixings. The stringent upper bound on \( x_D \) seem to provide severe constraint for QSA models. This is more or less true when squarks and gluino are as light as a few hundred GeV. However, as a result of the large mixing in squark sector in our case, the proximity of \( \Delta m_{Bu} \) to SM expectation leads to squarks at TeV scale, and sin \( 2\phi_1 \) may be affected in an interesting way. It is interesting that the scale determined from this leads to a \( D^0 \) meson mixing close to experimental hints. We eagerly await the experimental situation to clear up, i.e. whether the CLEO hint is due to \( \Delta \Gamma_D \) or \( \Delta m_D \).

VI. RADIATIVE \( B \) DECAYS

The effective Hamiltonian for \( b \rightarrow q \gamma, \, qg \) transitions, where \( q = d \) or \( s \), is given by

\[
H_{\text{eff.}} = \frac{G_F m_b}{\sqrt{2}} \frac{V_{tb} V_{tq}^*}{4 \pi^2} \left\{ e \left[ C_{7,\gamma} R + C_{7,\gamma}^* L \right] F^{\mu\nu} + \sum_{i=1}^{N} \left[ C_{8q} R + C_{8q}^* L \right] T^a G_{\mu\nu}^{a} \right\} \sigma_{\mu\nu} b,
\]

where we have neglected \( m_q \), \( C_{7,\gamma} = C_{7,\gamma}^{SM} + C_{7,\gamma}^\text{New} \) are the sum of SM and New Physics contributions, while \( C_{7,\gamma}^\text{New} \) come purely from New Physics. We are particularly interested in the case where \( C_{7,\gamma,8q} \) are large. The effects from the SUSY contributions are given by

\[
C_{7,\gamma}^{\text{New}} = C_{7,\gamma}^{(i)H^-} + C_{7,\gamma}^{(i)H^+} + C_{7,\gamma}^{(i)\tilde{H}^+} + C_{7,\gamma}^{(i)\tilde{H}^0},
\]

\[
C_{8q}^{\text{New}} = C_{8q}^{(i)H^-} + C_{8q}^{(i)H^+} + C_{8q}^{(i)\tilde{H}^+} + C_{8q}^{(i)\tilde{H}^0}.
\]

The Feynman diagrams are shown in Fig. 8.

A. Formulas

Charged Higgs Exchange:

\[
C_{7,\gamma}^{\text{H}} = -\frac{x_{1H}}{2} \left\{ \cot^2 \beta \left[ Q_u F_1(x_{1H}) + F_2(x_{1H}) \right] + \left[ Q_u F_3(x_{1H}) + F_4(x_{1H}) \right] \right\}, \tag{56}
\]

\[
C_{8q}^{\text{H}} = -\frac{x_{1H}}{2} \left\{ \cot^2 \beta F_1(x_{1H}) + F_3(x_{1H}) \right\}, \tag{57}
\]

where \( F_i(x) \) are loop functions and the explicit expressions can be found in Ref. [19].

Gluino Exchange:

\[
C_{7,\gamma}^{\text{Gluino}} = \frac{\pi \alpha_s}{\sqrt{2} G_F V_{tb} V_{td}} \frac{Q_d 2 C_2(R)}{\tilde{m}^2} \times \left\{ \delta_{dLL}^{13} g_2(x_{\tilde{g}\bar{q}}) - \frac{m_d}{m_b} \delta_{dLR}^{13} g_4(x_{\tilde{g}\bar{q}}) \right\}, \tag{58}
\]

\[
C_{8q}^{\text{Gluino}} = \frac{\pi \alpha_s}{\sqrt{2} G_F \tilde{m}^2 V_{tb} V_{td}^*} \times \left\{ \delta_{dLL}^{13} \left[ 2 C_2(R) - C_2(G) \right] g_2(x_{\tilde{g}\bar{q}}) - C_2(G) g_1(x_{\tilde{g}\bar{q}}) \right\} + C_2(G) g_3(x_{\tilde{g}\bar{q}}) \right\}, \tag{59}
\]

where \( Q_d \) is the down quark electric charge, \( C_2(G) = N = 3 \) and \( C_2(R) = (N^2 - 1)/(2N) = 4/3 \) are Casimirs, and the functions \( g_i(x) = -d/dx [x F_i(x)] \), i.e.

\[
g_1(x) = \frac{1 + 9x - 9x^2 - x^3 + 6x(1 + x)}{6(x - 1)^5} \ln x,
\]

\[
g_2(x) = \frac{1 - 9x - 9x^2 + 17x^3 - 6x^2(3 + x)}{12(x - 1)^5} \ln x,
\]

\[
g_3(x) = \frac{5 - 4x - x^2 + (2 + 4x)}{2(x - 1)^4} \ln x,
\]

\[
g_4(x) = \frac{-1 - 4x + 5x^2 - 2x(2 + x)}{2(x - 1)^4} \ln x. \tag{60}
\]

The chirality partners \( C_{7,\gamma,8q} \) are obtained by interchanging \( L \) and \( R \) in the \( \delta \)'s. We see that \( \delta_{LL} \) and \( \delta_{LR} \) contribute to \( C_{7,\gamma,8q}^R \), while \( \delta_{RR} \) and \( \delta_{RL} \) contribute to \( C_{7,\gamma,8q}^I \).
There is an enhancement factor $m_\tilde{q}/m_b$ that comes with $\delta_{LR, RL}$. The factor $m_b$ is from normalizing with respect to the SM result, and hence the smallness of the $b$ quark mass with respect to the gluino mass is the origin of this enhancement. Such enhancement was noted in our earlier study \cite{30} of $\tilde{s}-\tilde{b}$ mixings where $d$ sector was decoupled completely. It has also been invoked to generate $\varepsilon'/\varepsilon$ via an analogous $\delta_{LR}^2$ term \cite{34,30} under a horizontal $U(2)$ (hence non-Abelian) symmetry model. The mechanism is generic and has been discussed in Ref. \cite{31}, but SUSY with LR squark mixings gives a beautiful example.

