Supersymmetric contribution to the CP asymmetry of $B \to J/\psi \phi$ in the light of recent $B_s - \bar{B}_s$ measurements

Shaaban Khalil

1Ain Shams University, Faculty of Science, Cairo, 11566, Egypt.
German University in Cairo-GUC, New Cairo city, Egypt.

(Dated: December 26, 2017)

We derive new model independent constraints on the supersymmetric extensions of the standard model from the new experimental measurements of $B_s - \bar{B}_s$ mass difference. We point out that supersymmetry can still give a significant contribution to the CP asymmetry of $B \to J/\psi \phi$ that can be measured at the LHCb experiment. These new constraints on the LL and RR squark mixing severely restrict their possible contributions to the CP asymmetries of $B \to \phi K$ and $B \to \eta' K$.

Therefore, SUSY models with dominant LR flavor mixing is the only way to accommodate the apparent deviation of CP asymmetries from those expected in the standard model. Finally we present an example of SUSY non-minimal flavor model that can accommodate the new $\Delta M_{B_s}$ results and also induces significant CP asymmetries in $B_s \to J/\psi \phi$, $B \to \phi K$ and $B \to \eta' K$ processes.

1— Recently, the $D \emptyset$ [1] and CDF [2] collaborations have reported new results for the $B_s - \bar{B}_s$ mass difference:

$$17 \text{ ps}^{-1} < \Delta M_{B_s} < 21 \text{ ps}^{-1} \quad 90\% \text{ C.L.} \quad (D \emptyset),$$

$$\Delta M_{B_s} = 17.33^{+0.42}_{-0.23} \pm 0.07 \text{ ps}^{-1} \quad (CDF), \quad (1)$$

which seems consistent with the Standard Model (SM) predictions. In fact, the estimation of the SM value for $\Delta M_{B_s}$ contains large hadronic uncertainties. The $B_s^0 - B_s^0$ mass difference is defined as $\Delta M_{B_s} = 2M_{12}(B_s) = 2\langle |B_s^0| H_{\text{eff}}^{B=2}|B_s^0 \rangle $, where $H_{\text{eff}}^{B=2}$ is the effective Hamiltonian responsible for the $\Delta B = 2$ transition. In the SM, $H_{\text{eff}}^{B=2}$ is generated by the box diagrams with $W$-exchange. The best determination for $\Delta M_{B_s}^\text{SM}$ can be obtained from a ratio to the $\Delta M_{B_d}^\text{SM}$ in which some QCD corrections as well as $t$ quark mass dependence are cancelled out

$$\frac{\Delta M_{B_s}^\text{SM}}{\Delta M_{B_d}^\text{SM}} = \frac{M_{B_s} B_s f_{B_s}^2}{M_{B_d} B_d f_{B_d}^2} \frac{|V_{ts}|^2}{|V_{td}|^2}, \quad (2)$$

where $M_{B_d} = 5.28$ GeV and $M_{B_s} = 5.37$ GeV and the lattice calculations lead to $B_d f_{B_d}^2/(B_s f_{B_s}^2) = (1.15 \pm 0.06^{+0.07}_{-0.07})^2$. Since the $B_d^0 - B_d^0$ oscillation is mostly saturated by the SM contributions [4], we can assume that $\Delta M_{B_d}^\text{SM} = \Delta M_{B_d}^\text{exp} = (0.502 \pm 0.007)\text{ ps}^{-1}$. Finally, $|V_{ts}|^2/|V_{td}|^2$ can be given as a function of the angle $\gamma$ of the unitary triangle of Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix. In Fig. 1, we present the allowed range of $\Delta M_{B_d}^\text{SM}$ in terms of the angle $\gamma$ (measured from a pure SM process). Here we assume that $|V_{ts}|$ and $|V_{td}|$ are free of new physics and can be determined by the SM contribution to the semileptonic decay. Also, it is assumed that the angle $\beta$ is given by $\beta^\text{SM}$, measured from $B_d \to J/\psi K_s$. As can be seen from this figure, the new bounds on $\Delta M_B$ impose stringent constraints on the values of $\gamma^\text{SM}$. The lower bound of $D \emptyset$ result excludes values of $\gamma^\text{SM} > 70^\circ$. It is worth mentioning that the best fit for $\gamma^\text{SM}$ and $\Delta M_{B_s}^\text{SM}$, according to UTfit group is given by [5]:

$$\gamma^\text{SM} = 61.3 \pm 4.5, \Delta M_{B_s}^\text{SM} = (17.45 \pm 0.25)\text{ ps}^{-1}, \quad (3)$$

and according to CKMfitter group is given by [6]

$$\gamma^\text{SM} = 59.8^{+4.9}_{-4.1}, \Delta M_{B_s}^\text{SM} = 17.3^{+0.49}_{-0.29}. \quad (4)$$

