Anomaly mediated SUSY breaking model retrofitted for naturalness

Howard Baer\textsuperscript{1}, Vernon Barger\textsuperscript{2} and Dibyashree Sengupta\textsuperscript{1}

\textsuperscript{1}Department of Physics and Astronomy, University of Oklahoma, Norman, OK 73019, USA
\textsuperscript{2}Department of Physics, University of Wisconsin, Madison, WI 53706, USA

Abstract

Anomaly-mediated supersymmetry breaking (AMSB) models seem to have become increasingly implausible due to 1. difficulty in generating a Higgs mass $m_h \sim 125$ GeV, 2. typically unnatural superparticle spectra characterized by a large superpotential mu term and 3. the possibility of a wino-like lightest SUSY particle (LSP) as dark matter now seems to be excluded. In the present paper we propose some minor modifications to the paradigm model which solve these three issues. Instead of adding a universal bulk scalar mass to avoid tachyonic sleptons, we add distinct Higgs and matter scalar soft masses which then allow for light higgsinos. To gain accord with the measured Higgs mass, we also include a bulk trilinear soft term. The ensuing natural generalized AMSB (nAMSB) model then has a set of light higgsinos with mass nearby the weak scale $m(W, Z, h) \sim 100$ GeV as required by naturalness while the winos populate the several hundred GeV range and gluinos and squarks occupy the multi-TeV range. For LHC searches, the wino pair production followed by decay to same-sign diboson signature channel offers excellent prospects for discovery at high luminosity LHC along with higgsino pair production leading to soft dileptons plus jet(s)+MET. A linear $e^+e^-$ collider operating above higgsino pair production threshold should be able to distinguish the AMSB gaugino spectra from unified or mirage unified scenarios. Dark matter is expected to occur as a higgsino-like WIMP plus axion admixture.
1 Introduction

The discovery of $D$-branes in superstring models in the 1990s\cite{1} ushered in new avenues for particle physics model building. In the case of supersymmetry (SUSY), this was exemplified initially with the advent of models where the dominant contribution to soft SUSY breaking Lagrangian parameters originated from violations of the superconformal anomaly, in what became known as anomaly-mediated SUSY breaking models, or AMSB\cite{2,3}. The AMSB contributions to soft SUSY breaking terms are always present in gravity mediation, but since they occur at loop level, they are usually suppressed compared to tree-level contributions and hence had previously been mostly neglected. Randall and Sundrum (RS) constructed an extra-dimensional scenario where the AMSB soft term contributions were expected to be the dominant or nearly dominant terms. The initial idea was that the visible sector, usually assumed to be the Minimal Supersymmetric Standard Model or MSSM, would be located on one three-brane extending through an assumed extra-dimensional spacetime, while SUSY breaking would occur on a different brane. Thus, the SUSY breaking sector was in fact sequestered, or separated from the visible sector brane within the extra-dimensional spacetime. This setup suppressed tree-level SUSY breaking soft terms in the visible sector. But since gravity propagates in the bulk, the entire extra-dimensional spacetime, the anomaly-mediated contributions could dominate the visible sector soft terms.

The AMSB gaugino masses were calculated to be proportional to the corresponding gauge group beta functions times the gravitino mass

$$M_i = \frac{\beta_i}{g_i} m_{3/2}$$

with $\beta_i = \frac{g_i^2}{16\pi^2} b_i$, $b_i = (6, 6, 1, -3)$ and $i$ labels the gauge group. Taking into account the running gauge coupling values at the weak scale, then one expects gaugino masses in the ratio $M_1 : M_2 : M_3 \sim 3.3 : 1 : -9$ so that the winos are the lightest of the weak scale gauginos. This is in contrast to models with unified gaugino masses where the bino occurs as the lightest gaugino. The lightest neutral wino was then typically assumed to be the lightest SUSY particle (LSP) in AMSB with striking consequences for collider and dark matter signatures\cite{4,5,6}.