**Chargino Exchange:**

$$C_{\chi^+_{LR}, \chi^-} = \frac{m_w^2}{m_t^2 V_{tb}^* V_{tb}} \times \left\{ \left[ V_{j1} V_{j1}^* V_{mb}^\delta_{d,lL}^\dagger - V_{j1} V_{j2}^* \tilde{Y}_t V_{tb}^\delta_{d,lL}^\dagger \right] + V_{j2} V_{j2}^* \tilde{Y}_t V_{tb}^\delta_{d,lL}^\dagger \right\} +$$

$$V_{j2} U_{j2} \tilde{Y}_t V_{tb}^\delta_{d,lL}^\dagger \left[ \left[ g_1(x_{\tilde{u}} - \tilde{q}) + Q_u g_2(x_{\tilde{u}} - \tilde{q}) \right] + m_{\tilde{u}} \left[ g_3(x_{\tilde{u}} - \tilde{q}) + Q_u g_4(x_{\tilde{u}} - \tilde{q}) \right] \right],$$

(61)

$$C_{\chi^+_{LR}, \chi^-} = \tilde{Y}_q \frac{m_w^2}{m_t^2 V_{tb}^* V_{tb}} \times \left\{ U_{j2} U_{j2}^* \tilde{Y}_t V_{tb}^\delta_{d,lL}^\dagger \left[ g_1(x_{\tilde{u}} - \tilde{q}) + Q_u g_2(x_{\tilde{u}} - \tilde{q}) \right] + \right.$$

$$\left. + m_{\tilde{u}} \left[ g_3(x_{\tilde{u}} - \tilde{q}) + Q_u g_4(x_{\tilde{u}} - \tilde{q}) \right] \right\} +$$

$$V_{j2}^* U_{j2}^* \tilde{Y}_t V_{tb}^\delta_{d,lL}^\dagger - U_{j2}^* V_{j2}^* V_{mb}^\delta_{d,lL}^\dagger \right\},$$

(62)

where as before we sum over $l$, $m$ for three generations and $j$ for two chargino mass eigenstates. $C_{\chi^+_{LR}, \chi^-$ can be obtained by dropping $g_1, g_3$ from above equations and replacing $Q_u$ with 1. It is clear from these equations that $C_{\chi^+_{LR}, \chi^-}$ is suppressed by $\tilde{Y}_q$.

**Neutralino Exchange:**

$$C_{\chi^0_{LR}, \chi^-} = -\frac{Q_d m_w^2}{m_t^2 V_{tb}^* V_{tb}^*} \left\{ 2 G_{qL}^j G_{bL}^j \delta_{d,lL}^{33} g_2(x_{\tilde{u}}) \right.$$}

$$+ m_{\tilde{u}} \left[ -2 G_{qL}^j G_{bR}^j \delta_{d,lL}^{33} + \sqrt{2} G_{qL}^j H_{bL}^j \delta_{d,lL}^{33} \right] g_4(x_{\tilde{u}}) \right\},$$

(63)

where $j$ is summed over four neutralino mass eigenstates, and $C_{\chi^0_{LR}, \chi^-}$ can be obtained by replacing $Q_d \rightarrow 1$ in the above equation. We neglect terms with $H_{bL}^j$ and some LR mixing terms when there is no chiral enhancement. Similar to the gluino case, the chiral partners $C'$ are obtained by taking a conjugation in the chirality $L \leftrightarrow R$ and noting that $G_{L,R} \leftrightarrow - G_{R,L}$.

When running down to the B decay scale $\mu \approx m_b$, the leading order Wilson coefficients $C^{\text{SM}}_{LR}$ are given by \cite{32}.

$$C_{\chi^+_{LR}, \chi^-}(\mu = m_b) = -0.31 + \eta_7 C^{\text{SM}}_{\chi^+_{LR}, \chi^-}(M_{\text{SUSY}})$$

$$+ \frac{8}{3} (\eta_8 - \eta_7) C^{\text{SM}}_{\chi^+_{LR}, \chi^-}(M_{\text{SUSY}}),$$

(64)

where,

$$\eta_7 = \left( \frac{\alpha_s(M_{\text{SUSY}})}{\alpha_s(m_t)} \right)^{16/21} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_b)} \right)^{16/23},$$

$$\eta_8 = \left( \frac{\alpha_s(M_{\text{SUSY}})}{\alpha_s(m_t)} \right)^{14/21} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_b)} \right)^{14/23},$$

(65)

while for opposite chirality, which receives no SM contribution, one simply replaces $C^{\text{SM}}_{\chi^+_{LR}, \chi^-}$ by $C^{\text{SM}}_{\chi^+_{LR}, \chi^-}$ and set the constant terms to zero.

**B. Phenomenological Impact**

It is well known that $B \to X_s \gamma$ is a severe constraint on New Physics. The current experimental results are $Br(B \to X_s \gamma) = (2.85 \pm 0.35 \pm 0.22) \times 10^{-4}$ \cite{65}, $(3.37 \pm 0.53 \pm 0.42 \pm 0.50_{\text{model}}) \times 10^{-4}$ \cite{66}, $(3.11 \pm 0.82 \pm 0.72) \times 10^{-4}$ \cite{66}, from CLEO, Belle and ALEPH, respectively. It is known that the charged Higgs contribution interferes constructively \cite{24} with SM contribution $C^{\text{SM}}_{\chi^+_{LR}, \chi^-}(m_b) = -0.31$. By using Eq. (61), we have,

$$C_{\chi^+_{LR}, \chi^-} = 1 + (35\%\%, 22\%, 15\%, 11\%, 9\%\%),$$

(66)

for $m_{H^\pm} = (400, 600, 800, 1000, 1200)$ GeV. The rate can get enhanced by $\sim 80\%-20\%$. The result is insensitive to $\tan \beta$ within $2-50$, since the term without cot $\beta$ in Eq. (64) is dominant \cite{77}. If we require that the deviation from the SM rate to be less than 20\%, which is close to experimental error, we need $m_{H^\pm} \geq 1.2$ TeV, or we may need cancellations from other particles \cite{13}. The chargino contribution may partially cancel the charged Higgs contribution \cite{58}. This mechanism is still operative even if we have $m_{\tilde{g}}, m = 1.5$ TeV.

