Strongly Interacting Gauge Mediation at the LHC

Koichi Hamaguchi\textsuperscript{1,2}, Eita Nakamura\textsuperscript{1}, Satoshi Shirai\textsuperscript{1} and T. T. Yanagida\textsuperscript{1,2}

\textsuperscript{1} Department of Physics, University of Tokyo, Tokyo 113-0033, Japan
\textsuperscript{2} Institute for the Physics and Mathematics of the Universe, University of Tokyo, Chiba 277-8568, Japan

Abstract

Strongly interacting gauge mediation (SIGM) of supersymmetry breaking is very attractive, since it naturally predicts a light gravitino of mass $\lesssim O(10)$ eV, which causes no cosmological problem. We discuss various signatures of the SIGM in the early stage (a low integrated luminosity period) of the LHC experiments. We show, in particular, a possible way to discriminate it from the conventional gauge mediation by counting the number of high $P_T$ leptons.
1 Introduction

Strongly interacting gauge mediation (SIGM) of supersymmetry (SUSY) breaking \cite{1,2} is very attractive, since it naturally predicts a light gravitino of mass \(\lesssim \mathcal{O}(10)\) eV and such a light gravitino is free from all cosmological problems \cite{3,4,5,6}. The SIGM predicts most likely a relatively light gluino compared with the conventional gauge mediation \cite{7} and hence it may have more chance to be tested at the LHC. Therefore, an early SUSY discovery at the LHC with multiple jets and missing transverse momentum \(P_T\) may already indicate the SIGM.

Furthermore, the next-to-lightest SUSY particle (NLSP) is the bino-like neutralino except for a very special case, which dominantly decays into the gravitino emitting a high energy photon. Hence, the SUSY events of multiple jets plus missing \(P_T\) are accompanied with two high \(P_T\) photons. Therefore, the early discovery of events with multiple jets + missing \(P_T\) + two photons provides us with a crucial test of the SIGM.

In this letter, we discuss further tests of the SIGM at the LHC. For definiteness, we adopt an SIGM model to calculate the spectrum of SUSY particles. We find that the gluino becomes lighter than the wino in a large region of parameter space. We find that, as a consequence, the production of high \(P_T\) leptons is surprisingly suppressed at the LHC, contrary to the folklore of the generic SUSY phenomenology.

2 A strongly interacting gauge mediation model

The present model is an extension of the SUSY-breaking SU(5)\textsubscript{hid} model \cite{10} which has quark multiplets \(Q\) transforming as \(5^* + 10\) under the gauge group SU(5)\textsubscript{hid}. We introduce two pairs of massive messengers \(P_d + \bar{P}_d\) and \(P_{\ell} + \bar{P}_{\ell}\) which form \(5 + 5^*\) of the standard-model (SM) SU(5)\textsubscript{GUT}. A crucial assumption of the present SIGM model is that they also belong to \(5 + 5^*\) of the strongly interacting SU(5)\textsubscript{hid} \cite{2} and hence the messengers take

\footnote{Note that the gluino pair production cross section is a steeply falling function of the gluino mass \cite{8}. In the conventional gauge mediation the gluino mass is comparable to the squark mass, while the squark mass should be larger than approximately 1 TeV in order to satisfy the mass bound on the lightest Higgs particle \((m_{h_0} > 114.4\) GeV \cite{9}). (The mass of SM-like Higgs receives a radiative correction from quark-squark loop diagrams.) On the other hand, in the SIGM, squarks are predicted much heavier than the gluino and hence the gluino mass can be in the range below 1 TeV, satisfying the Higgs mass bound.}
part in the strong interactions in the hidden sector. (See Table 1.)

