Probing the Messenger of SUSY Breaking with Gaugino Masses

Won Sang Cho† and Kiwoon Choi‡

Department of Physics, KAIST, Daejeon 305-701, Korea
(Dated: February 1, 2008)

Abstract

Gaugino masses might provide useful information on the underlying scheme of supersymmetry breaking as they are least dependent on the unknown physics between the TeV scale and the high messenger scale of supersymmetry breaking. We discuss the pattern of low energy gaugino masses in various schemes of supersymmetry breaking together with the possibility to determine the gaugino masses at LHC.

PACS numbers: 12.60.Jv, 14.80.Ly

* Talk given by K. Choi at International Workshop on Theoretical High Energy Physics, March 2007, Roorkee, India
† email: wscho@hep.kaist.ac.kr
‡ email: kchoi@hep.kaist.ac.kr
I. INTRODUCTION

Low energy supersymmetry (SUSY) \cite{1} is one of the prime candidates for physics beyond
the standard model at the TeV scale. Most phenomenological aspects of low energy SUSY
are determined by the soft SUSY breaking terms in low energy effective lagrangian. Those
soft terms are generated at certain messenger scale $M_{\text{mess}}$ presumed to be higher than TeV,
and then receive quantum corrections due to the renormalization group (RG) evolution
and threshold effects that might occur at scales below $M_{\text{mess}}$. Among all the soft terms, the
MSSM gaugino masses $M_a$ ($a = SU(3), SU(2), U(1)$) appear to be the least model dependent
as they are related to the corresponding standard model (SM) gauge coupling constants $g_a$ in
a nontrivial manner. Specifically, $M_a/g_a^2$ do not run at the one-loop level, and also possible
intermediate threshold corrections to $M_a/g_a^2$ are severely constrained if one requires to keep
the successful gauge coupling unification at $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV. In this respect, analysis
of the gaugino mass pattern at TeV can be considered as a promising first step to uncover
the mediation mechanism of SUSY breaking at $M_{\text{mess}}$. In this talk, we discuss the possible
pattern of low energy gaugino masses which might be obtained in various SUSY breaking
schemes \cite{2} and also the possibility to determine the gaugino masses at LHC \cite{3, 4, 5}, aiming
to see what kind of information on SUSY breaking scheme can be extracted once the low
energy gaugino masses can be determined by future collider experiments.

II. GENERIC GAUGINO MASSES IN 4D SUPERGRAVITY

In 4D effective supergravity (SUGRA) with the cutoff scale $\Lambda$ which is chosen to be
just below the string or Kaluza-Klein (KK) or GUT threshold scale, the running gauge
couplings and gaugino masses at a scale $\mu$ below $\Lambda$ (but above the next threshold scale $M_{\text{th}}$)
are determined by the gauge coupling superfield $F_a(p^2)$ ($M_{\text{th}}^2 < p^2 < \Lambda^2$) in the quantum
effective action of gauge superfields:

$$\Gamma = \int d^4p d^4\theta \left( \frac{1}{4} F_a(p^2) W^a_{\alpha\dot{\alpha}} \frac{D^\alpha D_{\dot{\alpha}}}{16 p^2} W_a^a + \text{h.c} \right),$$

(1)
where $W^a$ denote the chiral gauge superfields and $D_\alpha$ is super-covariant derivative. At one-loop approximation, $\mathcal{F}_a$ is given by

$$
\mathcal{F}_a(p^2) = \text{Re}(f_a^{(0)}) - \frac{1}{16\pi^2} (3C_a - \sum_i C^i_a) \ln \left( \frac{CC^*\Lambda^2}{p^2} \right)
+ \frac{1}{8\pi^2} \sum_i C^i_a \ln \left( e^{-K_0/3}Z_i \right) + \frac{1}{8\pi^2} \Omega_a,
$$

