Z’-mediated Supersymmetry Breaking

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We consider a class of models in which supersymmetry breaking is communicated dominantly via a $U(1)^\prime$ gauge interaction, which also helps solve the $\mu$ problem. Such models can emerge naturally in top-down constructions and are a version of split supersymmetry. The spectrum contains heavy sfermions, Higgsinos, exotics, and $Z^\prime \sim 10 - 100 \text{ TeV}$; light gauginos $\sim 100 - 1000 \text{ GeV}$; a light Higgs $\sim 140 \text{ GeV}$; and a light singlino. A specific set of $U(1)^\prime$ charges and exotics is analyzed, and we present five benchmark models. Implications for the gluino lifetime, cold dark matter, and the gravitino and neutrino masses are discussed.

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I. INTRODUCTION AND MOTIVATION

To a large extent, the mediation mechanism of supersymmetry (SUSY) breaking determines the low energy phenomenology. A well-studied scenario is gravity mediation [1]. During the last couple of decades, in order to satisfy the increasingly stringent constraints from flavor changing neutral current measurements, many other mediation mechanisms, such as anomaly mediation [2], gauge mediation [3], and gaugino mediation [4], have been proposed (for a review, see [5]). In this letter, we present an alternative mechanism in which SUSY breaking is mediated by exotic gauge interactions, such as an additional $U(1)^\prime$. Concrete superstring constructions frequently lead to additional, non-anomalous, $U(1)^\prime$ factors in the low-energy theory (see, e.g., [6]) with properties allowing a $U(1)^\prime$-mediated SUSY breaking. Scenarios with an extra $U(1)^\prime$ involved in supersymmetry breaking mediation have been studied in various contexts [7]. Here, we study a new scenario where $Z^\prime$-meditation is the dominant source for both scalar and gaugino masses.

Another ingredient we would like to consider is the $\mu$-problem of the Minimal Supersymmetric Standard Model (MSSM). One class of solutions invokes a spontaneously broken Peccei-Quinn symmetry (see, e.g., [8]). From the point of view of top-down constructions it is common that such a symmetry is promoted to a $U(1)^\prime$ gauge symmetry [9]. Identifying this $U(1)^\prime$ with the mediator of SUSY breaking sets $\mu$ (as well as $\mu B$) to the scale of the other soft SUSY breaking parameters, which are of the right size whether or not the electroweak symmetry breaking is finely tuned.

In the setup we propose, schematically shown in Fig. 1, visible and hidden sector fields do not have direct renormalizable coupling with each other. At the same time, they are both charged under $U(1)^\prime$. A supersymmetry breaking $Z^\prime$-ino mass term, $M_{\tilde{Z}^\prime}^2$, is generated due to the $U(1)^\prime$ coupling to the hidden sector. The observable sector fields feel the supersymmetry breaking through their couplings to $U(1)^\prime$. The sfermion masses are of order $m_f^2 \sim M_{\tilde{Z}^\prime}^2/16\pi^2$. The $SU(3)_C \times SU(2)_L \times U(1)_Y$ gaugino masses are generated at higher loop order, $M_{1,2,3} \sim M_{Z^\prime}/(16\pi^2)^2$, which is 2-3 order of magnitudes lighter than the sfermions. LEP direct searches suggest electroweak-ino masses $> 100 \text{ GeV}$. We therefore expect that the sfermions are heavy, typically about $100 \text{ TeV}$. In this sense, this scenario can be viewed as a mini-version of split-supersymmetry [10]. In particular, one fine-tuning is needed to maintain a low electroweak scale. This scenario does not have flavor or CP violation problems due to the decoupling of the sfermions. One important difference from split-supersymmetry is the $\mu$-parameter, which is set by the scale of $U(1)^\prime$ breaking.

II. GENERIC FEATURES OF Z’-MEDIATED SUPERSYMMETRY BREAKING

The visible sector contains an extension of the MSSM. First, we introduce an extra $U(1)^\prime$ gauge symmetry. Second, the $\mu$ parameter is promoted into a dynamical field, $\mu H_u H_d \rightarrow S H_u H_d$. $S$ is a Standard Model singlet which is charged under the $U(1)^\prime$. Third, we include exotic matter multiplets with Yukawa couplings to $S$, $\sum_{i \in \{\text{exotics}\}} Y_{iSX_iX'_i}$. They are included to cancel the anomalies associated with the $U(1)^\prime$. Such exotics and couplings generically exist in string theory constructions.