We assume that the gauge coupling of the SU(5)\(_{\text{hid}}\) is strong at the messenger mass scale, since otherwise we may have a split SUSY spectrum with very heavy scalars [2]. Unfortunately, we can not provide a precise prediction of the masses of SUSY particles since the messenger particles are in the strong interactions when they decouple and the integration of the messengers is uncalculable. Therefore, we adopt in this letter the naive dimensional analysis (NDA) [11] to estimate the masses of SUSY particles.

To perform the NDA we assume, for simplicity, that the SUSY is broken mostly by a composite state \(\Phi_S\) which may be a bound state of the quarks \(Q\) (5* + 10). The \(\Phi_S\) is assumed to have both non-vanishing \(A\) and \(F\) terms to represent both the \(R\) symmetry and SUSY breakings, which are known to occur by the strong dynamics [10].

The integration of the strongly interacting SU(5)\(_{\text{hid}}\) sector induces a low-energy effective Kähler potential. The relevant term for the scalar masses in the SUSY standard model (SSM) is given by

\[
K \simeq -\frac{1}{16\pi^2} g_{\text{SM}}^4 \phi_{\text{SM}}^\dagger \phi_{\text{SM}} \frac{1}{16\pi^2} \left( g \Phi_S^\dagger (g \Phi_S) \right) \frac{1}{M_{d/\ell}^2},
\]

where \(\phi_{\text{SM}}\) denote the SM superfields, \(g_{\text{SM}}\) are the SM gauge couplings, and \(g\) is a constant of \(\mathcal{O}(4\pi)\) representing the strong dynamics [11]. \(M_{d/\ell}\) are the effective masses of hadrons \(\Phi_{d/\ell}\) which consist of at least one messenger quark \(P_{d/\ell}\) and SUSY-breaking quarks \(Q\). Here, contributions of higher dimensional terms are ignored, for simplicity. Eq. (1) results in

\[
m^2_\varphi \simeq \left( \frac{\alpha_{\text{SM}}}{4\pi} \right)^2 \frac{|g \langle F_S \rangle|^2}{M_{d/\ell}^2}.
\]

---

2 We assume that the SM gauge symmetry is not broken by the strong dynamics.
3 The \(\Phi_S\) may also contain covariant derivatives.
4 The \(R\) symmetry breaking results in an \(R\)-axion [12], but it is marginally consistent with the constraints (cf. Ref. [13]).
On the other hand, the Kähler potential relevant for the gaugino masses in the SSM is given by

\[ K \simeq g_{\text{SM}}^2 \text{Tr}[W_{\text{SM}}W_{\text{SM}}] \frac{1}{16\pi^2} \frac{(g\Phi_S^\dagger)(g\Phi_S)(g\Phi_S^\dagger)(gD^2\Phi_S)}{M_{d/\ell}^6} + \text{h.c.,} \]  

which leads to

\[ m_\lambda \simeq \frac{\alpha_{\text{SM}}}{4\pi} \frac{|g\langle F_S \rangle|^2 \cdot g\langle \Phi_S^\dagger \rangle g\langle F_S \rangle}{M_{d/\ell}^6}. \]  

Notice that the gaugino masses arise not at \( O(F_S) \) but at \( O(F_S^3) \), since there is no direct coupling between the SUSY breaking fields and the messenger fields [2, 14].

As for the masses of the messengers \( P_d \) and \( P_\ell \) we consider \( m_d = m_\ell \) at the GUT scale.\(^5\) It is very important for our analysis that the SSM gauge interactions increase the value of \( m_d \) more than \( m_\ell \) at low energies. We find, by solving one-loop renormalization group (RG) equations for the messenger masses, \( m_d \simeq 2.5 \times m_\ell \) at the mass of the messenger \( P_d \). Below the threshold of the messenger \( P_d \), only the messenger \( P_\ell \) receives the mass renormalization from the strong SU(5)\(_{\text{hid}} \) interactions and hence the disparity in the messenger masses becomes milder. Furthermore, the hadrons \( \Phi_{d/\ell} \) contain the dynamical quarks \( Q \) besides the messenger quark \( P_{d/\ell} \) and hence their mass ratio \( M_d/M_\ell \) is smaller than the mass ratio of messenger quarks \( m_d/m_\ell \). Since the mass ratio is uncalculable due to the strong dynamics, we simply introduce one parameter \( \kappa_1 = M_d/M_\ell \) and consider a region of \( 1 \lesssim \kappa_1 \lesssim 2 \).