Therefore, it is expected that the experimental measurements in Eq. (1) provide important constraints on any new physics beyond the SM [7]. In this letter, we study the constraints imposed on supersymmetric model due to these experimental limits. We derive model independent bounds on the relevant SUSY mass insertions. Then we analyze the implications of these constraints on the supersymmetric contribution to the CP asymmetry in $B_s \to J/\psi \phi$ process. Finally, we consider the SUSY non-minimal flavor model studied in Ref. [8], as an example for SUSY model, that can accommodate the new
$\Delta M_{B_s}$ results and also induces significant CP asymmetry in $B \to J/\psi \phi$ which can be measured at the LHCb experiment.

2- In supersymmetric theories, the effective Hamiltonian $H_{\Delta B=2}^\text{eff}$ receives new contributions through the box diagrams mediated by gluino, chargino, neutralino, and charged Higgs. It turns out that gluino exchanges give the dominant contributions \cite{Chung}. The most general effective Hamiltonian for $\Delta B = 2$ processes, induced by gluino exchange through $\Delta B = 2$ box diagrams, can be expressed as

$$H_{\text{eff}}^{\Delta B = 2} = \sum_{i=1}^{5} C_i(\mu) Q_i(\mu) + \sum_{i=1}^{3} \tilde{C}_i(\mu) \tilde{Q}_i(\mu) + h.c. , \quad (5)$$

where $C_i(\mu)$, $\tilde{C}_i(\mu)$ and $Q_i(\mu)$, $\tilde{Q}_i(\mu)$ are the Wilson coefficients and the local operators normalized at the scale $m_b$, respectively, which can be found in Ref.\cite{Chung}. As in the $B_d$ system, the effect of SUSY can be parameterized by a dimensionless parameter $r_s$ and a phase $\theta_s$ defined as follows:

$$r_s e^{i\theta_s} = \sqrt{\frac{\mathcal{M}_{12}(B_s)}{\mathcal{M}_{12}^\text{SM}(B_s)}}, \quad (6)$$

where $\mathcal{M}_{12}(B_s) = \langle B_s^0 | H_{\text{eff}}^{\Delta B = 2} | B_s^0 \rangle \equiv \mathcal{M}_{12}^{\text{SM}} + \mathcal{M}_{12}^{\text{SUSY}}$. Thus, the total $B_s - \bar{B}_s$ mass difference is given by $\Delta M_{B_s} = 2\mathcal{M}_{12}^2(B_s)$, and the down squark mass insertions from second and third generations, $(\delta_{23}^d)_{LL}$, where $A$ and $B$ stand for left ($L$) or right ($R$) handed mixing. A general expression for $R_s = \mathcal{M}_{12}^2/\mathcal{M}_{12}^{\text{SM}}$ has been given in Ref.\cite{Chung} as follows:

$$R_s = a_1(m_q \bar{q}, x)\left[ (\delta_{LL}^d)^2 + (\delta_{RR}^d)^2 + \frac{3}{4} a_2(m_q \bar{q}, x) [(\delta_{LL}^d)^2 + (\delta_{RR}^d)^2] \right]$$

$$+ a_4(m_q \bar{q}, x) \left( \delta_{RR}^d (\delta_{LL}^d)_{23} \right) + a_4(m_q \bar{q}, x) \left( \delta_{LL}^d (\delta_{RR}^d)_{23} \right), \quad (7)$$

where the coefficients $|a_1| \simeq O(1)$, $|a_2| < |a_3| < |a_4| \simeq O(100)$. For instance, with $m_q = 300$ and $x = 1$, one finds

$$R_s = 7.2 [(\delta_{LL}^d)^2 + (\delta_{RR}^d)^2] + 129.8 [(\delta_{LL}^d)^2 + (\delta_{RR}^d)^2]$$

$$- 205.7 [(\delta_{LL}^d)_{23}(\delta_{RR}^d)_{23}] - 803.8 [(\delta_{LL}^d)_{23}(\delta_{RR}^d)_{23}]. \quad (8)$$