FIG. 1: $Z^\prime$-mediated supersymmetry breaking.
A. Features of the Spectrum

We parameterize the hidden sector supersymmetry breaking by a spurion field \( X = M + \theta^2 F \). At the scale \( \Lambda_S \), supersymmetry breaking is assumed to generate a mass \( M_{Z'} \sim g_{Z'}^2(F/M)/16\pi^2 \) for the fermionic component of the \( Z' \) vector superfield.

We assume that all the chiral superfields in the visible sector are charged under \( U(1)' \), so all the corresponding scalars receive soft mass terms at 1-loop,

\[
m_{f_i}^2 \sim \frac{g_{Z'}^2 Q_{f_i}^2}{16\pi^2} M_{Z'}^2 \log \left( \frac{\Lambda_S}{M_{Z'}} \right) \sim (100 \text{ TeV})^2 \tag{1}\]

where \( g_{Z'} \) is the \( U(1)' \) gauge coupling and \( Q_{f_i} \) is the \( U(1)' \) charge of \( f_i \), which we take to be of order unity.

The \( SU(3)_c \times SU(2)_L \times U(1)_Y \) gaugino masses can only be generated at 2-loop level since they do not directly couple to the \( U(1)' \) gaugino,

\[
M_a \sim \frac{g_2^2 g_\alpha^2}{(16\pi^2)^2} M_{Z'} \log \left( \frac{\Lambda_S}{M_{Z'}} \right) \sim 10^2 - 10^3 \text{ GeV} \tag{2}
\]

where \( g_\alpha \) is the gauge coupling for the gaugino \( \lambda_\alpha \). It is straightforward to verify that this is indeed the leading \( U(1)' \) contribution to the gaugino mass. In particular, kinetic mixing induced by loops of visible sector fields does not contribute significantly due to chiral symmetries.

The gravitino mass \( m_{3/2} \sim F/M_P \) depends strongly on the scale of supersymmetry breaking. Requiring MSSM gaugino masses \( \gtrsim 100 \text{ GeV} \) and assuming \( \sqrt{F}, M \) and \( \Lambda_S \) to be of the same order of magnitude, we find \( \sqrt{F} \sim 10^7 - 10^{11} \text{ GeV} \). This is very different from gauge mediated supersymmetry breaking, where the lower scale \( \sim 10 - 1000 \text{ TeV} \) typically implies a gravitino much lighter than the other superpartners. Here, the scale is constrained logarithmically by the requirement of radiative symmetry breaking. Therefore, the gravitino mass is exponentially sensitive to the choice of model parameters.

We also expect contributions to gaugino masses through gravity mediation of the order \( F/M_P \), which could be of the same order as Eq. (2). However, its contribution to scalar masses \( \sim F^2/M_P^2 \) is negligible compared with the \( Z' \)-mediation. Therefore, we expect the hierarchy between scalar and gaugino masses to be generic.

B. Symmetry breaking and fine-tuning

The \( U(1)' \) gauge symmetry must be broken by the singlet’s VEV \( \langle S \rangle \). We assume this is triggered by radiative corrections to the soft mass \( m_{2}^2 \), especially through Yukawa couplings to exotics. Therefore, successful radiative breaking of \( U(1)' \) usually requires that those couplings are not small. \( \langle S \rangle \) is parametrically only an order of magnitude smaller than \( M_{Z'} \). It is therefore reasonable to first determine \( \langle S \rangle \) ignoring the Higgs doublets, and then to consider the Higgs potential for the doublets regarding \( \langle S \rangle \) as fixed.