$$C_{\chi^+_{LR}, \chi^-} = 1 \pm \left\{ \begin{array}{c|c}
13 (0.6) & 12 (0.6) \\
17 (0.9) & 3 (0.8) \\
18 (0.5) & 0.3 (0.4) \\
8 (0.05) & 0.2 (0.2) \\
2 (0.1) & 0.1 (0.1)
\end{array} \right\} \left( \frac{\delta_{LR}^{33}}{m_t/m} \right) \left( \frac{V_{tb}^* V_{mb}^{\dagger \delta_{d,lL}^{33}}}{\lambda^2} \right),$$

(67)
for $\mu = \pm (100, 500, 1000)$ GeV, and $\tan \beta = 50$ (2). The $\delta_{uRL}^{3}$ term comes from chiral enhancement (with $m_{\tilde{C}} / m_{b}$ factor). The sign of the coefficient in front of $\delta_{uRL}^{3}$ is the same as the sign of $\mu$. The coefficient for $|\mu| = 500$ GeV is greater than $|\mu| = 100$ GeV, because of chiral enhancement. The coefficient will drop below 0.1 for $|\mu| \geq 1.2$ TeV. LR and RL mixings without chiral enhancement are negligible, their contributions being only about $10^{-5} - 6$ of $C_{7}^{SM}$. The term with $V_{ts}^{*} V_{tb} \delta_{uLL}^{3}$ is also from the chiral enhancement term, while the last term is not chirally enhanced. Due to the smallness of $Y_{\gamma}$, $|C_{7}^{T_{\gamma}} / C_{7}^{SM}|$ is below 1% for expected mixing angles. For $\tan \beta = 2$, terms are suppressed by the $Y_{\gamma}(\tan \beta = 2) / Y_{\gamma}(\tan \beta = 50) \sim 1/22$ factor, except for the last term.

The sign of RL stop mixings is anti-correlated to $\mu$. This can provide needed cancellations for low $m_{H^{-}}$, even if $M_{2}$ and $\tilde{m}$ are large. For $|\tilde{m}| = 100, 500$ GeV, if the signs of second and third terms of Eq. (67) are negative, the charged Higgs mass can be as low as 300, 400 GeV, for a rate within 20% from SM expectation. Even for $|\mu| = 1$ TeV, the cancellation may lower $m_{H^{-}}$ to 600 GeV. The cancellation, however, requires some degree of fine-tuning. Without such cancellations, we would need to require $|\mu|, m_{H^{-}} > 1$ TeV if we allow deviations from the SM rate to be within 20%.

For $b \to d \gamma$ decay, we should replace $V_{ts}^{*} V_{tb}$ and $\lambda^{2}$ in Eq. (67) by $V_{td}^{*} V_{tb}$ and $\lambda^{2} e^{i \phi_{t}}$, respectively, while Eq. (66) remains unchanged. For $V_{ts}^{*} V_{tb} \delta_{uLL}^{3}$ real and negative, it would cancel against charged Higgs contribution for low $\mu$. However, this need not be the case. For example, for $|\mu| = 100$ GeV, $m_{H^{-}} = 400$ GeV, the cancellation mechanism gives $Br(B \to X_{\gamma})$ within experimental error, while $Br(b \to d \gamma)$ is enhanced by a factor of 2. But in both $b \rightarrow s \gamma, d \gamma$ decays, the chargino contribution to asymmetry $a_{\gamma}^{\rho, \gamma}$ of Eq. (3) is within 2%.

To obtain large asymmetry $a_{\gamma}^{\rho, \gamma}$, we need a sizable $\sin 2 \theta_{\gamma}$ (see Eq. (3)) which requires $C_{7}^{\rho, \gamma}$ and $C_{7}^{\gamma}$ to be of comparable size. To achieve this the new Physics must have large $C_{7}^{\rho, \gamma}$ but a relatively small contribution to $C_{7}^{\gamma}$, since the latter already receives a large SM contribution. For example, in case of gluinos, we need large $\delta_{dRR}^{3}$ and/or $\delta_{dRL}^{3} m_{3} / m_{b}$ and small $\delta_{dLL}^{3}$ and $\delta_{dLR}^{3} m_{3} / m_{b}$. It is interesting that this indeed can be realized in AHS models. For similar reasons we do not expect a large modification of $C_{7}^{\gamma}$ from squark mixings as was noted in our discussion of charged Higgs effects.

To use the formula of $a_{\gamma}^{\rho, \gamma}$ we also need to know the phase. The SM phase in $B \to X_{\gamma}$ is rather complicated since $u$ and $c$ quark contributions at NLO are not CKM suppressed. However, as shown in [31] the NLO contribution are found to increase the rate by 10% and the long-distance contribution from intermediate $u$ quarks in the penguin is expected to be small. For exclusive modes, some model estimations (such as from Light-Cone QCD sum rule) give long-distance effect in $B \to \rho \gamma$ and $B \to \omega \gamma$ at about $O(15\%)$ [37]. Even though long distance physics may enter, it does not enhance $C_{7}^{\rho, \gamma}$ [37]. For charge $B$ decays the dominant long distance contribution come from weak annihilation diagrams (giving $|C_{7}^{\rho, \gamma} / C_{7}^{SM}| \sim 4\%$), which is, however, absent in the case of $B^{0} \rightarrow \rho \gamma$ decay.