where $f_a^{(0)}$ are the tree-level holomorphic gauge kinetic function, $C_a$ and $C^i_a$ are the quadratic Casimir of the gauge multiplet and the gauge-charged matter superfield $Q_i$, respectively, and $C$ is the chiral compensator of 4D SUGRA. Here $K_0(X_I, X^*_I)$ is the Kähler potential of generic SUSY breaking (moduli or matter) superfields $X_I$ which have nonzero $F$-components $F^I$, $Z_i(X_I, X^*_I)$ is the Kähler metric of $Q_i$, and $\Omega_a$ include the string, KK and GUT threshold corrections as well as the (regularization scheme-dependent) field-theoretic one-loop part: $\frac{1}{8\pi^2} C_a \ln[\text{Re}(f_a^{(0)})]$. In the one-loop approximation, $\Omega_a$ are independent of the external momentum $p^2$, thus independent of $C$ as a consequence of the super-Weyl invariance. However $\Omega_a$ generically depend on SUSY breaking fields $X_I$, and a full determination of their $X_I$-dependence requires a detailed knowledge of the UV physics above $\Lambda$. From the above gauge coupling superfield, one easily finds that the running gauge couplings and gaugino masses at $\mu$ ($M_{\text{th}} < \mu < \Lambda$) are given by

$$
\frac{1}{g_a^2(\mu)} = \mathcal{F}_a|_{C = e^{K_0/6}, p^2 = \mu^2}
= \text{Re}(f_a^{(0)}) - \frac{1}{16\pi^2} \left[ (3C_a - \sum_i C^i_a) \ln \left( \frac{\Lambda^2}{\mu^2} \right) + (C_a - \sum_i C^i_a)K_0 + 2 \sum_i C^i_a \ln Z_i \right] + \frac{1}{8\pi^2} \Omega_a,
$$

$$
\frac{M_a(\mu)}{g_a^2(\mu)} = \left. F^A\partial_A\mathcal{F}_a \right|_{C = e^{K_0/6}, p^2 = \mu^2}
= F^I \left[ \frac{1}{2} \partial_I f_a^{(0)} - \frac{1}{8\pi^2} \sum_i C^i_a \partial_I \ln(e^{-K_0/3}Z_i) + \frac{1}{8\pi^2} \partial_I \Omega_a \right] - \frac{1}{16\pi^2} (3C_a - \sum_i C^i_a) \frac{F^C}{C},
$$

where $F^A = (F^C, F^I)$, $\partial_A = (\partial_C, \partial_I)$, $\frac{F^C}{C} = m_{3/2}^a + \frac{1}{3} F^I \partial_I K_0$, and $C = e^{K_0/6}$ corresponds to the Einstein frame condition.

The above expression of $M_a/g_a^2$ is valid at any scale between $M_{\text{th}}$ and $\Lambda$. However, depending upon the SUSY breaking scenario, $M_a/g_a^2$ can receive important threshold correction at
the next threshold scale $M_{\text{th}}$. To see how $M_a/g_a^2$ are modified by threshold effect, let us assume \{Q_i\} $\equiv$ \{Φ + Φc, Q_x\} and Φ + Φc get a supersymmetric mass of the order of $M_{\text{th}}$, while $Q_x$ remain to be massless at $M_{\text{th}}$. Then Φ + Φc can be integrated out to derive the low energy parameters at scales below $M_{\text{th}}$. The relevant couplings of Φ + Φc at $M_{\text{th}}$ can be written as

$$
\int d^4\theta CC^*e^{-K_0/3}(Z_\Phi\Phi^*\Phi + Z_\Phi^c\Phi^c\Phi^c) + \left(\int d^2\theta C^3\lambda_\Phi X_\Phi\Phi^c\Phi + \text{h.c}\right), \tag{4}
$$

where $X_\Phi$ is assumed to have a vacuum value $\langle X_\Phi \rangle = M_\Phi + \theta^2 F_X$ for which the physical mass of Φ + Φc is given by $M_\Phi = \lambda_\Phi CX_\Phi/\sqrt{e^{-2K_0/3}Z_\Phi Z_\Phi^c}$. Then, integrating out Φ + Φc yields a threshold correction to $F_a$:

$$
\Delta F_a(M_{\text{th}}) = -\frac{1}{8\pi^2} \sum_\Phi C_\Phi a \ln\left(\frac{M_\Phi M_\Phi^c}{M_{\text{th}}^2}\right), \tag{5}
$$

which results in the threshold correction to gaugino mass at $M_{\text{th}}$:

$$
M_a(M_{\text{th}}^-) - M_a(M_{\text{th}}^+) = g_a^2(M_{\text{th}})F_A\partial_A\Delta F_a
= -\frac{g_a^2(M_{\text{th}})}{8\pi^2} \sum_\Phi C_\Phi a \left(\frac{F_\Phi^c}{C} + \frac{F_X^c}{M_\Phi} - F_I^I \partial_I \ln(e^{-2K_0/3}Z_\Phi Z_\Phi^c)\right). \tag{6}
$$

Including this threshold, one finds

$$
\left(\frac{M_a}{g_a^2}\right)_{M_{\text{th}}^-} = \left(\begin{array}{c}
F_I^I \left[\frac{1}{2} \partial_I f_a^{(0)} - \frac{1}{8\pi^2} \sum_x C_x^a \partial_I \left(e^{-K_0/3}Z_x\right) + \frac{1}{8\pi^2} \partial_I \Omega_a\right]
- \frac{1}{8\pi^2} \sum_\Phi C_\Phi a \frac{F_X^c}{M_\Phi} - \frac{1}{16\pi^2} (3C_a - \sum_x C_x^a) \frac{F_\Phi^c}{C},
\end{array}\right) \tag{7}
$$

where $\sum_x$ denotes the summation over \{Q_x\} which remain as light matter fields at $M_{\text{th}}^-$.

One can repeat the above procedure, i.e. run down to the lower threshold scale, integrate out the massive fields there, and then include the threshold correction to gaugino masses until one arrives at the TeV scale. Then one finally finds

$$
\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} = \tilde{M}_a^\text{moduli} + \tilde{M}_a^gauge + \tilde{M}_a^\text{conformal} + \tilde{M}_a^\text{konishi} + \tilde{M}_a^\text{UV} \tag{8}
$$
where

\[
\tilde{M}_a^{\text{moduli}} = \frac{1}{2} F^I \partial_I f_a^{(0)},
\]

\[
\tilde{M}_a^{\text{gauge}} = -\frac{1}{8\pi^2} \sum_\Phi C_\Phi \frac{F^{X_\Phi} M_\Phi}{M_\Phi},
\]

\[
\tilde{M}_a^{\text{conformal}} = \frac{1}{16\pi^2} b_a \frac{F^C C}{C},
\]

\[
\tilde{M}_a^{\text{konishi}} = -\frac{1}{8\pi^2} \sum_m C_m^a F^I \partial_I \ln(e^{-K_0/3 Z_m}),
\]

\[
\tilde{M}_a^{\text{UV}} = \frac{1}{8\pi^2} F^I \partial_I \Omega_a,
\]

(9)

where \(\sum_m\) denotes the summation over the light matter multiplets \(\{Q_m\}\) at the TeV scale, \(\sum_\Phi\) denotes the summation over the gauge messenger fields \(\Phi + \Phi^c\) which have a mass lighter than \(\Lambda\) but heavier than TeV, and \(b_a = -3C_a + \sum_m C_m^a\) are the one-loop beta-function coefficients at TeV. Here \(\tilde{M}_a^{\text{moduli}}\) denotes the moduli-mediated tree level value of \(M_a/g_a^2\) \[8\], \(\tilde{M}_a^{\text{gauge}}\) is the intermediate scale gauge threshold due to gauge-charged massive particles with a mass between \(\Lambda\) and TeV \[9\], \(\tilde{M}_a^{\text{conformal}}\) is the anomaly-mediated contribution determined by the conformal anomaly at TeV \[7\], \(\tilde{M}_a^{\text{konishi}}\) is a piece determined by the Konishi anomaly \[10\], and finally \(\tilde{M}_a^{\text{UV}}\) contains the UV thresholds at string, KK and GUT scales.