To generate the electroweak scale \( \Lambda_{EW} \) we must fine-tune one linear combination of the two Higgs doublets to be much lighter than its natural scale. The full mass matrix for the two Higgs doublets is,

\[
\mathcal{M}_H^2 = \begin{pmatrix} m_H^2 & -A_H \langle S \rangle \\ -A_H \langle S \rangle & m_{1/2}^2 \end{pmatrix}
\]

\[
m_H^2 = m_{H_u}^2 + g_{Z'}^2 Q_S Q_S \langle S \rangle^2 + \lambda_S^2 \langle S \rangle^2
\]

\[
m_{1/2}^2 = m_{H_d}^2 + g_{Z'}^2 Q_S Q_1 \langle S \rangle^2 + \lambda_S^2 \langle S \rangle^2. \tag{3}
\]

Generically, one can tune various elements in \( \mathcal{M}_H^2 \) to obtain one small eigenvalue \( \sim \Lambda_{EW}^2 \). The up-type Higgs mass term can be driven small or negative due to the large top Yukawa coupling. One typically finds solutions by tuning \( |m_{1/2}^2| \ll m_H^2 \sim g_{Z'}^2 M_{Z'}/16\pi^2 \). The trilinear term is smaller, \( A_H \sim \lambda_S^2, M_{Z'}/16\pi^2 \sim \lambda \times 10 \text{ TeV} \), so integrating out \( H_d \) will not shift the smaller eigenvalue significantly. \( \tan \beta \) is well approximated by \( \tan \beta = m_{1/2}^2/A_{H} \langle S \rangle \sim 10 - 100 \). There is a single Standard Model-like Higgs scalar, with mass in the range 140 GeV. The remaining Higgs particles are at a scale of order \( \sim 100 \text{ TeV} \). The Higgs mass is somewhat heavier than the typical prediction of the MSSM, due to the \( U(1)' \) \( D \) term and the running of the effective quartic coupling from \( M_{Z'} \), down to the electroweak scale.

It is possible to tune with all the parameters, such as \( g_{Z'} \) and \( \lambda \), of the same order. In addition, there is an interesting limit when \( g_{Z'} \ll \lambda \). Generically, we expect \( \langle S \rangle \sim M_{Z'}/4\pi \). The singlino mass is \( \sim g_{Z'}^2 Q_S^2 \langle S \rangle^2/M_{Z'} \sim g_{Z'}^2 M_{Z'}/16\pi^2 \ll M_{Z'} \). Moreover, since \( m_{1/2}^2 \propto g_{Z'}^2 M_{Z'}/16\pi^2 \), to fine-tune \( m_{1/2}^2 \sim \Lambda_{EW}^2 \) we expect the parameters to be chosen so that the singlet’s VEV is even smaller \( \langle S \rangle \sim (g_{Z'}/\lambda) M_{Z'}/4\pi \). Therefore, it is possible to have the singlino be very light \( m_3 \sim (10^{-3} - 10^{-5}) M_{Z'} \). In certain cases, the \( Z' \) gauge boson, \( M_{Z'} \sim g_{Z'} Q_S \langle S \rangle \), could even be light enough to be produced at the LHC.

III. MODEL PARAMETERS AND LOW-ENERGY SPECTRUM

The free parameters are \( g_{Z'}, \lambda \), the exotic Yukawa couplings, the \( U(1)' \) charges, \( M_{Z'} \), and the supersymmetry breaking scale \( \Lambda_S \). The charges are chosen to cancel all the anomalies. A minimal choice, which also leads to a light wino (\( M_2 < M_1 \)), involves the introduction of 3 families of colored exotics (D) and two uncolored \( SU(2) \)-singlet families (E). Normalizing the down-type Higgs charge to unity, \( Q_1 = 1 \), we are left with two independent parameters, which we choose to be the up-type Higgs and the left-handed quark charges, \( Q_2 \) and \( Q_3 \) respectively. Several additional constraints need to be satisfied by the


| $Q_2$  | 1   | 2   | 3   | 4   | 5   |
|-------|-----|-----|-----|-----|-----|
| $Q_Q$ | $-\frac{1}{3}$ | $-\frac{1}{3}$ | $-\frac{1}{3}$ | $-\frac{1}{3}$ | $-\frac{1}{3}$ |
| $g_{\tilde{u}}$ | 0.45 | 0.23 | 0.23 | 0.23 | 0.06 |
| $\lambda$ | 0.5 | 0.8 | 0.8 | 0.8 | 0.3 |
| $Y_D$ | 0.6 | 0.7 | 0.8 | 0.8 | 0.4 |
| $Y_E$ | 0.6 | 0.6 | 0.6 | 0.6 | 0.1 |