In Fig. 9, we show gluino and neutralino contributions to the asymmetry coefficient $\sin 2 \theta_{\gamma}$ vs. $\tilde{m}$, with gluino mass $m_{\tilde{g}} = 1.5$ TeV. We use $\mu = 1$ TeV and $\tan \beta = 2$. The asymmetry is generated mainly from the RL mixing induced gluino penguins $\sim 8\%$ and can reach 10% when including RR mixing contributions. The asymmetry can be measured at B Factories, and at future hadron collider B detectors such as LHCb or BTeV. Since $\sin 2 \theta_{\gamma} \sim 2 |C_{7}^{\rho, \gamma} / C_{7}^{\gamma}|$ for small $\delta_{d}$, we obtain $|C_{7}^{\rho, \gamma} / C_{7}^{SM}|$ of about 4%, for $\tilde{m} \sim 1.5$ TeV, which is slightly larger than the estimation $\sim 2\%$ for long distance effects from charm penguin [38].

![Figure 9](image_url)

**FIG. 9.** Asymmetry coefficient $\sin 2 \theta_{\gamma}$ vs. $\tilde{m}$, for gluino mass $m_{\tilde{g}} = 1.5$ TeV, where solid (dashed) curve correspond to $\delta_{dLR}^{3}$ and $\delta_{dRR}^{3}$ having opposite (same) phase.

### C. Non-standard C-Term and $\tan \beta$ Enhancement

In a previous study, we obtain large or even maximal asymmetry rather easily in $b \rightarrow s \gamma$ with sub-TeV super-particle mass scale [39]. Here, the high SUSY scale as required by meson mixings leads to too severe a suppression in $1/G_{F} m^{2}$, as can be seen from Eqs. (58), (59). We find, however, that it is still possible to have large $a_{\rho, \gamma}$, when considering *non-standard soft breaking terms* [40]. Non-standard soft breaking terms can survive without inducing quadratic divergence if there is no gauge singlet particles in the low energy spectrum. For scales below the horizontal symmetry breaking scale and the masses of $S_{i}$, we are left with particles of the minimal supersymmetry standard model (MSSM). Therefore, by using a low energy effective theory of SUSY, it is legitimate to include these non-standard soft breaking terms, in particular, a non-holomorphic trilinear term, which is called the C-term.

Besides a (standard) A-term, $A_{d}(H_{d}) Y^{/} \tilde{D}_{L} \tilde{D}_{R}^{*}$, we now allow $(\tilde{M}_{d}^{2})_{LR}$ to have a non-standard C-term,
decay is defined as $d\gamma$ and $d\phi$, hence suppresses the asymmetry while enhancing the rate up to a factor 1.8. The enhancement factor for the third one is below 10%.

$C(H_0')Y'L_DD_R'$. It is natural that $A_d \sim C \sim \tilde{m}$, hence $(\tilde{M}_L^{21})_{LR} \sim \tilde{m} \tilde{M}_L^{21}$ tan $\beta$. In this way, one gains a tan $\beta \equiv |\langle H_0'\rangle/\langle H_d\rangle|$ enhancement factor, while $(\tilde{M}_L^{21})_{LR} \sim \tilde{m} \tilde{M}_L^{21}$ is unaffected. $M_q$, and hence $U_{qLR}$, is also unchanged, so the previous result for $D^0$ mixing remains unchanged. Some zeros in $(\tilde{M}_L^{21})_{LR}$ may also be lifted since these $C$-terms are no longer holomorphic, but they are still suppressed. We note that the $\delta_{dRL}$ contribution to kaon mixing remains protected by the smallness of $M_L^{21}/\tilde{m}$. Likewise, for $B_d$ and $B_s$ mixings, tan $\beta$ enhancement of $\delta_{dRL,RL}$ is insufficient to overcome $m_q/\tilde{m}$ suppression and $\delta_{dRL}$ still dominates.

We illustrate in Figs. 10 and 11 the ratio $\text{Br}(B \rightarrow X_d\gamma)/\text{Br}(B \rightarrow X_d\gamma)_{SM}$ and the coefficient sin $2\theta$ relevant for mixing dependent CP violation, with respect to the average squark mass $\tilde{m}$ for $m_{\tilde{b}} = 1.5$ TeV. The solid, dashed and dotted curves correspond to tan $\beta = 50$, 20 and 2, respectively. The branching ratio can be enhanced by a factor of 5 with respect to the SM value. This can be easily understood by noting that, before introducing the $C$-term, the RL mixing induced gluino diagrams give $|C_{7γ}/C_{7γ}^{SM}| \sim 4\%$ for $\tilde{m} = 1.5$ TeV. Adding the non-standard $C$-term enhances $\delta_{dRL}$ by tan $\beta$ and hence $|C_{7γ}/C_{7γ}^{SM}|$ is brought up to 0.04 tan $\beta$. A factor of 5 enhancement in rate follows for tan $\beta = 50$. Note that sin 2$\theta$ reaches maximum for $\tilde{m} \sim 2.6$ TeV. The reason is simply because $C_{7γ}$ dominates over $C_{7γ}$ for lower $\tilde{m}$ scale, hence suppresses the asymmetry while enhancing the rate significantly. Since the phase combination $\sin[2\Phi_B - \phi(C_{7γ} - \phi(C_{7γ})]$ in general should not vanish even if $\phi(C_{7γ})$ vanishes (because of non-vanishing $\Phi_B$), $a_{u,\rho,\gamma}$ could be sizable, which would unequivocally indicate the presence of New Physics.