In addition, in AMSB the soft breaking scalar masses were computed to be

$$m_f^2 = -\frac{1}{4} \left\{ \frac{d\gamma}{dg} \beta_g + \frac{d\gamma}{df} \beta_f \right\} m_{3/2}^2$$

where $\beta_f$ is the beta function for the corresponding superpotential Yukawa coupling and anomalous dimension $\gamma = \partial \ln Z/\partial \ln \mu$ with $Z$ the wave function renormalization constant and $\mu$ is the running energy scale. The AMSB contribution to trilinear soft SUSY breaking terms is given by

$$A_f = \frac{\beta_f}{f} m_{3/2}$$

where $f$ is the corresponding Yukawa coupling.

For some assumed value of gravitino mass $m_{3/2} \sim 50-100$ TeV, then all the AMSB soft terms are comparable to each other with values near to the weak scale as required by phenomenology.
An annoyance is that the slepton masses turn out to be tachyonic with negative mass-squared leading to an electric charge breaking minimum for the scalar potential. It was suggested by RS[2] that additional bulk contributions to scalar masses, which are comparable to the AMSB contributions, could be present to alleviate this problem. An assortment of other solutions to the negative slepton mass problem were also devised[7].

To gain concrete phenomenological predictions for AMSB at colliding beam and dark matter detection experiments, a minimal AMSB model (mAMSB) was devised wherein a common bulk contribution $m_0^2$ was appended to all AMSB scalar mass-squared values[5, 6]. Once the weak scale soft terms were determined, then the superpotential $\mu$ term was tuned so as to maintain the measured value of the $Z$ boson mass via the scalar potential minimization conditions. Thus, the parameter space of the mAMSB model was given by

$$m_0, \ m_{3/2}, \ tan\beta, \ sign(\mu).$$

Expectations for LHC searches within the mAMSB construct have been presented in Ref’s [8, 9]. Searches for direct chargino pair production in mAMSB with disappearing tracks from long-lived but ultimately unstable wino-like charginos[4] have been presented by Atlas[10].

The minimal AMSB model has provided a beautiful and compelling framework for new physics searches. It has been especially appreciated for containing solutions to the SUSY flavor problem (since the sfermions of each generation acquire common masses) and the gravitino problem (since gravitinos are so heavy that they decay much more quickly than the TeV-scale gravitinos which are expected in usual SUGRA models). While wino-like WIMPs are thermally underproduced in the mAMSB model, it was hypothesized by Moroi and Randall[11] that non-thermal WIMP production from, for instance, decay of light moduli fields could augment the relic abundance of dark matter and bring its mass abundance into accord with measured values.

While the mAMSB model is a well-motivated and beautiful construct, recently it has suffered several setbacks on the phenomenological front.

- The first of these was the discovery of the Higgs boson at a mass value $m_H \approx 125$ GeV. In the mAMSB model, the trilinear soft terms given by Eq. 3 are generally not large enough to lift the predicted value of $m_H$ into the 125 GeV range unless sparticle masses are very heavy – in the vicinity of tens of TeV[12, 13, 14]. Such heavy sparticle masses exacerbate the so-called Little Hierarchy problem which arises from the growing mass gap between the measured value of the weak scale and the sparticle mass scale.

- The second setback arises from non-observation of sparticles at the CERN Large Hadron Collider (LHC). While one solution to this issue is to simply posit that the mAMSB sparticles are heavier than experimental limits, this also makes the theory increasingly unnatural[18] and hence increasingly implausible.

- A third setback arose on the dark matter front. In mAMSB, where a wino-like WIMP is expected to comprise the dark matter, the model has come into conflict with new stringent limits from direct and indirect dark matter detection experiments. Searches for WIMPs at underground noble liquid experiments– which test the spin-independent (SI) direct detection (DD) rate– apparently exclude about half the remaining mAMSB
parameter space\textsuperscript{29}. Meanwhile, indirect WIMP detection (IDD) searches– via observation of gamma rays arising from WIMP-WIMP annihilation into hadrons followed by e.g. $\pi^0 \to \gamma\gamma$ decay– have placed severe limits on wino dark matter. The Fermi-LAT/MAGIC collaboration\textsuperscript{30}, via a search for gamma rays from dwarf spheroidal galaxies, now seems to require $m(\text{wino}) \gtrsim 700$ GeV. Along with this, the HESS experiment\textsuperscript{31}, from 254 hours (10 years) of observation of continuum gamma rays arising from the galactic center, now requires $m(\text{wino}) \gtrsim 1200$ GeV. If Sommerfeld enhancement effects are included in the WIMP-WIMP annihilation rate, then wino-like WIMPs seem to be excluded over their entire mass range\textsuperscript{32, 33, 29}. At first sight, such limits from IDD dark matter searches would seem to exclude models like mAMSB with wino-like WIMP dark matter\textsuperscript{4}.