Formulae (8) and (9) give the most general description of gaugino masses and its origin from the underlying schemes \[8\]. Depending upon the SUSY breaking scenario, \(M_a/g_a^2\) are dominated by some of these five contributions. The SM gauge coupling constants at TeV have been measured with the (approximate) result: \(g_1^2 : g_2^2 : g_3^2 \simeq 1 : 2 : 6\). As a result, once the gaugino mass ratios at TeV are measured, the ratios of \(M_a/g_a^2\) at TeV can be experimentally determined, which will allow us to test SUSY breaking schemes using the predicted pattern of low energy gaugino masses.

As \(f_a^{(0)}\) determine the gauge coupling constants at \(M_{GUT}\), it is expected that \(\tilde{M}_a^{\text{moduli}}\) are universal in most cases realizing the gauge coupling unification at \(M_{GUT}\). In compactified string theory or higher dimensional SUGRA, the tree level gauge kinetic functions are generically given by \(f_a^{(0)} = \sum_I k_{aI} X_I\), where \(X_I\) correspond to the dilaton and/or moduli superfields and \(k_{aI}\) are rational numbers. In models realizing gauge coupling unification, \(k_{aI}\) are universal for the SM gauge group \(a = SU(3), SU(2), U(1)\), which would give universal \(\tilde{M}_a^{\text{moduli}} = \frac{1}{2} \sum_I k_{aI} F^I\).

The intermediate gauge thresholds \(\tilde{M}_a^{\text{gauge}} = -\frac{1}{8\pi^2} \sum_\Phi C_\Phi \frac{F^{X_\Phi} M_\Phi}{M_\Phi}\) accompany the additional
running of gauge couplings from $M_{\text{GUT}}$ to $M_{\Phi}$: $\Delta(1/\ g_a^2) = \frac{1}{4\pi^2} \sum_{\Phi} c_{a}^{\Phi} \ln(M_{\text{GUT}}/M_{\Phi})$, indicating that $\tilde{M}_a^{\text{gauge}}$ with $M_{\Phi} \ll M_{\text{GUT}}$ are required to be universal also to keep the gauge coupling unification at $M_{\text{GUT}}$. On the other hand, there is no good reason to expect that the string, KK and GUT thresholds $\tilde{M}_a^{\text{UV}}$ are universal. In fact, the UV thresholds encoded in $\frac{1}{8\pi^2}\Omega_a$ are most model-dependent, and difficult to compute. If this part gives an important contribution to $M_a/g_a^2$, it is difficult to make a model-independent statement about the gaugino masses.

With the above observation, one can consider the following three distinctive patterns of low energy gaugino masses which can result from theoretically well motivated setup.

**mSUGRA pattern**: The scenario which has been discussed most often in the literatures is that $(M_a/g_a^2)_{\text{TeV}}$ are dominated by $\tilde{M}_a^{\text{moduli}}$ or $\tilde{M}_a^{\text{gauge}}$ which are assumed to be universal:

$$\left(\frac{M_a}{g_a^2}\right)_{\text{TeV}} \simeq \tilde{M}_a^{\text{moduli}} \text{ or } \tilde{M}_a^{\text{gauge}},$$

leading to the following low energy gaugino mass ratios:

$$M_1 : M_2 : M_3 \simeq 1 : 2 : 6 \quad (11)$$

which will be termed mSUGRA pattern in the following. Schemes giving the mSUGRA pattern of gaugino masses include the dilaton and/or moduli dominated SUSY breaking scenarios realized in various compactified string theories [8] with a large string and compactification scales near $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV, gaugino mediation scenario [11], and also gauge mediation scenario [9] with a messenger scale $M_{\text{mess}} \ll M_{\text{GUT}}$.