The CP violating partial rate asymmetry $A_{CP}$ in $b \rightarrow d\gamma$ decay is defined as

$$A_{CP} = \frac{\Gamma(b \rightarrow d\gamma) - \bar{\Gamma}(\bar{b} \rightarrow \bar{d}\gamma)}{\Gamma(b \rightarrow d\gamma) + \bar{\Gamma}(\bar{b} \rightarrow \bar{d}\gamma)}$$

where $\bar{C}_{7γ}$ are coefficients for $\bar{b}$ decay. To have nonzero $A_{CP}$, apart from CP phases, one also needs absorptive parts. In the model under consideration, these can come only from the SM contribution with $u$ and $c$ quarks in the loop. The $A_{CP}$ is smaller than the SM one since $\delta_{LR}$ which contributes to $C_{7γ}$ is much smaller than $\delta_{RL}$, while there is no strong phase in $C_{7γ}$ to contribute to CP violation. Therefore New Physics only dilutes the $A_{CP}$ in this case by contributing to the total rate in the denominator of Eq. (68). That is, $A_{CP}$ is reduced by a rate enhancement factor for large tan $\beta$.

These figures hold also for $b \rightarrow s\gamma$ for the other choice of using Eq. (29) with $\tilde{s}\bar{b}$ but no $\tilde{d}\bar{b}$ mixing. Allowing for 20% rate uncertainty for the measured $\text{Br}(B \rightarrow X_s\gamma)$, we see that for the tan $\beta = 20$ case, $\tilde{m} \geq 3$ TeV is allowed, while for heavier squark $\tilde{m} = 5$ TeV the full range of $2 \lesssim \tan \beta \lesssim 50 \sim m_{\tilde{b}}/m_{\tilde{b}}$ is allowed. In these cases sin 2$\theta$ can go up to $\sim 0.6$. For lighter $\tilde{m}$ such as 1.5 TeV, the enhancement factor for large tan $\beta$ starts to break the good agreement between SM and the experimental value of $\text{Br}(B \rightarrow X_s\gamma)$, hence it seems one cannot have both large tan $\beta$ and $\tilde{m}$, $m_{\tilde{b}}$ too light (approaching TeV).

The interesting case of “strange-beauty” squark [48] where large mixing dependent CP is possible without non-standard C-terms is discussed in next section.

VII. DISCUSSION

We offer to discuss a few miscellaneous items.

A. Re($\epsilon'/\epsilon$) and EDM Constraints

Unlike Ref. [34], the AHS model presented here cannot be responsible for the large value of Re($\epsilon'/\epsilon$). This
is because of heavy masses and the suppressed value of \( \delta_{dLR,RL}^{12} \), as shown in Table II. However, \( \text{Re}(\varepsilon'/\varepsilon) \) does not provide further constraint on \( \delta' \), since we already satisfy the most severe case, \( \varepsilon_K < 2.268 \times 10^{-3} \). For further discussion on the issue of \( \text{Re}(\varepsilon'/\varepsilon) \) in the context of AHS models, see Ref. \[71\].

There are other horizontal symmetry models that lead to the pattern of quark mass ratios and mixings. In particular, the non-Abelian U(2) horizontal model \[72\] gives suppressed \( \delta_{dLR,RL}^{12} \) and can evade the kaon mixing constraint by a U(2) symmetry with relatively light masses. Mixing angles in fermion mass matrices are in general of the order of the square root of mass ratios \[72\].

\[
M_q = U_{qL}^{\dagger} M_{\text{diag}} U_{qR},
\]

\[
U_{qL,R} = \begin{pmatrix} \frac{1}{s_{qL,R}} & \frac{s_{qL,R}}{s_{qL,R}^2} & 0 \\ -\frac{s_{qL,R}^2}{s_{qL,R}} & \frac{1}{s_{qL,R}} & \frac{s_{qL,R}^2}{s_{qL,R}} \\ \frac{s_{qL,R}^2}{s_{qL,R}} & -\frac{s_{qL,R}}{s_{qL,R}^2} & \frac{1}{s_{qL,R}} \end{pmatrix},
\]

where

\[
s_{qL,R}^2 = \left( \frac{m_2}{m_3} \right)_q, \quad s_{qL}^2 = s_{qR}^2 = \sqrt{\left( \frac{m_1}{m_2} \right)_q}.
\]

The model leads to large \( \text{Re}(\varepsilon'/\varepsilon) \) \[72\]. The D–D mixing is small in the same way that the kaon mixing constraint is evaded. Mixing in \( \tilde{b} \sim \tilde{d} \) is relatively small, but the sparticle scale can be relatively light since kaon and D meson mixing constraints are evaded. This model may also lead to FCNC effects in B system \[73\].

Let us now consider the constraint from the electric dipole moment (EDM) of the neutron. It is well known that the EDMs of the electron and atoms give severe constraint on SUSY phases \[74\]. This is a common problem to all SUSY models and is quite independent from FCNC processes considered in this work. The problem should not be worse in our case, and in fact the TeV sparticle scale should loosen the constraint compared with usual considerations.

For the neutron EDM, there are contributions from electric dipole operator, the color dipole operator and the dimension six purely gluonic operator. We expect the first one to be dominant while the others may give comparable contributions and may in some cases loosen the constraint through cancellations \[76\]. In addition, there are two loop contributions \[74\]. These contributions could become important in the absence of large one loop contributions, such as in SUSY models with massive first and second generation squarks \[78\]. But since we have potentially large one loop contributions, we use only the electric dipole operator to estimate the order of magnitude bounds. For a more complete recent study of EDM constraint on SUSY models, see Ref. \[73\].