Another important point is that the powers of \( M_{d/\ell} \) are different between gauginos and scalars. Assuming that \( g\langle \Phi_S \rangle = \Lambda \) and \( g\langle F_S \rangle = \Lambda^2 \), we obtain

\[ m^2_\varphi \simeq \left( \frac{\alpha_{\text{SM}}}{4\pi} \right)^2 \frac{\Lambda^4}{M_{d/\ell}^2}, \quad m_\lambda \simeq \frac{\alpha_{\text{SM}}}{4\pi} \frac{\Lambda^7}{M_{d/\ell}^6}. \]  

Here, \( \Lambda \) denotes the hadron mass scale. We introduce one more parameter \( \kappa_2 = M_d/\Lambda \), which is of order unity and satisfies \( \kappa_2 \gtrsim 1 \). Note that \( F_S = \Lambda^2/g \simeq \Lambda^2/4\pi \) yields a relatively small gravitino mass compared with the case of \( F_S \simeq \Lambda^2 \) (i.e., \( g \simeq 1 \)).

\(^5\)\( m_{d/\ell} \) are the masses of constituent messenger quarks \( P_{d/\ell} \), while \( M_{d/\ell} \) are the masses of composite hadrons \( \Phi_{d/\ell} \).

\(^6\)Here, we assume \( m_d = \mathcal{O}(100) \text{ TeV} \) to realize \( m_{3/2} = \mathcal{O}(1) \text{ eV} \) and \( m_{\text{gaugino}} = \mathcal{O}(100) \text{ GeV} \) as we shall see later.
3 Spectrum of the SUSY particles

The low energy spectrum of the SSM particles can be obtained by using Eq. (5) at the messenger mass scale and then RG-evolving their masses from the messenger mass scale down to the weak scale. As discussed in the previous section, we introduce two parameters \( \kappa_1 = M_d/M_\ell \) and \( \kappa_2 = M_d/\Lambda \) to represent the uncertainties arising from the strong dynamics. We investigate the regions

\[
1 \leq \kappa_1 \leq 2 \quad \text{and} \quad \kappa_2 \geq 1.
\]

The explicit mass formulae at the messenger mass scale are

\[
m_\tilde{g} = \frac{\alpha_3 \Lambda}{4\pi \kappa_2^6}, \quad m_\tilde{W} = \frac{\alpha_2 \Lambda}{4\pi \kappa_2^6} \kappa_1^6, \quad m_\tilde{B} = \frac{\alpha_1 \Lambda}{4\pi \kappa_2^6} \left[ \frac{1}{3} + \frac{1}{2} \kappa_1^6 \right]
\]

for the gauginos and

\[
m_\varphi^2 = \frac{\Lambda^2}{\kappa_2^2} \left\{ C_{\text{SU(3)}_C}^\varphi \left( \frac{\alpha_3}{4\pi} \right)^2 + C_{\text{SU(2)}_L}^\varphi \left( \frac{\alpha_2}{4\pi} \right)^2 \kappa_1^2 + C_{\text{U(1)}_Y}^\varphi \left( \frac{\alpha_1}{4\pi} \right)^2 \left[ \frac{1}{3} + \frac{1}{2} \kappa_1^2 \right] \right\}
\]

for the scalar particles, where \( \alpha_i = g_i^2/4\pi \) are the SM gauge couplings and \( C_{\text{SU(3)}_C}^\varphi, C_{\text{SU(2)}_L}^\varphi \) and \( C_{\text{U(1)}_Y}^\varphi \) are the quadratic Casimir invariants of the corresponding gauge groups for the field \( \varphi \). We should mention as a reminder that there exist \( O(1) \) uncertainties in these formulae.