We stress that the universality of $\tilde{M}_a^{\text{moduli}}$ which is essential for the mSUGRA pattern heavily relies on the assumption of high scale gauge coupling unification. In models without gauge coupling unification, $k_{aI}$ and thus $\tilde{M}_a^{\text{moduli}}$ are generically non-universal and highly model-dependent. However, in some case, one might be able to extract information on $k_{aI}$ for $X_I$ providing a dominant source of SUSY breaking, thereby make a certain prediction on low energy gaugino masses. A nontrivial example of this kind is the large volume compactification of Type IIB string theory proposed in [12]. In the model of [12], moduli are stabilized at a vacuum with the string scale $M_{\text{st}} \sim 10^{11}$ GeV, and $f_a^{(0)} = k_aT_s + h_aS$, where $T_s$ is the volume modulus of small 4-cycle and $S$ is the IIB dilaton with $|F^S| \ll |F^{T_s}|$. As $\text{Re}(f_a^{(0)}) \simeq 1/g_a^2(M_{\text{st}})$, in such intermediate string scale scenario, $k_a$ and $h_a$ can not be constrained by gauge coupling unification. However, the model of [12] has $k_{SU(3)} = k_{SU(2)}$. 


regardless of the values of $g_a^2(M_{st})$, while $k_{U(1)}$ and $h_a$ are generically non-universal independent parameters. As $F_S$ and $F^C$ ($C =$ SUGRA compensator) are negligible in the model of [12], $k_{SU(3)} = k_{SU(2)}$ leads to $M_2 : M_3 \simeq 1 : 3$, while the ratio with $M_1$ depends on the unknown $k_{U(1)}/k_{SU(2)}$.

**Mirage pattern:** Another interesting scenario is that $(M_a/g_a^2)_{\text{TeV}}$ are dominated by $\tilde{M}_a^\text{conformal}$ and universal $\tilde{M}_a^\text{moduli}$ (or $\tilde{M}_a^\text{gauge}$) which are comparable to each other:

$$\left( \frac{M_a}{g_a^2} \right)_{\text{TeV}} \simeq \left( \tilde{M}_a^\text{moduli} \text{ or } \tilde{M}_a^\text{gauge} \right) + \tilde{M}_a^\text{conformal},$$

leading to

$$M_1 : M_2 : M_3 \simeq (1 + 0.66\alpha) : (2 + 0.2\alpha) : (6 - 1.8\alpha),$$

where $\alpha$ is a *positive parameter of order unity* defined as

$$\frac{g_{\text{GUT}}^2 b_a \ln(M_{\text{Pl}}/m_{3/2})}{16\pi^2} \equiv \frac{\tilde{M}_a^\text{conformal}}{\tilde{M}_a^\text{moduli}} \text{ or } \frac{\tilde{M}_a^\text{conformal}}{\tilde{M}_a^\text{gauge}},$$

where $b_a = (33/5, 1, -3)$ are the MSSM beta function coefficients. This pattern is termed mirage pattern as $M_a$ are unified at the mirage messenger scale [13]:

$$M_{\text{mirage}} = M_{\text{GUT}}(m_{3/2}/M_{\text{Pl}})^{\alpha/2}. $$

Examples giving the mirage pattern of gaugino masses include the KKLT compactification [14] of Type IIB string theory with the MSSM gauge fields living on $D7$ branes [15], deflected anomaly mediation scenario proposed in [16], and also some variants of KKLT setup [17].

Mirage pattern might be considered as a smooth interpolation between the mSUGRA pattern ($\alpha = 0$) and the anomaly pattern ($\alpha = \infty$). For a positive $\alpha = \mathcal{O}(1)$ which is predicted to be the case in most of the schemes yielding the mirage pattern, gaugino masses are significantly more degenerate than those in mSUGRA and anomaly patterns. Different schemes giving the same mirage pattern of gaugino masses can be distinguished from each other by sfermion masses. For instance, in mirage mediation scheme [13, 15] resulting from KKLT-type string compactification [14], the 1st and 2nd generations of squark and slepton masses show up the same mirage unification at $M_{\text{mirage}}$, while the sfermion masses in deflected anomaly mediation scenario have a different structure [16].