The neutron EDM can be expressed as

\[
\frac{d_n}{|e|} = \frac{2Q_{\alpha} m_2}{3m_1^2} g_4(x_{\tilde{g}|q}) \text{Im} \delta_{dLR}^{11},
\]

and for chargino \[75\],

\[
\frac{d_{\alpha}}{|e|} = \frac{2}{4\pi m_2^2} \left[ g_4(x_{\tilde{\chi}^+_{\alpha} q}) + (Q_\tilde{d} - Q_\alpha) g_3(x_{\tilde{\chi}^+_{\alpha} q}) \right] \text{Im} \eta_{\tilde{d}} \delta_{m}^{\chi^+,\chi^-},
\]

\[
\eta_{\tilde{d}}^{\chi^+} = -\tilde{Y}_\alpha V_{ij}^{*} \left( U_{ij}^0 V_{dLR}^{\dagger m} V_{am}^{*} - U_{ij}^0 V_{uLR}^{\dagger m} V_{um}^{*} \tilde{Y}_{am} \right),
\]

where \( l, m \) are summed over three generations of down type squarks. By interchanging \( Q_{\alpha, d} \), \( U, V_{uLR}^{\dagger m} V_{am} \), \( \tilde{Y}_{am} \) with \( Q_{\alpha, u} \), \( V, V^{*}_{uLR} \), respectively, one can obtain \( d_\alpha \). Finally, for neutralino contributions, one finds \[77\]

\[
\frac{d_{\nu}}{|e|} = \frac{2}{4\pi m_2^2} \left[ g_4(x_{\tilde{\chi}^0_{\alpha} q}) \text{Im} \eta_{\tilde{\nu}} \delta_{m}^{\chi^0},
\]

\[
\eta_{\tilde{\nu}}^{\chi^0} = -2G_{R}^{(5)} \tilde{Y}_{qR}^{\dagger l} \tilde{Y}_{lR}^{\dagger l} - \sqrt{2}G_{R}^{(5)} \tilde{Y}_{qR}^{\dagger l} H_{L}^{\dagger l} + \sqrt{2}H_{L}^{\dagger l} \tilde{Y}_{qR}^{\dagger l} H_{R}^{\dagger l},
\]

To illustrate the constraint on squark mixing phases we take chargino and neutralino mixing matrices to be real, and discuss the phase of \( \mu \) later. Requiring \( d_{\nu} \) to be less than the current experimental bound of \( 0.63 \times 10^{-25} \text{ e cm} \), we obtain

\[
|\text{Im} \delta_{dLR}^{11}| \leq (2.8, 3.4, 6.5) \times 10^{-6},
\]

\[
|\text{Im} \delta_{uLR}^{11}| \leq (5.5, 6.7, 12.6) \times 10^{-6},
\]

for \( \tilde{m} = 1.5 \text{ TeV}, |\mu| = 100–1000 \text{ GeV} \), and \( x_{\tilde{g}|q} = (0.3, 1, 4) \), respectively. These bounds are consistent with Ref. \[77\]. The bounds come dominantly from gluino contributions and hence insensitive to \( \tan \beta \) and \( |\mu| \). From chargino contribution alone with above parameter space and \( \tan \beta = 50 \) \[2\], we have

\[
|\text{Im} \left( \sum_{lm} V_{uLR}^{\dagger m} V_{dLR}^{\dagger n} \tilde{Y}_{am} \right)| \leq 0.43 - 0.54 (0.39 - 0.48),
\]

\[
|\text{Im} \left( \sum_{lm} V_{uLR}^{\dagger m} V_{dLR}^{\dagger n} \tilde{Y}_{am} \right)| \leq 0.11 (0.10).
\]

The AHS model gives \( |\text{Im} \delta_{dLR}^{11}| \sim m_4 A_{d}(1 + \{ \tan \beta \}) - \mu \tan \beta / \tilde{m}^2 \sim |m_4/\tilde{m}| \tan \beta \sim 8.4 \times 10^{-9} (\tan \beta/50) \). Thus, for large \( \tan \beta \), we need \( \arg(\delta_{dLR}^{11}) \) to be less than 0.1 to satisfy the EDM constraint. EDM from Mercury atom gives \( d_{H_{g}} < 2.1 \times 10^{-28} \text{ e cm} \). The bounds on \( |\text{Im} \delta_{uLR}^{11}| \) are one order of magnitude smaller than that from the neutron EDM bounds \[78\].
If $\mu$ is complex, it will contribute to $\arg(\delta_{dLR}^{11})$ as $-\arg(\mu)\tan\beta m_\mu/\mu^2$. For large tan $\beta$ and $|\mu| \sim \bar{m}$, we need $\arg(\mu)$ to be less than 0.1 (0.01) from the neutron (Mercury) EDM constraint. One should be more concerned, however, with the presence of $\delta_{dRL(11)}^{11}$ in Eq. (76). Take $\delta_{dRL}^{11}$ for example, we note that it is $\sim O(1)$ and is not suppressed by quark mass like $\delta_{dRL(11)}^{11}$. Thus, it will lead to a severe constraint on $\arg(\mu)$. For $m_\mu = \bar{m} = 1.5$ TeV, tan $\beta = 2$ and $|\mu| = 100$–1000 GeV, we need to have $\arg(\mu) \leq 0.03$–0.012. The bound is roughly inversely proportional to tan $\beta$. For large tan $\beta$, say 50, we need $\arg(\mu) \leq 5 \times 10^{-4}$ from the neutron EDM constraint. This constraint is quite severe, even for $m_\mu, \bar{m}$ at TeV scale. However, the very strong constraint on $\arg(\mu)$ from EDM consideration is a well known problem (see for example, Ref. [73]), and is not aggravated by considerations of FCNC induced by squark mixings, which has been the main focus of our study.