Note that $r^2_s = 1 + R_s |^2$. From the experimental upper bound on $\Delta M_{B_s}$ in Eq.\cite{Chung}, one can derive an upper bound on the mass insertions involved in Eq.\cite{Chung}. In order to find conservative upper bounds, we set the SM contribution to its best fit value, namely $\Delta M_{B_s}^{\text{SM}} = 17.5$ ps$^{-1}$. In this case, the $|R_s|$ should satisfy the following bound:

$$|R_s| = \left| \frac{\Delta M_{B_s}}{\Delta M_{B_s}^{\text{SM}}} \right| - 1 \lesssim 4/17. \quad (9)$$

It is worth mentioning that if one assumes that $\Delta M_{B_s}^{\text{SM}} \simeq 21$ ps$^{-1}$, the above bound remains valid. In table 1 we present our results for the upper bounds on $(\delta_{LL}^d)_{23}$ mass insertions from their individual contributions to $B_s - \bar{B}_s$ mixing for $m_q = 300$ GeV and $x$ varies from 0.25 to 2. As can be seen from Eq.\cite{Chung} that the constraints imposed on the mass insertions are symmetric under changing $L \leftrightarrow R$. Therefore, we present in table 1 the upper bounds on one combination of the mass insertions.

Table 1: Upper bounds on $(\delta_{LL}^d)_{23}$, $(\delta_{RR}^d)_{23}$ from $\Delta M_{B_s} < 21$ ps$^{-1}$ for $m_q = 300$ GeV.

| $x$ | $(\delta_{LL}^d)_{23}$ | $(\delta_{RR}^d)_{23}$ | $\sqrt{(\delta_{LL}^d)_{23}^2 + (\delta_{RR}^d)_{23}^2}$ | $\sqrt{(\delta_{LL}^d)_{23}^2 + (\delta_{RR}^d)_{23}^2}$ |
|-----|-----------------|-----------------|-----------------|-----------------|
| 0.25 | 0.074           | 0.035           | 0.018           | 0.014           |
| 0.5  | 0.11            | 0.037           | 0.024           | 0.015           |
| 1    | 0.17            | 0.04            | 0.032           | 0.016           |
| 1.5  | 0.27            | 0.43            | 0.039           | 0.017           |
| 2    | 0.46            | 0.046           | 0.046           | 0.018           |

Three comments on the results of table 1 are in order: 1) the constraints obtained on $(\delta_{LL}^d)_{23}$ are the strongest known constraints on these mass insertions, since other processes based on $b \to s$ transition, like $B \to X_s \gamma$, leave them unconstrained \cite{Chung}. In fact with these constraints, one can verify that the $LL(RR)$ contributions to $B \to \phi K$, $B \to \eta' K$ and $B \to \pi K$ are diminished and become insignificant. Therefore, $LR$ contribution remains as the only candidate for saturating any deviation from the SM results in the CP asymmetries or branching ratios of these processes \cite{Chung}. 2) The upper bounds on $LR(RL)$ mass insertions from the $B_s - \bar{B}_s$ are less stringent than those derived from the experimental limits of the branching ratio of $B \to X_s \gamma$ \cite{Chung}. Therefore, from the LHCb experiment, the final state of $B_s \to J/\psi \phi$ is not a CP eigenstate, but a superposition of CP odd and even states which can, however, be dis-
entangled through an angular analysis of their products. This angular distribution yields to a tiny direct CP violation. Thus, the CP asymmetry of the $B_s$ and $\bar{B}_s$ meson decay to $J/\psi \phi$ is given by

$$a_{J/\psi \phi}(t) = \frac{\Gamma(\bar{B}_s^0(t) \to J/\psi \phi) - \Gamma(B_s^0(t) \to J/\psi \phi)}{\Gamma(\bar{B}_s^0(t) \to J/\psi \phi) + \Gamma(B_s^0(t) \to J/\psi \phi)} = S_{J/\psi \phi} \sin(\Delta M_{B_s} t),$$

(10)

where $S_{J/\psi \phi}$ is the mixing-induced CP asymmetry. In the SM, the mixing CP asymmetry $S_{J/\psi \phi}$ is given by

$$S_{J/\psi \phi} = \sin 2\beta_{\text{SM}} \sin [2 \arg(V_{tb}V_{ts}^*)] \simeq -2\lambda^2 \eta \simeq O(10^{-2}).$$

(11)

Such small CP asymmetry in the SM gives the hope that if a sizable value of $S_{J/\psi \phi}$ is found in future experiments (in particular at LHCb experiment), then it would be an immediate signal for a new physics effect.