| $(\delta)$ | $2 \times 10^6$ | $7 \times 10^6$ | $6 \times 10^4$ | $2 \times 10^7$ | $8 \times 10^4$ |
| $\tan \beta$ | 20 | 29 | 33 | 45 | 60 |
| $M_1$ | 2700 | 735 | 650 | 760 | 270 |
| $M_2$ | 710 | 195 | 180 | 340 | 123 |
| $M_3$ | 4300 | 1200 | 1100 | 540 | 200 |
| $m_{\tilde{H}}$ | 140 | 140 | 140 | 140 | 140 |
| $m_{\tilde{Q}_2}$ | $10^3$ | $5 \times 10^4$ | $4 \times 10^4$ | $8 \times 10^4$ | $4 \times 10^4$ |
| $m_{\tilde{L}_3}$ | $3 \times 10^5$ | $10^5$ | $10^5$ | $2 \times 10^5$ | $10^5$ |
| $m_{\tilde{\chi}^0_1/2}$ | 890 | 3600 | 810 | 3 | 0.1 |
| $m_{\tilde{g}}$ | 4300 | 230 | 160 | 31 | 4 |
| $m_{\tilde{\chi}^\pm_1}$ | $7 \times 10^4$ | $1.5 \times 10^4$ | $1.3 \times 10^4$ | $5600$ | $2100$ |

TABLE I: Model inputs and superpartner spectrum of five representative models. Masses are in GeV. We fix $M_{\tilde{Z}_2}=10^6$ GeV. The masses of the first two generations of squarks and sleptons are typically larger than that of the third. The input parameters $\lambda$, $g_{\tilde{u}}$, and $Y_D, Y_E$ are defined at $\Lambda_S$. The spectra are calculated using exact Renormalization Group Equations (RGE) (see, e.g., [12]). There is a theoretical uncertainty due to multiple RGE thresholds. This mainly affects $m_{\tilde{H}}$, leading to a several GeV uncertainty. The gravitino mass is calculated by $m_{\tilde{g}}=\Lambda_S^3/M_F$ assuming $\Lambda_S \sim \sqrt{F}$. There could be deviations from this relation in some SUSY breaking models which could lead to a gravitino mass that is different by up to a couple orders of magnitude (typically lower). For details, see [11].

The exotic matter in this model is very heavy and does not enter any collider phenomenology.

We have found several regions in the $(Q_Q, Q_2)$ space where a solution satisfies all the constraints. A detailed scan will be presented in a forthcoming publication [11]. The results exhibit a variety of patterns for the low energy spectrum. In Table I we display five representative models. Different ordering of the MSSM gauginos and singlino masses could give rise to very different phenomenology. The singlino mass typically has more variation since it is determined by fine-tuning. The appearance of a light $Z'$ in the spectrum, shown in model 5 (with $\sigma \times B(R(Z' \rightarrow t\bar{t}) \gtrsim 10^6$ fb), could result in a spectacular signal and help untangle the underlying model. This generically happens in the case where the singlino is very light.

A wino as the lightest supersymmetric particle (LSP) and its nearly degenerate charged partner (the degeneracy is lifted at one-loop by about 160 MeV [13] and allows the decay $\tilde{W}^+ \rightarrow \tilde{W}^0 + \pi^+$, which results in a 4 cm displaced vertex) have been studied extensively [14], especially in connection with anomaly mediated models [2]. It can annihilate efficiently into gauge bosons. For pure thermal production the dark matter density is too low for the several hundred GeV mass range we have assumed. However, it can be considerably larger for non-standard cosmological scenarios.

Due to small mixings, at most of the order $\lambda v/\mu \tan \beta$, the decay chain involving the singlino and wino will have a long life-time which could result in a displaced vertex. For example, depending on whether the decay is two or three-body, the life-time for $S \rightarrow h^{(*)} + W^+$ or $W \rightarrow h^{(*)} + S$ is in the range of $10^{-11} - 10^{-19}$ s. This could give an interesting signature in case of $M_2 > M_3$, or $M_5 > M_2$ if the $Z'$ is light enough and has an appreciable branching ratio for decay into the singlino.