**Anomaly pattern:** If there is no singlet with nonzero $F$-component or all SUSY breaking fields are sequestered from the visible gauge fields, $(M_a/g_a^2)_{\text{TeV}}$ are dominated by $\tilde{M}_a^\text{conformal}$. 


leading to

\[ M_1 : M_2 : M_3 \simeq 3.3 : 1 : 9 \]

which is termed anomaly pattern. One stringy example giving the anomaly pattern would be the KKLT compactification with the MSSM gauge fields on D3 branes [15].

We finally note that there are schemes in which \((M_\mu / g_\mu^2)_{\text{TeV}}\) receive an important contribution from the UV thresholds \(\tilde{M}_\mu^{\text{UV}}\) at string or GUT scale [18]. Gaugino masses in such scheme are the most model-dependent, and one needs to know the details of the model around the string or GUT scale in order to determine the low energy gaugino masses.

III. MEASURING GAUGINO MASSES AT LHC

For \(R\)-parity conserving SUSY model with neutralino LSP, if gluino or squarks are light enough to be copiously produced at LHC, some superparticle masses can be experimentally determined by analyzing the various invariant mass distributions for the decay products of the gluino or squark decays [3, 4]. The three gaugino mass patterns discussed in the previous section can be clearly distinguished by their prediction of the gluino to LSP neutralino mass ratio: \(m_{\tilde{g}} / m_{\chi_1} \gtrsim 6\) for mSUGRA pattern and \(m_{\tilde{g}} / m_{\chi_1} \gtrsim 9\), while \(m_{\tilde{g}} / m_{\chi_1}\) can be significantly smaller than 6 in mirage pattern. (Note that a nonzero Higgsino component in the LSP \(\chi_1^0\) makes the LSP mass \(m_{\chi_1}\) smaller than the smallest gaugino mass.)

There are several LHC observables providing information on gaugino masses, which are expected to be available under a mild assumption on SUSY spectra. Let us suppose that the gluino has a mass lighter than 2 TeV and \(\chi_1^0\) has a sizable Bino or Wino component, so that there will be a copious production of gluino pairs at LHC, subsequently decaying into four jets and two LSPs: \(\tilde{g} \tilde{g} \rightarrow q_1 q_2 \tilde{\chi}_1^0 q_3 q_4 \tilde{\chi}_1^0\). One observable useful for the determination of \(\{m_{\tilde{g}}, m_{\chi_1}\}\) is the \(M_{T2}\) variable [5, 19] of this gluino pair decay:

\[
M_{T2}^2(\tilde{g} \rightarrow qq\chi_1) \equiv \min_{p_{T1}^q, p_{T1}^{\tilde{q}} = p_{T2}^{\text{miss}}} \left[ \max\{m_T^2(p_T^{q_1}, p_T^{q_2}, p_T^{\tilde{q}_1}), m_T^2(p_T^{q_3}, p_T^{q_4}, p_T^{\tilde{\chi}_1})\} \right],
\]
where

\[ m_T^2(p_T^{q_1},p_T^{q_2},p_T^\chi) \equiv m_{\chi_1}^2 + 2(E_T^{q_1} E_T^\chi - p_T^{q_1} \cdot p_T^\chi) + 2(E_T^{q_2} E_T^\chi - p_T^{q_2} \cdot p_T^\chi) \]