To expand upon the fine-tuning/naturalness issue, we here adopt the most conservative fine-tuning measure, $\Delta_{EW}$\textsuperscript{15, 16}. The quantity $\Delta_{EW}$ measures how well the weak scale MSSM Lagrangian parameters match the measured value of the weak scale. By minimizing the MSSM weak scale scalar potential to determine the Higgs field vevs, one derives the well-known expression relating the $Z$-boson mass to the SUSY Lagrangian parameters:

$$m_Z^2/2 = m_{H_d}^2 + \Sigma_d^d - (m_{H_u}^2 + \Sigma_u^u) \tan^2 \beta \over \tan^2 \beta - 1 - \mu^2 \simeq -m_{H_u}^2 - \Sigma_u^u(t_{1,2}) - \mu^2. \quad (5)$$

Here, $\tan \beta = v_u/v_d$ is the ratio of Higgs field vacuum-expectation-values and the $\Sigma^u$ and $\Sigma^d$ contain an assortment of radiative corrections, the largest of which typically arise from the top squarks. Expressions for the $\Sigma^u$ and $\Sigma^d$ are given in the Appendix of Ref.\textsuperscript{16}. Thus, $\Delta_{EW}$ compares the maximal contribution on the right-hand-side (RHS) of Eq. (5) to the value of $m_Z^2/2$. If the RHS terms in Eq. (5) are individually comparable to $m_Z^2/2$, then no unnatural fine-tunings are required to generate $m_Z = 91.2$ GeV\textsuperscript{4}.

The main requirements for low electroweak fine-tuning ($\Delta_{EW} \lesssim 30$)\textsuperscript{1} are the following.

- $|\mu| \sim 100 - 300$ GeV\textsuperscript{23, 24} (the lighter the better) where $\mu \gtrsim 100$ GeV is required to accommodate LEP2 limits from chargino pair production searches.

- $m_{H_u}^2$ is driven radiatively to small– not large– negative values at the weak scale\textsuperscript{15, 16}.

- The top squark contributions to the radiative corrections $\Sigma^u(t_{1,2})$ are minimized for TeV-scale highly mixed top squarks\textsuperscript{15}. This latter condition also lifts the Higgs mass to $m_h \sim 125$ GeV. For $\Delta_{EW} \lesssim 30$, the lighter top squarks are bounded by $m_{\tilde{t}_1} \lesssim 3$ TeV\textsuperscript{16, 22}.

- The gluino mass, which feeds into the stop masses at one-loop and hence into the scalar potential at two-loop order, is bounded by $m_{\tilde{g}} \lesssim 6$ TeV\textsuperscript{16, 22}.

\textsuperscript{1} A possibility which avoids these constraints consists of mixed wino/axion dark matter\textsuperscript{34}.

\textsuperscript{2} Other measures include $\Delta_{BG} = \max \frac{p_i \delta m^2}{m_Z^2}$ where $p_i$ are fundamental parameters of the theory\textsuperscript{17}. In a theory where all soft terms are interdependent (such as AMSB or SUGRA or GMSB) then $\Delta_{BG}$ reduces to $\Delta_{EW}$\textsuperscript{18}. Sometimes $\Delta_{HS} = \delta m_h^2/m_h^2$ is used\textsuperscript{19} where $\delta m_h^2 \sim -\frac{3f_t^2}{8\pi^2} (m_{\tilde{Q}_3}^2 + m_{\tilde{L}_3}^2 + A_t^2) \ln (\Lambda^2/m^2_{\text{SUSY}})$ with $f_t$ the top Yukawa coupling, $\Lambda$ is as high as $m_{\text{GUT}}$ and $m_{\text{SUSY}} \sim 1$ TeV. This measure has been oversimplified by neglecting the $m_{H_u}^2$ contribution to its own running so as not to allow for radiatively driven naturalness, where large high scale soft terms are driven by radiative corrections to natural values at the weak scale\textsuperscript{20, 21}.