Before looking at the numerical results, we note some general features of the mass spectrum which can be read from the above formulae. The parameter \( \kappa_2 \) dictates the hierarchy between the gaugino masses and the scalar masses. As the value of \( \kappa_2 \) increases, the scalars become heavier than the gauginos. The parameter \( \kappa_1 \) determines the relative mass relations between \( \text{SU(3)}_C \) charged particles (\( \tilde{g} \) and \( \tilde{q} \)) and the \( \text{SU(3)}_C \) neutral particles (\( \tilde{\ell}, \tilde{\chi} \) and \( H_{u,d} \)). As the value of \( \kappa_1 \) increases, \( \text{SU(3)}_C \) neutral particles become heavier. In the limit \( \kappa_1, \kappa_2 \to 1 \), we recover the GUT relation for the gaugino masses and \( m_{\text{scalar}} \sim m_{\text{gaugino}} \) as in the conventional gauge mediation model. \( \Lambda \) provides the overall scale for the soft SUSY breaking masses.

As remarked earlier, the gravitino is very light in our model. Its mass is given by

\[
m_{3/2} = \frac{\langle F_S \rangle}{\sqrt{3}M_P} \simeq \frac{\Lambda^2}{4\pi \sqrt{3}M_P},
\]

\[\text{In our normalization of hypercharge, } \alpha_1 = \frac{5}{3} \frac{g' \cos \theta_W}{4\pi} \text{ with } g' \cos \theta_W = e \text{ and } C_{\text{U(1)}_Y}^\varphi = \frac{1}{3} \left( Q_{\text{em}} - T_{\text{SU(2)}}^3 \right)^2.\]
where $M_P = 2.44 \times 10^{18}$ GeV is the reduced Planck mass. We see that $m_{3/2} < 10$ eV corresponds to $\Lambda \lesssim 730$ TeV.

An important issue is the occurrence of the electroweak (EW) symmetry breaking.\footnote{We do not discuss the origin of the $\mu$- and $B$-terms (the $\mu$-problem) in the present model and take the $\tan \beta$ as a free parameter.} In our model, the Higgs scalars can be relatively heavy and accordingly, we must check whether the squared Higgs masses become so small at the weak scale that the EW symmetry breaking occurs. In fact, as the value of $\kappa_1$ increases, the soft SUSY breaking Higgs masses squared $m_{H_{u,d}}^2$ become larger and there is an upper bound on $\kappa_1$ (for each value of $\kappa_2$) above which no EW symmetry breaking occurs.

Now we present the numerical results. We calculate the SUSY-particle masses by using the one-loop RG equations for the SSM parameters. In Figs. 1 and 2, we show the gluino mass ($m_{\tilde{g}}$), the first two lightest neutralino masses ($m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_2^0}$) and the lighter stau mass ($m_{\tilde{\tau}_1}$) as functions of $\kappa_1$. Two examples for $(\kappa_2, \Lambda) = (1.35, 280$ TeV) and $(1.8, 900$ TeV) are shown in Figs. 1 and 2 respectively. Here, $\tan \beta = 10$ in both cases. The upper bound of $\kappa_1$ in each figure corresponds to the point where $\mu = 0$ above which EW symmetry breaking does not occur.

The general features of the mass spectrum discussed above can be seen in those examples. The NLSP is the lightest neutralino in most of the parameter region.\footnote{The lightest neutralino is bino-like and the second lightest neutralino is wino-like in a large region of the parameter space where $\kappa_1$ is not near the upper bound value. As $\kappa_1$ approaches the upper bound, $\mu$ becomes much smaller than $m_{\tilde{B}}$ and $m_{\tilde{W}}$, and the higgsino components dominate in the first two lightest neutralino, whose masses approach zero when $\kappa_1$ is near the upper bound.} The GUT relation among the gaugino masses, which holds in the conventional gauge mediation models, is violated for $\kappa_1 > 1$. In particular, for $\kappa_1 \gtrsim 1.2$, $m_{\tilde{W}}$ becomes larger than $m_{\tilde{g}}$.