### B. Radiative $c \to u\gamma$ and $t \to c\gamma$ Decays

It is of interest to check radiative flavor changing neutral current processes in up type quark decays, since quark-squark alignment has shifted flavor violation to the up-type sector. It is well known that the short distance one-loop $c \to u\gamma$ amplitude is very small in the SM, due to the CKM suppression and the small $m_\mu^2/M_\mu^2$ factor. The amplitude can be raised by 2 orders of magnitude when one considers leading logarithmic QCD corrections involving operator mixings, and further raised by another 3 orders of magnitude when including non-CKM suppressed two-loop diagrams [31]. It is also known that long distance effects are in general large [32].

We can obtain SUSY contribution by using formulas similar to $b \to q\gamma$,

$$H_{eff.} = -\frac{G_F}{\sqrt{2}4\pi}m_c \bar{u} [c_{\gamma}\gamma L + c_{\gamma}'\gamma L] \sigma_{\mu\nu}F^{\mu\nu}c$$

$$-\frac{G_F}{\sqrt{2}4\pi}m_c \bar{u} [c_{\gamma}\gamma L + c_{\gamma}'\gamma L] \sigma_{\mu\nu}T^\alpha G^{\mu\nu}\alpha c.$$  \hspace{1cm} (81)

Note that here we do not factor out the CKM factor from $c_{\gamma}^{l'},8_g s$ (hence use lower case symbol) in the Hamiltonian. For chargino contributions we have,

$$c_{\gamma}\gamma \sim = \frac{m_\chi^2}{\tilde{m}_\chi^2}$$

$$\times \left\{ U_{\gamma j1}^D U_{\gamma j1}^U V_{ul}^d \tilde{Y}_{\gamma e} e_{\gamma j1}+ U_{\gamma j1}^D U_{\gamma j1}^U V_{ul}^d \tilde{Y}_{\gamma e} e_{\gamma j1} \right\}$$

$$-U_{\gamma j1}^D U_{\gamma j1}^U V_{ul}^d \tilde{Y}_{\gamma e} e_{\gamma j1}+ U_{\gamma j1}^D U_{\gamma j1}^U V_{ul}^d \tilde{Y}_{\gamma e} e_{\gamma j1} \right\}$$

$$\times \left\{ \frac{Q_4g_4(x_{\gamma j1})}{g_3(x_{\gamma j1})} - \frac{m_\gamma}{m_c} \left( U_{\gamma j1}^D U_{\gamma j1}^U V_{ul}^d \tilde{Y}_{\gamma e} e_{\gamma j1} \right) \right\}$$

$$= \frac{G_F}{\sqrt{2}4\pi}m_c \bar{u} [c_{\gamma}\gamma L + c_{\gamma}'\gamma L] \sigma_{\mu\nu}T^\alpha G^{\mu\nu}\alpha c.$$  \hspace{1cm} (81)

For the parameter space considered in the previous case, the chargino loop gives $B_r(c \to u\gamma) \sim 10^{-9}$, dominated by chiral enhanced LL mixing. Therefore, the tan $\beta$ enhancement effects are small and unable to overcome the heavy superparticle decoupling effects. This is also true in other up-type FCNC processes, such as $t \to c\gamma$ to be discussed later. Charged Higgs contribution is small since we do not have the top quark in the loop.

Formulas for gluino and neutralino contributions are similar to previous sections with a trivial modification on neutralino mixing matrices $G_{L,R}$ and $H_{L,R}$. For $m_\tilde{t} = m_\tilde{\tau} = 1.5$ TeV, $|\mu| = 100$–1000 GeV, tan $\beta = 2$–50 and $V_{ul}^d \tilde{Y}_{\gamma e} e_{\gamma j1} \sim \lambda$ the chargino loop gives $|c_{\gamma}\gamma| \sim 10^{-3}$, which is one order of magnitude below the Cabibbo favored two-loop amplitude. Contributions from LR and RL mixings are smaller by a few orders of magnitude compared to the LL mixing contribution. Therefore, for the $t \to c\gamma$ case, the SM result is very small, $B_r(t \to c\gamma) \sim 10^{-13}$ [33]. In this model, by using similar formulas for $c_{\gamma}\gamma$, we have,

$$\Gamma(t \to c\gamma) = \frac{G_F^2\alpha}{32\pi^3}m_\tau^2(|c_{\gamma}\gamma|^2 + |c_{\gamma}'\gamma|^2).$$  \hspace{1cm} (83)

For the parameter space considered in the previous case, the chargino loop gives $B_r(t \to c\gamma) \sim 10^{-9}$, dominated by chiral enhanced LL mixing. Therefore, it is not sensitive to the non-standard soft breaking term. The rate is still unobservable. The gluino contribution is as small as the SM one. This is in contrast to generic MSSM with non-universal soft squark masses where the rate can be close to experimental bounds [84].

We see that, even though the flavor violation is shifted to the up-type sector, we still do not have large FCNC nor CP violation in $t, c$ decays. This is because of heavy $\tilde{m}, m_\tau$ masses as required by meson mixings, and absence of enhancement mechanisms in gluino and neutralino diagrams.
C. Light Strange-Beauty Squark?

In this paper we have focused on the general case of naturally large $d_R$-$b_R$ or $s_R$-$b_R$ mixings as a consequence of Abelian flavor symmetry with SUSY. As we have seen, kaon FCNC constraints require texture zeros to remove $d$-$s$ mixing. This is done by QSA, which shifts the source of Cabibbo angle to up-type sector. It then follows that both the $K^0$ mixing constraint and $D^0$ mixing bound demand TeV scale SUSY particles. In the case of $d_R$-$b_R$ mixing (mutually exclusive with $s_R$-$b_R$ mixing), $B_d$ mixing constraint also implies TeV scale sparticle masses, and one could get maximal sin $2\phi_1$ as suggested by recent Belle result [1].