There is a wide range of possible gravitino masses, $m_{\tilde{\chi}^\pm_1} \sim 10^{-3} - 10^4$ GeV. With typical assumptions about cosmology, $m_{\tilde{\chi}^\pm_1}$ is strongly constrained by Big Bang Nucleosynthesis (BBN). If the gravitino is not the LSP, we typically require either it to be heavy ($> 10$ TeV) so it decays before BBN, or that the reheating temperature is less than about $10^5 - 10^7$ GeV [15]. In the case that the gravitino is the LSP and the next to lightest supersymmetric particle (NLSP) is the wino, we require the gravitino to be lighter than about $100$ MeV [16]. It is particularly problematic when the singlino is the NLSP since its decay to the gravitino is further suppressed, unless the singlino density is strongly diluted by some late time entropy generation. We also note that decaying into a light gravitino, $m_{\tilde{\chi}^\pm_1} \sim$ MeV, is not observable on collider time scales since the NLSP is neutral.

Since the squarks are heavy the gluino decays off-shell [10]. Its life-time is very sensitive to $g_{\tilde{u}}$ and is given by,

$$\tau_g = 4 \times 10^{-16} \text{sec} \left( \frac{m_{\tilde{Q}}}{10^2 \text{ TeV}} \right)^4 \left( \frac{1 \text{ TeV}}{M_3} \right)^5 \sim \frac{1}{g_{\tilde{u}}^2}. \quad (4)$$

Even though the life-time is long enough for the gluino to hadronize it is too short to result in a displaced vertex.

Since the scalars are heavy, one-loop contributions to most flavor observables (such as $b \rightarrow s \gamma$) are highly suppressed. There are also two loop contributions to EDM and muon $g-2$. However, those are suppressed as compared with the Split SUSY scenario [10] since the Higgsinos are heavy and the singlino-wino mixing is small. The exotic matter in this model is very heavy and does not enter any collider phenomenology.

IV. DISCUSSION

In this letter, we discussed the generic feature of supersymmetry breaking dominantly mediated by an extra $U(1)'$. We have used a $U(1)'$ which forbids a $\mu$ term. Such a requirement gives additional constraints and predicts interesting low energy phenomenology, such as the existence of a light singlino and $Z'$ in various regions of the parameter space. However, $Z'$-mediation is possible in a wider range of $U(1)'$ models, such as $U(1)_{B-L}$. We
expect the hierarchy between the soft scalar masses and the gaugino masses to be generic, although the detailed pattern of soft terms could be quite different. Considering $Z'$-mediation in a broader range of models is certainly worth pursuing.

The model presented here does not provide a seesaw mechanism for neutrino mass. However, in a simple variant the $U(1)'$ symmetry forbids Dirac Yukawa couplings $Y_
u H_u L
u^c$ at the renormalizable level, but allows them to be generated by a higher-dimensional operator [17],

$$W_
u = c_{
u} \frac{S}{M_p} H_u L \nu^c.$$ (5)

This naturally yields small Dirac neutrino masses of order $(0.01 c_{\nu})$ eV for $(S) = 100$ TeV.

There are several scenarios for the decays and lifetimes of the heavy exotic particles [18] and for gauge unification. These depend on the details of the $U(1)'$ charge assignments, and will be discussed in [11].

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[1] A. H. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. 49, 970 (1982); R. Barbieri, S. Ferrara and C. A. Savoy, Phys. Lett. B 119, 343 (1982); N. Ohta, Prog. Theor. Phys. 70, 542 (1983); H. P.Nilles, M. Srednicki and D. Wyler, Phys. Lett. B 120, 346 (1983); E. Cremmer, P. Fayet and L. Girardello, Phys. Lett. B 122, 41 (1983); L. J. Hall, J. D. Lykken and S. Weinberg, Phys. Rev. D 27, 2359 (1983); S. K. Soni and H. A. Weldon, Phys. Lett. B 126, 215 (1983).

[2] L. Randall and R. Sundrum, Nucl. Phys. B 557, 79 (1999) [arXiv:hep-th/9810155]; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998) [arXiv:hep-ph/9810442].

[3] M. Dine and A. E. Nelson, Phys. Rev. D 48, 1277 (1993) [arXiv:hep-ph/9303230]; M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D 51, 1362 (1995) [arXiv:hep-ph/9408384]; M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D 53, 2658 (1996) [arXiv:hep-ph/9507348]; For a review of gauge mediation, see G. F. Giudice and R. Rattazzi, Phys. Rept. 322, 419 (1999) [arXiv:hep-ph/9801271].