4 \textbf{Test for $m_{\tilde{W}} > m_{\tilde{g}}$ at the LHC}

If the SIGM is the case and the neutralino is the NLSP, its experimental signatures are high $P_T$ photons, high $P_T$ multiple jets and large missing $P_T$. However, the conventional
Figure 1: Mass spectrum for $\Lambda = 280$ TeV and $\kappa_2 = 1.35$. Here, scalar quark masses are approximately 1.5 TeV. In the shaded region, EW breaking does not occur.

Figure 2: Mass spectrum for $\Lambda = 900$ TeV and $\kappa_2 = 1.8$. Here, scalar quark masses are approximately 3.5 TeV. In the shaded region, EW breaking does not occur.
minimal gauge mediated SUSY breaking (mGMSB) models with ultralight gravitino LSP and neutralino NLSP will also have similar signatures. Here, we discuss possible ways to discriminate between those two models.

We have seen that the SIGM predicts heavier scalar particles and violation of GUT relation for the gaugino masses. Therefore, we can straightforwardly discriminate the SIGM models from the mGMSB ones by measuring the SUSY particles’ masses. However, it may require a large integrated luminosity. In the following, we propose another way to discriminate the SIGM at an earlier stage of the LHC experiments.

At the LHC, if the SIGM is the case, SUSY particles are mainly produced through:

\[ pp \rightarrow \tilde{g} + \tilde{g} + X. \]  

In the case that \( m_{\tilde{W}} > m_{\tilde{g}} \), the gluinos dominantly decay into \( \tilde{B} + q + \bar{q} \), \( \tilde{B} + g \) and \( \tilde{B} + t + \bar{t} \) if not kinematically forbidden. Then, the produced binos \( \tilde{B} \) dominantly decay into \( \gamma + \tilde{G}_{3/2} \), and \( Z^0 + \tilde{G}_{3/2} \) if not kinematically forbidden (\( \tilde{G}_{3/2} \) denotes the gravitino). Thus if \( m_{\tilde{W}} > m_{\tilde{g}} \) is the case, there are no high \( P_T \) leptons except for ones which come from \( Z^0 \)'s and \( t \)'s decays.

On the other hand, in the case of mGMSB, many lepton production channels exist since the wino (and sleptons) are lighter than the gluino. Therefore, we can distinguish between mGMSB and SIGM by counting the number of high \( P_T \) leptons. We consider two examples, by taking an SIGM \((m_{3/2} = 10 \text{ eV}, \kappa_1 = 1.35, \kappa_2 = 1.5, \tan \beta = 10)\) and an mGMSB \((F/M_{\text{mess}} = 80 \text{ TeV}, M_{\text{mess}} = 160 \text{ TeV}, N_5 = 1, \tan \beta = 10)\). The mass spectrums are shown in Figs. 3 and 4. These spectrums are calculated by ISAJET7.72 [15]. To simulate LHC signatures for these models, we use programs Herwig 6.5 [16] and AcerDET-1.0 [17].

We take the events cuts as follows:

- \( \geq 4 \) jets with \( P_T > 50 \text{ GeV} \) and \( P_{T,1,2} > 100 \text{ GeV} \).
- \( \geq 2 \) photons with \( P_T > 10 \text{ GeV} \) and \( P_{T,1} > 20 \text{ GeV} \).
- \( M_{\text{eff}} > 500 \text{ GeV} \), where

\[ M_{\text{eff}} = \sum_{\text{jets}} P_{T,j} + P_{T,\text{miss}}. \]  

\(^{10}\) This mGMSB example has a light gluino, but it does not satisfy the Higgs mass bound. (See the discussion in Sec. [1]) We take this model point just as a demonstration, for a comparison to the SIGM.
Figure 3: SIGM mass spectrum for $m_{3/2} = 10$ eV, $\kappa_1 = 1.35$, $\kappa_2 = 1.5$ and $\tan \beta = 10$.