\textsuperscript{3} The onset of fine-tuning for $\Delta_{EW} \gtrsim 30$ is visually displayed in Ref.\textsuperscript{22}. \end{footnotesize}
Figure 1: Plot of points from a scan over mAMSB parameter space in the $\Delta_{EW}$ vs. $m_h$ plane.

In Fig. 1 we show the results of a scan over mAMSB model parameter space in the $\Delta_{EW}$ vs. $m_h$ plane. We use Isajet 7.87[25] to generate the mAMSB spectra. We have scanned over

- $m_0 : 1 - 10$ TeV,
- $m_{3/2} : 80 - 1000$ TeV,
- $\tan \beta : 4 - 58$,

with $\mu > 0$.

From Fig. 1 we see that the minimal value of $\Delta_{EW}$ occurs around 100 so that indeed the model is fine-tuned in the electroweak sector at least at the $\sim 1\%$ level. The lowest $\Delta_{EW}$ points occur at $m_{3/2} \sim 100$ TeV where $m_{\tilde{g}} \sim 2$ TeV, just beyond the current LHC $m_{\tilde{g}}$ mass limit[36]. While many of these points have $m_h \sim 122$ GeV, to gain $m_h \sim 125$ GeV the value of $\Delta_{EW}$ jumps to $\gtrsim 6000$.

To improve upon this situation, in this paper we present a retrofitted phenomenological AMSB model which is a generalization of mAMSB and which addresses the three issues discussed above. Indeed, in the original RS paper[2], the authors actually advocated for the modifications we present here. It was only when some simplifications were implemented in the original minimal AMSB model that these features were abandoned[5, 6]. The two generalizations to mAMSB include the following:

1. independent bulk contributions $m_{H_u}^{2}(bulk)$ and $m_{H_d}^{2}(bulk)$ to the soft SUSY breaking Higgs masses as opposed to matter scalar bulk masses $m_0^{2}(1, 2)$ (for first/second generation matter scalars) and $m_0^{2}(3)$ (for third generation matter scalars) and
2. inclusion of bulk contributions $A_0$ to the trilinear soft terms.

These two modest changes in the AMSB model will allow each of the three issues above to be circumvented. However, we will also see that the anticipated collider phenomenology and dark matter expectations will be very different. After bringing the model into accord with the measured Higgs mass and naturalness, the LSP will no longer be a wino-like neutralino, but instead a higgsino-like neutralino. If we posit that the SUSY $\mu$ problem is solved via the Kim-Nilles mechanism\cite{37} (a supersymmetrized version of the DFSZ axion model\cite{38} which allows for $\mu \ll m_{soft}$) then dark matter is expected to consist of an axion plus higgsino-like WIMP admixture\cite{39}.

In Sec. 2 we make explicit our modified AMSB soft term formulae. We also present aspects of the anticipated natural AMSB spectra where now the LSP is expected to be a higgsino-like neutralino but where the lightest gaugino is still expected to be wino-like. Since the model can now be rendered natural, we dub the resultant model as nAMSB, or natural anomaly-mediation, to distinguish it from the previously explored minimal AMSB model. We present some benchmark spectra and a nAMSB model line. In Sec. 3 we discuss consequences of the nAMSB model for collider and dark matter searches. In Sec. 4 we summarize and present our conclusions.

2 Natural Anomaly Mediated SUSY Breaking Model (nAMSB)

2.1 Soft terms for nAMSB

In this Section, we propose several minor modifications of the mAMSB model which will allow for naturalness along with a Higgs mass $m_h \simeq 125$ GeV.