[4] Z. Chacko, M. A. Luty, A. E. Nelson and E. Ponton, JHEP 0001, 003 (2000) [arXiv:hep-ph/9911233]; D. E. Kaplan, G. D. Kribs and M. Schmaltz, Phys. Rev. D 62, 035010 (2000) [arXiv:hep-ph/9910293].

[5] J. D. Chung, L. L. Everett, G. L. Kane, S. F. King, J. D. Lykken and L. T. Wang, Phys. Rept. 407, 1 (2005) [arXiv:hep-ph/0312378].

[6] R. Blumenhagen, M. Cvetic, P. Langacker and G. Shiu, Ann. Rev. Nucl. Part. Sci. 55, 71 (2005) [arXiv:hep-th/0502005]; P. Binetruy, G. L. Kane, J. D. Lykken and B. D. Nelson, J. Phys. G 32, 129 (2006) [arXiv:hep-th/0509157].

[7] B. A. Dobrescu, Phys. Lett. B 403, 285 (1997) [arXiv:hep-ph/9703390]; D. E. Kaplan, F. Lepeintre, A. Masiero, A. E. Nelson and A. Riotto, Phys. Rev. D 60, 055003 (1999) [arXiv:hep-ph/9806430]; H. C. Cheng, B. A. Dobrescu and K. T. Matchev, Phys. Lett. B 439, 301 (1998) [arXiv:hep-ph/9807246]; H. C. Cheng, B. A. Dobrescu and K. T. Matchev, Nucl. Phys. B 543, 47 (1999) [arXiv:hep-ph/9811316]; L. L. Everett, P. Langacker, M. Plumacher and J. Wang, Phys. Lett. B 477, 233 (2000) [arXiv:hep-ph/0001073].

[8] E. Accomando et al., arXiv:hep-ph/0608079.

[9] D. Suematsu and Y. Yamagishi, Int. J. Mod. Phys. A 10, 4521 (1995) [arXiv:hep-ph/9411239]; M. Cvetic, D. A. Demir, J. R. Espinosa, L. L. Everett and P. Langacker, Phys. Rev. D 56, 2861 (1997) [Erratum-ibid. D 58, 119905 (1998)] [arXiv:hep-ph/9703317].

[10] N. Arkani-Hamed and S. Dimopoulos, JHEP 0506, 073 (2005) [arXiv:hep-th/0405159]; N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, Nucl. Phys. B 709, 3 (2005) [arXiv:hep-ph/0409232]; P. Gambino, G. F. Giudice and P. Slavich, Nucl. Phys. B 726, 35 (2005) [arXiv:hep-ph/0506214].

[11] P. Langacker, G. Paz, L. T. Wang and I. Yavin, arXiv:0801.3693 [hep-ph].

[12] S. P. Martin and M. T. Vaughn, Phys. Rev. D 50, 2282 (1994) [arXiv:hep-ph/9311340].

[13] D. M. Pierce, J. A. Bagger, K. T. Matchev and R. J. Zhang, Nucl. Phys. B 491, 3 (1997) [arXiv:hep-ph/9606211].

[14] C. H. Chen, M. Drees and J. F. Gunion, Phys. Rev. D 55, 330 (1997) [Erratum-ibid. D 60, 039901 (1999)] [arXiv:hep-ph/9607421]; J. L. Feng, T. Moroi, L. Randall, M. Strassler and S. Su, Phys. Rept. 83, 1731 (1999) [arXiv:hep-ph/9904250]; U. Chattopadhyay, A. C. Gonzalez-Pumplin, D. D. Dam, P. Konar and D. P. Roy, Phys. Rev. D 75, 073014 (2007) [arXiv:hep-ph/0610077]; M. Ibe, T. Moroi and T. T. Yanagida, Phys. Lett. B 644, 355 (2007) [arXiv:hep-ph/0610277].

[15] K. Kohri, T. Moroi and A. Yotsuyanagi, Phys. Rev. D 73, 123511 (2006) [arXiv:hep-ph/0507245].

[16] J. L. Feng, S. Su and F. Takayama, Phys. Rev. D 70, 075019 (2004) [arXiv:hep-ph/0404231].

[17] P. Langacker, Phys. Rev. D 58, 093017 (1998) [arXiv:hep-ph/9805281].

[18] J. Kang, P. Langacker and B. D. Nelson, arXiv:0708.2701 [hep-ph].