The $s_R$-$b_R$ mixing case has several special features worthy of note. First, it is maximal, largely because $V_{cb} \sim m_s/m_b \sim \lambda^2$. Second, unlike $\Delta m_{B_d}$ which is precisely measured already, we only have a lower bound on $\Delta m_{B_s}$. In fact, data hints at $\Delta m_{B_s} > \Delta m_{B_s}^{SM}$. From the latter, one cannot draw the conclusion that $B_s$ mixing data demand TeV scale sparticles. From the former, it is intriguing that, in fact, one has a mechanism for the possibility of one light squark. As pointed out in Ref. [19], the “democratic” nature of the 2-3 sub-matrix of $M_{U_R}$ in Eq. (17) not only induces maximal $s_R$-$b_R$ mixing, it could also drive one mass eigenstate $\tilde{s}_1$, dubbed the “strange-beauty” squark because it carries both flavors equally, to be much lighter by level splitting. What is rather surprising is that, having $\tilde{s}_1$ as light as 100 GeV does not make visible impact on the $b \rightarrow s\gamma$ rate. Thus, the light $\tilde{s}_1$ scenario survives one of the strongest known constraints on new physics! This has phenomenological bearings.

The general average squark mass scale $\bar{m}$ is still fixed by $K^0$ and $D^0$ mixings at TeV. But with some tuning in the $M_{U_R}$ matrix, for example $\bar{m}_{23}/\bar{m}_2 \sim 1$ to $\chi^3$ order, $\tilde{s}_1$ can be brought down to 100 GeV. With such large squark mass splittings, the formulas in previous sections do not apply, but one can still follow Ref. [19]. In fact, we find that the Br($b \rightarrow s\gamma$) constraint itself can be easily satisfied even if $m_{\tilde{s}_1} \rightarrow 0$. We note three major consequences of experimental interest: i) Because of low $\tilde{s}_1$ mass, sizable $\chi^2_R$ is generated. Although it is subdominant in $b \rightarrow s\gamma$ rate, it allows the mixing dependent CP asymmetry, e.g. in $B_s \rightarrow \phi\gamma$, to go up to 60%. There is no need to resort to nonstandard C-terms in this case. ii) A light $\tilde{s}_1$ squark further enriches $B_s$ mixing and its associated CP phase. $\Delta m_{B_s}$ could be close to or larger than the SM expectation, and a non-vanishing CP phase in $B_s$ mixing can be measured readily in the moderate $x_{B_s}$ case. iii) A light $\tilde{s}_1$ squark clearly offers itself for direct search at future colliders, in a model where sparticles are otherwise at TeV scale. In fact, one neutralino, the bino, could also be rather light. A possible decay hence search scenario is $\tilde{s}_1 \rightarrow (b,s)+\chi^0_1$ with equal probability of $s$ and $b$ quarks in decay final state. Note that other predictions, such as sizable $x_D$ from SUSY, still hold.

This special scenario, perhaps a bit tuned, seems tailor made for spectacular measurements at the Tevatron and the future LHC. More details and discussions can be found in Ref. [18].

VIII. CONCLUSION

In this work, we make a complete one-loop analysis in SUSY AHS models on FCNC concerning $B_d$, $B_s$, $K^0$, $D^0$ mixings and $b \rightarrow d\gamma, s\gamma$ decays. We find that $B_d$ ($B_s$) and $D^0$ mixings all receive sizable SUSY contributions even with TeV scale superparticles.

Large off-diagonal elements involving the third generation in the fermion mass matrices follow naturally in AHS models. Hence, flavor mixings involving $d_R$ are naturally prominent. It could be the source for near maximal sin $2\phi_1$ given by recent experiments. For $m_{\tilde{g}}$ and $m_{\tilde{q}}$ at TeV scale, the effects could be comparable to SM in $B_s$ (or $B_d$) mixing, leading to $\phi_{B_d} \neq \phi_1$ (or $\phi_{B_s} \neq 0$), while $K^0$ mixing and $\varepsilon_K$ require quark-squark alignment to make $M_{12}^{\tilde{q}}$ and $M_{13}^{\tilde{q}}$ vanish. This shifts $V_{us}$ to $u$ sector, and $\tilde{u}_L-\tilde{c}_L$ mixing with masses $\sim$ TeV gives $D^0$ mixing that is tantalizing close to recent hints from data. With the same squark mixings, the chargino induced contributions to the kaon mixing also points to a TeV scale for superparticle masses, independent of $\Delta m_{B_s}$ considerations. There is a special variant where a “strange-beauty” squark is driven light by maximal $s_R$-$b_R$ mixing, which can give rise to even more astounding phenomena, but with little impact on $B_d$ system. Otherwise, TeV is in general the preferred sparticle scale in this model.

With such heavy gluino and squarks, one has few other low energy phenomena, and prospects for direct production are depressing. It is possible to have $\sim 10\%$ mixing-dependent asymmetry, $a_{M^{\tilde{q}\gamma}}$ in $b \rightarrow s\gamma$ and $d\gamma$ transitions. In addition, these asymmetries are sensitive to nonstandard soft breaking terms via $\tan\beta$ enhancement, and asymmetries could be up to 60% in $b \rightarrow s\gamma$ when the Br($B \rightarrow X_{s\gamma}$) constraint is taken into account. If we insist on non-vanishing $M_{12}^{\tilde{q}}$, the kaon mixing constraint requires, indirectly, that $s$ flavor is almost decoupled from the other down flavors. It is then possible that one has SUSY effects in $B_d$, $D^0$ but not $B_s$ mixings, while $a_{M^{\tilde{q}\gamma}}$ and $a_{\phi^{\gamma\gamma}}$ could be maximal with rate enhancements up to a factor of five. Alternatively, effects could concentrate in $B_s$ system and $b \rightarrow s\gamma$ (plus $D^0$ mixing). The phenomenology outlined here can be tested at B factories and the Tevatron in the next few years even if the New Physics scale is so high such that direct searches show no effect.

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