In the presence of SUSY contribution, the CP asymmetry $S_{J/\psi \phi}$ is given by

$$S_{J/\psi \phi} = \sin 2\beta_{\text{SUSY}} = \sin (2\beta_s + 2\theta_s),$$

(12)

where $\theta_s$ is given in Eq. (6) as $2\theta_s = \arg(1 + R_\eta)$. Therefore, the value of $S_{J/\psi \phi}$ depends on the magnitude of $R_\eta$ which, as emphasized above, is constrained from $\Delta M_{B_s}$ to be less than or equal to 4/17. In this respect, it is easy to show that the maximum value of $S_{J/\psi \phi}$ that one may obtain from SUSY contributions to the $B_s - \bar{B}_s$ mixing is given by

$$S_{J/\psi \phi} \simeq 0.24.$$  

(13)

It is important to note that due to the stringent constraints on $(\delta_{LR}^d)_{23}$ from $b \to s\gamma$: $|\delta_{LR}^d| \leq 0.016$, the LR (RL) supersymmetric contribution to $S_{J/\psi \phi}$ is very restricted. It implies that $S_{J/\psi \phi} < 0.02$, which is too small to be observed at the Tevatron or the LHC. Therefore, the LR and RL contributions can not provide significant contribution to $B_s$ mixing or to the mixing CP asymmetry of $B_s \to J/\psi \phi$.

On the other hand, the LL and RR mass insertions can generate sizable and measurable values of $S_{J/\psi \phi}$. For instance, $(\delta_{LL}^d)_{23} \simeq 0.17 e^{i\pi/4}$ yields to $R \simeq 0.24 e^{i\pi/2}$ which implies that $\sin 2\beta_s \simeq 0.24$. However, as mentioned above, it is important to note that since the minimum value of the mass insertion $(\delta_{LL}^d)_{23}$ is of order $10^{-2}$, too small to be observed at the Tevatron, its contributions are considered simultaneously in determining the $\Delta M_{B_s}$ and $\sin 2\beta_s$.

4—We now consider the impact of the $\Delta M_{B_s}$ constraints derived above on the mixed CP asymmetries in $B_d \to \phi K$ and $B_d \to \eta'K$ processes, which at the quark level are also based on $b \to s$ transition. The BaBar and Belle results for these asymmetries lead to the following averages:

$$S_{\phi K} = 0.47 \pm 0.19, \quad S_{\eta'K} = 0.48 \pm 0.09,$$  

(14)

which display about 1σ and 2.5σ deviation from the SM predictions, respectively.

The SUSY contributions to the decay amplitudes of $B_d \to \phi K$ and $B_d \to \eta'K$ are given by

$$A_{\phi K} = -i \sqrt{2} m_B^2 F_+^{B_d \to K} f_{\phi} \sum_{i=1}^{12} H_i(\phi)(C_i + \bar{C}_i),$$

$$A_{\eta'K} = -i \sqrt{2} m_B^2 F_+^{B_d \to K} f_{\eta'} \sum_{i=1}^{12} H_i(\eta')(C_i - \bar{C}_i),$$  

(15)

where the $C_i$ are the corresponding Wilson coefficients to the local operators of $b \to s$ transition. $C_i$ as functions of the mass insertions $(\delta_{LL}^d)_{23}$ and $(\delta_{LR}^d)_{23}$ can be found in Ref. [12]. Here the QCD factorization mechanism is adopted to determine the hadronic matrix elements and as in Ref. [12] they can be parameterized in terms of the parameters $H_i(\phi)$ and $H_i(\eta')$ which are given in Ref. [12]. In terms of SUSY contributions, the CP asymmetry $S_{\phi(\eta')K}$ can be written as

$$S_{\phi(\eta')K} = \sin 2\beta + 2|R_{\phi(\eta')}| \cos \delta_{\phi(\eta')} \sin \theta_{\phi(\eta')} \sin 2\beta,$$  

(16)

where

$$R_{\phi(\eta')} = \left( \frac{A_{\text{SUSY}}}{A_{\text{SM}}} \right)_{\phi(\eta')K}, \quad \theta_{\phi(\eta')} = \arg \left( \frac{A_{\text{SUSY}}}{A_{\text{SM}}} \right)_{\phi(\eta')K}$$

and $\delta_{\phi(\eta')}$ is the strong phase. Thus, in order to derive $S_{\phi(\eta')K}$ toward their central values of the average experimental results in Eq. (14), $|R_{\phi(\eta')}| \geq 0.2$ should be satisfied. For a gluino mass and average squark mass of order $\tilde{m} = m_{\tilde{g}} = 500$ GeV, one finds

$$R_{\phi} = -0.14 e^{-i0.1}(\delta_{LL}^d)_{23} - 127 e^{-i0.08}(\delta_{LR}^d)_{23} + L \leftrightarrow R,$$  