For gaugino masses, we maintain the usual formulae:

\begin{align}
M_1 &= \frac{33}{5} \frac{g_1^2}{16\pi^2} m_{3/2}, \\
M_2 &= \frac{g_2^2}{16\pi^2} m_{3/2}, \\
M_3 &= -3 \frac{g_3^2}{16\pi^2} m_{3/2}.
\end{align}

\[\text{(6)-(8)}\]
Third generation soft SUSY breaking scalar squared masses are given by

\[
m^2_{\tau_3} = \left(-\frac{88}{25} g_1^4 + 8 g_3^4 + 2 f_t \beta_f \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2(3),
\]

(9)

\[
m^2_{D_3} = \left(-\frac{22}{25} g_1^4 + 8 g_3^4 + 2 f_b \beta_f \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2(3),
\]

(10)

\[
m^2_{Q_3} = \left(-\frac{11}{50} g_1^4 - \frac{3}{2} g_2^4 + 8 g_3^4 + f_t \beta_{f_t} + f_b \beta_{f_b} \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2(3),
\]

(11)

\[
m^2_{L_3} = \left(-\frac{99}{50} g_1^4 - \frac{3}{2} g_2^4 + f_{t} \beta_{f_t} \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2(3),
\]

(12)

\[
m^2_{E_3} = \left(-\frac{198}{25} g_1^4 + 2 f_{e} \beta_{f_e} \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_0^2(3),
\]

(13)

while first/second generation scalar squared masses are given by similar formulae but where the associated Yukawa couplings may be safely ignored and the bulk sfermion mass is changed from \(m_0^2(3) \rightarrow m_0^2(1,2)\).

For soft SUSY breaking Higgs masses, we propose (in accord with Ref. [2]) that each Higgs doublet receive an independent bulk mass contribution so that

\[
m^2_{H_u} = \left(-\frac{99}{50} g_1^4 - \frac{3}{2} g_2^4 + 3 f_t \beta_{f_t} \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_{H_u}^2(bulk),
\]

(14)

\[
m^2_{H_d} = \left(-\frac{99}{50} g_1^4 - \frac{3}{2} g_2^4 + 3 f_b \beta_{f_b} + f_{e} \beta_{f_e} \right) \frac{m_{3/2}^2}{(16\pi^2)^2} + m_{H_d}^2(bulk).
\]

(15)

The freedom of independent bulk Higgs soft masses \(m_{H_u}^2(bulk)\) and \(m_{H_d}^2(bulk)\) may be traded using the electroweak minimation conditions for the alternative weak scale inputs \(\mu\) and \(m_A\) (as in the NUHM2 SUSY model[26]).

Using this flexibility, we again scan over AMSB parameters as in Sec. [1] but now also including

- \(\mu : 100 - 500\) GeV and
- \(m_A : 0.25 - 10\) TeV.

The results are plotted again in the \(\Delta_{EW} vs. m_h\) plane and shown in Fig. [2]. From the figure, we see that now many points have dropped into the natural area where \(\Delta_{EW} < 30\). However, almost all these points also have \(m_h \lesssim 122\) GeV.

Thus, following Ref. [2], we propose adding as well a bulk contribution to the trilinear soft terms. Then the \(A\)-parameters are given by

\[
A_t = \frac{\beta_{f_t} m_{3/2}}{f_t 16\pi^2} + A_0,
\]

(16)

\[
A_b = \frac{\beta_{f_b} m_{3/2}}{f_b 16\pi^2} + A_0, \text{ and}
\]

(17)

\[
A_\tau = \frac{\beta_{f_\tau} m_{3/2}}{f_\tau 16\pi^2} + A_0.
\]

(18)
Figure 2: Plot of points in the $\Delta_{EW}$ vs. $m_h$ plane from a scan over AMSB parameter space with added bulk Higgs soft terms but without bulk $A_0$ terms.

The quantities $\hat{\beta}_{f_i}$ that enter the expressions for scalar masses and $A$-parameters are given by the standard expressions

\begin{align}
\hat{\beta}_{f_t} &= 16\pi^2 \beta_t = f_t \left( -\frac{13}{15} g_1^2 - 3g_2^2 - \frac{16}{3} g_3^2 + 6f_t^2 + f_b^2 \right), \\
\hat{\beta}_{f_b} &= 16\pi^2 \beta_b = f_b \left( -\frac{7}{15} g_1^2 - 3g_2^2 - \frac{16}{3} g_3^2 + f_t^2 + 6f_b^2 + f_f^2 \right), \\
\hat{\beta}_{f_\tau} &= 16\pi^2 \beta_\tau = f_\tau \left( -\frac{9}{5} g_1^2 - 3g_2^2 + 3f_b^2 + 4f_f^2 \right).
\end{align}