(19)

for \( E_T = \sqrt{|p_T|^2 + m^2} \). Here, \( p_T^{q_1} \) and \( p_T^{q_2} \) denote the transverse momentum of the quark jets from the one gluino decay, \( p_T^{\text{miss}} \) is the observed missing transverse momentum, and we have ignored the light quark masses. If squark also has a mass comparable to the gluino mass, so that squark pairs can be copiously produced, the \( M_{T2} \) variable for the squark pair decay \( \tilde{q}\tilde{q} \rightarrow q_1\chi_1 q_2\chi_1 \) provides an information on \( \{m_{\tilde{q}},m_{\chi_1}\} \):

\[ M_{T2}^2(\tilde{g} \rightarrow q\chi_1) \equiv \min_{p_T^{\chi_1}+p_T^{\chi_2}=p_T^{\text{miss}}} \left[ \max\{m_T^2(p_T^{q_1},p_T^{\chi_1}),m_T^2(p_T^{q_2},p_T^{\chi_2})\} \right], \]

(20)

where

\[ m_T^2(p_T^{q_1},p_T^{\chi}) \equiv m_{\chi_1}^2 + 2(E_T^{q_1} E_T^\chi - p_T^{q_1} \cdot p_T^\chi). \]

(21)

The above two \( M_{T2} \) variables will be available at LHC as long as both the gluino and squark pairs are copiously produced at LHC, and a sizable fraction of them decay into the LSP pair plus quark jets. One still needs further information to determine the gluino to LSP mass ratio. Such additional information can be provided by measuring the maximal invariant mass \( M_{qq}^{\text{max}} \) of two jets in the final products of the gluino decay \( \tilde{g} \rightarrow \tilde{q}q \rightarrow \chi_1 qq \). The observed two \( M_{T2} \) variables \([18]\) and \([20]\) will tell us which of the gluino and squark is heavier than the other. Then, under an appropriate event selection cut, the measured \( M_{qq}^{\text{max}} \) corresponds to

\[ M_{qq}^{\text{max}} \simeq \left[ \frac{(m_{\tilde{g}}^2 - m_{\tilde{q}}^2)(m_{\tilde{q}}^2 - m_{\chi_1}^2)}{m_{\tilde{q}}^2} \right]^{1/2} \text{ for heavier gluino} \]

or \( m_{\tilde{g}} - m_{\chi_1} \) for heavier squark.

(22)

Combining this information with those from the two \( M_{T2} \) variables \([18]\) and \([20]\), one can determine \( \{m_{\tilde{g}},m_{\tilde{q}},m_{\chi_1}\} \) and thus the gluino to LSP mass ratio.

In regard to the mass of the second lightest neutralino \( \chi_2^0 \), a particularly interesting possibility is that \( \chi_2^0 \) is heavier than slepton, e.g. \( m_{\chi_2} - m_{\tilde{l}} \gtrsim 10 \text{ GeV} \) for which the lepton from \( \chi_2^0 \rightarrow \tilde{l}l \) is energetic enough to pass the selection cut, so the following cascade decay of squark is available \([3]\):

\[ \tilde{q} \rightarrow q\chi_2^0 \rightarrow q\tilde{l}^\pm l^\mp \rightarrow q\chi_1^0 l^+ l^- . \]

(23)
In such case, one can look at the edges of the invariant mass distributions of $ll$, $llq$ and $lq$ to determine $m_{\chi_2}$ and $m_{\tilde{t}}$. If $\chi_2^0$ is also gaugino-like, the obtained value of $m_{\chi_2}$ will allow the full determination of the gaugino mass ratios. Even when the slepton is heavier than $\chi_2^0$, so the decay $\chi_2 \rightarrow \tilde{l}l$ is not open, one can still determine $m_{\chi_2} - m_{\chi_1}$ using the dilepton invariant mass distribution in the 3-body decay $\chi_2 \rightarrow \chi_1 l^+ l^-$. 

Acknowledgments

This work is supported by the KRF Grant funded by the Korean Government (KRF-2005-201-C00006), the KOSEF Grant (KOSEF R01-2005-000-10404-0), and the Center for High Energy Physics of Kyungpook National University. We thank K. S. Jeong and H. P. Nilles for useful discussions.