(17)

and

$$R_{\eta'} = -0.07 e^{i0.24}(\delta_{LL}^d)_{23} - 64(\delta_{LR}^d)_{23} - L \leftrightarrow R.$$  

(18)

It is now clear that the $\Delta M_{B_s}$ constraints play a crucial role in reducing the LL and RR contributions to the $S_{\phi(\eta')K}$. By implementing the bounds in table 1, one can easily observe that LL( RR) contribution leads to $|R_{\phi(\eta')}| \sim O(10^{-2})$ which yields a negligible effect on $S_{\phi(\eta')}$ and one can safely conclude that the LL and RR mass insertions can not provide an explanation to any
deviation in $S_{\bar{\phi}(\bar{\psi}')} \phi(\bar{\psi}')$ results. On the other hand, the contribution of $|\delta_{LR}^{\phi}|^2 |^2$ is less constrained by $\Delta M_{B_s}$ and large effects in $|R_{\bar{\phi}(\bar{\psi}')}|$ that could drive $S_{\bar{\phi}(\bar{\psi}')} R$ toward 0.4 can be achieved.

5- The above results show that $S_{J/\psi \phi}$ and $S_{\bar{\phi}(\bar{\psi}')} R$ are dominated by different mass insertions: LL and LR/RL respectively. As emphasized in Ref. [8], these two mass insertions can be enhanced simultaneously in SUSY models with intermediate/large $\tan \beta$ and a simple non-minimal flavor structure, where the scalar mass of the first two generations is different from the scalar mass of the third generation. In particular, let us consider the following soft SUSY breaking terms are assumed at the GUT scale.

$$
M_1 = M_2 = M_3 = M_{1/2}, \quad A^u = A^d = A_0 e^{i \phi_A},
$$

$$
M^2_U = M^2_0 = m_0^2, \quad m^2_{H_1} = m^2_{H_2} = m_0^2,
$$

$$
M^2_Q = \begin{pmatrix} m_0^2 & m_0^2 \\ m_0^2 & a^2 m_0^2 \end{pmatrix}. \quad (19)
$$

The parameter $a$ measures the non-universality of the squark masses. It is worth mentioning that the EDM constraints on the CP violating phase $\phi_A$ of the trilinear coupling is less severe than the constraints imposed on the other SUSY CP phases and can be of order $O(0.1)$. [8]

Using the relevant renormalization group equations, one can explore these parameters from GUT scale to the electroweak scale, where we impose the electroweak symmetry breaking conditions and calculate the squark mass matrices. Then we determine the numerical values of the corresponding mass insertions. For instance, for $a = 5$, $\tan \beta = 15$ and $m_3 \sim m_q \sim 500$ GeV, one finds that $|\delta_{LR}^{\phi}|^2 \simeq 0.18$ which leads to $\Delta M_{B_s} \simeq 19$ ps$^{-1}$. Also with a proper choice for the phase $\phi_A$, one can get $\arg(\delta_{LR}^{\phi}) \simeq 0.7$ which implies that $S_{J/\psi \phi} \simeq 0.1$ which can be measured by the LHCb experiment. Note that in this scenario the phases of the mass insertions are due to a combined effect of the SM phase in the CKM mixing matrix and the SUSY CP phase $\phi_A$. However, for the LL mass insertion the main effect is due the CKM phase, see Ref. [8] for more details.