Figure 4: mGMSB mass spectrum $F/M_{\text{mess}} = 80$ TeV, $M_{\text{mess}} = 160$ TeV, $N_5 = 1$ and $\tan \beta = 10$. 
Figure 5: The lepton number distributions for mGMSB and SIGM.

- $P_{T,\text{miss}} > 0.2 M_{\text{eff}}$.

Under these cuts, we see that the standard model backgrounds are almost negligible.

The number distributions for high $P_T$ leptons (with $P_T > 20$ GeV) are shown in Fig. 5. In Fig. 5, we can see there are very few leptons produced at the LHC in the SIGM. The photon cuts reduce the lepton production from $\tilde{B}$'s decay, and the decay $\tilde{g} \rightarrow \tilde{B} + t + \bar{t}$ is kinematically forbidden in the SIGM example.

In general, larger $\kappa_1$ results in the smaller production rate for high $P_T$ leptons. To see the relation of $\kappa_1$ to the lepton production rate, we define $R$ as

$$R \equiv \frac{\# \text{ of events after the cuts with at least one lepton } (P_T > 20 \text{ GeV})}{\# \text{ of all events after the cuts}}. \quad (12)$$

For the conventional mGMSB with the parameters given above, we obtain $R = 0.40$. In Figs. 6 and 7, we show the $\kappa_1$ dependence of $R$ and $(m_{\tilde{g}}, m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_2^0})$ in the SIGM ($m_{3/2} = 10$ eV, $\kappa_2 = 1.5$ and $\tan \beta = 10$), respectively. To compute $R$, 20000 SUSY events are generated for each value of $\kappa_1$. Fig. 6 clearly shows that there are less lepton production for larger $\kappa_1$. It is also confirmed that in the case $m_{\tilde{W}} \gtrsim m_{\tilde{g}}$ almost no leptons are produced. Thus, we may distinguish between mGMSB and SIGM by counting the number of high $P_T$ leptons.
Figure 6: The relation of $\kappa_1$ to $R$. SIGM parameters are the same as in Fig. 3 except for $\kappa_1$. Error bars represent only the statistical errors.

Figure 7: The relation of $\kappa_1$ to $m_{\tilde{g}}$ and $m_{\tilde{\chi}^0_{1,2}}$. 
5 Discussion

The strongly interacting gauge mediation (SIGM) predicts gaugino masses without the GUT relation, in particular a relatively light gluino compared with a conventional gauge mediation. We have calculated the SUSY mass spectrum by taking some explicit examples, and shown that the SIGM with such a mass spectrum can be discriminated simply by counting the number of high $P_T$ leptons.

If many events with two photons, multiple jets and missing energy are discovered at the LHC, it naturally points to gauge mediation models with a neutralino NLSP and the gravitino LSP. If the number of such events is large, which suggests the light gluino or squark, it may already indicate an unconventional gauge mediation mass spectrum even at an early stage of the LHC experiments.\footnote{\textsuperscript{11}We can then test a peculiar mass hierarchy among gauginos by simply counting the number of high $P_T$ leptons, as investigated in this letter.}

We have also found that the gluino can even be the NLSP for some parameter region. In this case the main SUSY event signal will be two jets + missing energy. In such a case, it will be challenging to identify the LSP (gravitino) and the NLSP (gluino).

Acknowledgements

This work was supported by World Premier International Center Initiative (WPI Program), MEXT, Japan. The work of TTY is supported in part by the Grant-in-Aid for Science Research, Japan Society for the Promotion of Science, Japan (No. 1940270). The work by KH is supported by JSPS (18840012). The work of SS is supported in part by JSPS Research Fellowships for Young Scientists.