[1] H. P. Nilles, Phys. Rept. 110, 1 (1984); H. E. Haber and G. L. Kane, Phys. Rept. 117, 75 (1985).
[2] K. Choi and H. P. Nilles, JHEP 0704, 006 (2007) [arXiv:hep-ph/0702146].
[3] G. Weiglein et al., Phys. Rept. 426, 47 (2006) [arXiv:hep-ph/0410364].
[4] W. S. Cho, Y. G. Kim, K. Y. Lee, C. B. Park and Y. Shimizu, JHEP 0704, 054 (2007) [arXiv:hep-ph/0703163].
[5] W. S. Cho, K. Choi, Y. G. Kim and C. B. Park, in preparation.
[6] V. Kaplunovsky and J. Louis, Nucl. Phys. B 422, 57 (1994) [arXiv:hep-th/9402005]; N. Arkani-Hamed, G. F. Giudice, M. A. Luty and R. Rattazzi, Phys. Rev. D 58, 115005 (1998) [arXiv:hep-ph/9803290].
[7] 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];
[8] V. S. Kaplunovsky and J. Louis, Phys. Lett. B 306, 269 (1993) [arXiv:hep-th/9303040]; A. Brignole, L. E. Ibanez and C. Munoz, Nucl. Phys. B 422, 125 (1994) [Erratum-ibid. B 436, 747 (1995)] [arXiv:hep-ph/9308271], H. P. Nilles, M. Olechowski and M. Yamaguchi, Phys. Lett. B 415, 24 (1997) [arXiv:hep-th/9707143]; K. Choi, H. B. Kim and C. Munoz, Phys. Rev. D 57, 7521 (1998) [arXiv:hep-th/9711158].
[9] For a review, see G. F. Giudice and R. Rattazzi, Phys. Rept. 322, 419 (1999) arXiv:hep-ph/9801271.

[10] K. Konishi, Phys. Lett. B135, 439 (1984).

[11] D. E. Kaplan, G. D. Kribs and M. Schmaltz, Phys. Rev. D 62, 035010 (2000) arXiv:hep-ph/9911293; Z. Chacko, M. A. Luty, A. E. Nelson and E. Ponton, JHEP 0001, 003 (2000) arXiv:hep-ph/9911323.

[12] V. Balasubramanian, P. Berglund, J. P. Conlon and F. Quevedo, JHEP 0503, 007 (2005) arXiv:hep-th/0502058; J. P. Conlon, F. Quevedo and K. Suruliz, JHEP 0508, 007 (2005) arXiv:hep-th/0505076.

[13] K. Choi, K. S. Jeong and K. i. Okumura, JHEP 0509, 039 (2005) arXiv:hep-ph/0504037.

[14] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, Phys. Rev. D 68, 046005 (2003) arXiv:hep-th/0301240.

[15] K. Choi, A. Falkowski, H. P. Nilles, M. Olechowski and S. Pokorski, JHEP 0411, 076 (2004) arXiv:hep-th/0411066; K. Choi, A. Falkowski, H. P. Nilles and M. Olechowski, Nucl. Phys. B718, 113 (2005) arXiv:hep-th/0503216.

[16] A. Pomarol and R. Rattazzi, JHEP 9905, 013 (1999) arXiv:hep-ph/9903448.

[17] O. Lebedev, H. P. Nilles and M. Ratz, Phys. Lett. B 636 (2006) 126 arXiv:hep-th/0603047; K. Choi and K. S. Jeong, JHEP 0608, 007 (2006) arXiv:hep-th/0605108; JHEP 0701, 103 (2007) [arXiv:hep-th/0611279].

[18] As recent examples of such scheme, see B. Acharya, K. Bobkov, G. Kane, P. Kumar and D. Vaman, Phys. Rev. Lett. 97, 191601 (2006) arXiv:hep-th/0606262; R. Dermisek, H. D. Kim and I.-W. Kim, JHEP 0610, 001 (2006) arXiv:hep-ph/0607169.

[19] C. G. Lester and D. J. Summers, Phys. Lett. B 463, 99 (1999).