Concerning the mass insertion $\delta_{LR}^{\phi}$, it is expected to be negligible due to the universality of the trilinear couplings. However, with intermediate/large $\tan \beta$, the double mass insertion is quite important and it gives the dominant effect as follows [8].

$$
(\delta_{LR}^{\phi})_{23,23} = (\delta_{LR}^{\phi})_{23} + (\delta_{LR}^{\phi})_{23} (\delta_{LR}^{d})_{23}, \quad (20)
$$

where $(\delta_{LR}^{d})_{23} \simeq \frac{m_A}{m} \frac{e^{i \phi_A}}{m \tan \beta}$. Since $(\delta_{LR}^{d})_{23}$ is negligible, $(\delta_{LR}^{\phi})_{23}$ is given by

$$
(\delta_{LR}^{\phi})_{23} \simeq (\delta_{LR}^{\phi})_{23} \frac{m_0}{m} \tan \beta. \quad (21)
$$

The parameter $\mu$ is determined by the electroweak conditions and it is found to be of order the squark mass. The phase of $\mu$ set to zero to overcome the EDM constraints. Since $(\delta_{LR}^{\phi})_{23} \simeq 0.18$, the value of $(\delta_{LR}^{\phi})_{23}$ is of order $10^{-2}$ which is sufficient to reduce the CP asymmetries $S_{B_s K}$ and $S_{B_s K}$ from the SM result $\sin 2\beta \simeq 0.7$ to their central values of average experimental results.

6- To conclude, we have considered the supersymmetric contributions to the $B_s - \bar{B}_s$ mixing. We derived new model independent constraints on the magnitude of the mass insertions $(\delta_{AB})_{23}$, where $\{A, B\} = \{L, R\}$, from the new experimental measurements of $\Delta M_{B_s}$. We showed that by implementing these constraint, the SUSY contribution, through the LL mixing, can enhance the CP asymmetry of $B_s \to J/\psi \phi$ up to 0.24, which can be observed at the LHCb experiment. We also emphasized that the new constraints exclude the SUSY models with large RR flavor mixing and severely restrict the LL contributions to the CP asymmetries of $B \to \phi K$ and $B \to \eta' K$. Therefore, SUSY models with dominant LR flavor mixing is the only way to accommodate the apparent deviation of CP asymmetries from those expected in the standard model. Finally we studied an example of SUSY non-minimal flavor model and intermediate/large $\tan \beta$. We showed that in this model the new $\Delta M_{B_s}$ results and also the CP asymmetries in $B_s \to J/\psi \phi$, $B \to \phi K$ and $B \to \eta' K$ processes can be simultaneously saturated.

[1] V. Abazov [D0 Collaboration], arXiv:hep-ex/0603029
[2] G. Gomez-Ceballos [CDF Collaboration], Talk at FPCP 2006, http://fpcp2006.triumf.ca/agenda.php
[3] N. Yamada et al. [JLQCD Collaboration], Nucl. Phys. Proc. Suppl. 106 (2002) 397; D. Becirevic, V. Gimenez, G. Martinelli, M. Papinutto and J. Reyes, JHEP 0204 (2002) 025; S. M. Ryan, Nucl. Phys. Proc. Suppl. 106 (2002) 86.
[4] E. Gabrielli and S. Khalil, Phys. Rev. D 67, 015008 (2003).
[5] UTfit Collaborations, http://utfit.roma1.infn.it
[6] CKMfit Collaborations, http://ckmfitter.in2p3.fr
[7] P. Ball and R. Fleischer, arXiv:hep-ph/0604249
F. J. Foster, K. i. Okumura and L. Roszkowski, arXiv:hep-ph/0604121
Z. Ligeti, M. Papucci and G. Perez, arXiv:hep-ph/0604112
M. Blanke, A. J. Buras, D. Guadagnoli and C. Tarantino, arXiv:hep-ph/0604057;
M. Endo and S. Mishima, arXiv:hep-ph/0603251
M. Ciuccioni and L. Silvestrini, arXiv:hep-ph/0603114.
[8] S. Khalil, Phys. Rev. D 72, 055020 (2005);
[9] P. Ball, S. Khalil and E. Kou, Phys. Rev. D 69, 115011 (2004).
[10] S. Khalil, Phys. Rev. D 72, 035007 (2005).
[11] S. Abel and S. Khalil, Phys. Lett. B 618, 201 (2005);
J. Hisano and Y. Shimizu, Phys. Rev. D 70, 093001 (2004); M. Endo, M. Kakizaki and M. Yamaguchi, Phys. Lett. B 583, 186 (2004).
[12] A. S. Dighe, I. Dunietz, H. J. Lipkin and J. L. Rosner,
Phys. Lett. B 369, 144 (1996).

[13] E. Gabrielli, K. Huitu and S. Khalil, Nucl. Phys. B 710, 139 (2005).

[14] S. Khalil and E. Kou, Phys. Rev. D 67, 055009 (2003);

Phys. Rev. Lett. 91, 241602 (2003).

[15] S. Abel, S. Khalil and O. Lebedev, Nucl. Phys. B 606, 151 (2001).