References

[1] K. I. Izawa and T. Yanagida, Prog. Theor. Phys. 114 (2005) 433 \texttt{arXiv:hep-ph/0501254}.

\footnote{\textsuperscript{11}Note that the gluino and squarks are heavy in the conventional GMSB models. See footnote\textsuperscript{11}}
[2] M. Ibe, Y. Nakayama and T. T. Yanagida, Phys. Lett. B 649 (2007) 292 [arXiv:hep-ph/0703110];
M. Ibe, Y. Nakayama and T. T. Yanagida, [arXiv:0804.0636 [hep-ph]].

[3] H. Pagels and J. R. Primack, Phys. Rev. Lett. 48 (1982) 223;
M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese and A. Riotto, Phys. Rev. D 71 (2005) 063534 [arXiv:astro-ph/0501562].

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

[5] M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B 638 (2006) 8 [arXiv:hep-ph/0603265];
M. Endo, F. Takahashi and T. T. Yanagida, Phys. Rev. D 76 (2007) 083509 [arXiv:0706.0986 [hep-ph]].

[6] M. Endo, K. Hamaguchi and F. Takahashi, Phys. Rev. Lett. 96 (2006) 211301 [arXiv:hep-ph/0602061];
S. Nakamura and M. Yamaguchi, Phys. Lett. B 638 (2006) 389 [arXiv:hep-ph/0602081].

[7] See for a review, G. F. Giudice and R. Rattazzi, Phys. Rept. 322 (1999) 419 [arXiv:hep-ph/9801271].

[8] W. Beenakker, R. Hopker, M. Spira and P. M. Zerwas, Nucl. Phys. B 492 (1997) 51 [arXiv:hep-ph/9610490].

[9] R. Barate et al. [LEP Working Group for Higgs boson searches], Phys. Lett. B 565, 61 (2003) [arXiv:hep-ex/0306033];
S. Schael et al. [LEP Working Group for Higgs boson searches], Eur. Phys. J. C 47, 547 (2006) [arXiv:hep-ex/0602042].

[10] I. Affleck, M. Dine and N. Seiberg, Phys. Lett. B 137 (1984) 187;
Y. Meurice and G. Veneziano, Phys. Lett. B 141 (1984) 69;
H. Murayama, Phys. Lett. B 355 (1995) 187 [arXiv:hep-th/9505082].
[11] M. A. Luty, Phys. Rev. D 57 (1998) 1531 [arXiv:hep-ph/9706235];
    A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. B 412 (1997) 301
    [arXiv:hep-ph/9706275].

[12] A. E. Nelson and N. Seiberg, Nucl. Phys. B 416 (1994) 46 [arXiv:hep-ph/9309299];
    J. Bagger, E. Poppitz and L. Randall, Nucl. Phys. B 426 (1994) 3
    [arXiv:hep-ph/9405345].

[13] L. J. Hall and T. Watari, Phys. Rev. D 70 (2004) 115001 [arXiv:hep-ph/0405109].

[14] N. Arkani-Hamed, G. F. Giudice, M. A. Luty and R. Rattazzi, Phys. Rev. D 58
    (1998) 115005 [arXiv:hep-ph/9803290].

[15] F. E. Paige, S. D. Protopopescu, H. Baer and X. Tata, [arXiv:hep-ph/0312045].

[16] G. Marchesini, B. R. Webber, G. Abbiendi, I. G. Knowles, M. H. Seymour and
    L. Stanco, Comput. Phys. Commun. 67 (1992) 465;
    G. Corcella et al., JHEP 0101 (2001) 010 [arXiv:hep-ph/0011363];
    G. Corcella et al., [arXiv:hep-ph/0210213].

[17] E. Richter-Was, [arXiv:hep-ph/0207355].