The first two generations of squark and slepton masses are given by the corresponding formulae above with the Yukawa couplings set to zero. Eq. (6)-(18) serve as RGE boundary conditions at $Q = m_{GUT}$. The nAMSB model is therefore characterized by the parameter set,

\begin{equation}
\begin{align*}
m_0(1,2), \ m_0(3), \ m_{3/2}, \ A_0, \ \tan \beta, \ \mu, \ \text{and} \ m_A.
\end{align*}
\end{equation}

To see the effect of including the bulk $A_0$ trilinear soft term, we adopt a nAMSB benchmark point with parameters $m_{3/2} = 135$ TeV, $m_0(1,2) = 13$ TeV, $m_0(3) = 5$ TeV, $\mu = 200$ GeV and $m_A = 2$ TeV with $\tan \beta = 10$. In Fig. 3(a), we show the value of $\Delta_{EW}$ as we vary $A_0$. For no bulk trilinear, with $A_0 = 0$, then $\Delta_{EW} \sim 70$ and the model requires EW fine-tuning at the 1.4% level. As $A_0$ varies and becomes large positive or negative, then large mixing in the stop sector leads to a reduction in both $\Sigma_u^{\alpha}(\tilde{t}_{1,2})$ values. For $A_0 \sim +5$ TeV, then $\Delta_{EW}$ drops to as low as
Figure 3: Frame a): $\Delta_{EW}$ vs. $A_0$ for $m_{3/2} = 135$ TeV, $m_0(1, 2) = 13.5$ TeV, $m_0(3) = 5$ TeV, $\mu = 200$ GeV and $m_A = 2000$ GeV. In frame b), we plot $m_h$ vs. $A_0$ for the same parameters.

10. In frame b), we show the variation in $m_h$ versus $A_0$. With no bulk contribution to $A$ terms, then $m_h \sim 120$ GeV. As $A_0$ increases to $\sim +5$ TeV, then the added stop mixing increases $m_h$ until it reaches the $\sim 125$ GeV level.

In Fig. 4 we show the nAMSB spectra plot from our benchmark point where now we adopt $A_0 = +5.4$ TeV. From the plot, we see that the $W$, $Z$ and $h$ are clustered around the $\sim 100$ GeV scale with the higgsinos $\tilde{W}_1^\pm$ and $\tilde{Z}_{1,2}$ clustered not too far away at $\sim 200$ GeV as required by naturalness. Meanwhile, first/second generation matter sfermions lie in the multi-TeV range at $\sim 13$ TeV. For the gauginos, we have $m_{\tilde{g}} \sim 3$ TeV, well beyond current LHC limits which at present require $m_{\tilde{g}} \gtrsim 2$ TeV. What is characteristic about nAMSB is the rather light winos $\tilde{W}_2^\pm$ and $\tilde{Z}_3$ with mass $\sim 400$ GeV. The bino $\tilde{Z}_4$ has mass $\sim 1.2$ TeV. For the top squarks, we find them to be highly mixed by the large $A_t$ term with $m_{\tilde{t}_1} \sim 1.4$ TeV and $m_{\tilde{t}_2} \sim 3.5$ TeV. As we shall see, the nAMSB mass spectrum leads to very different expectations for LHC signatures as compared to the old mAMSB model.

The precise benchmark point mass values are listed numerically in Table 1 along with various calculated dark matter and $B$-decay observables. For this point, the thermal WIMP abundance of higgsino-like WIMP comes in (from IsaReD [28]) at $\Omega_{\tilde{Z}_1}^{TP} h^2 \sim 0.009$, a factor 13.3 below the measured abundance. In the case of nAMSB, we also expect the presence of a SUSY-DFSZ axion which would likely make up the remaining dark matter abundance. A complete calculation requires an eight-coupled Boltzmann equation computation [58]. The WIMP detection rates are also given, but in this case they must be scaled down by factors of $\xi \equiv \Omega_{\tilde{Z}_1} h^2/0.12$ for the SI and SD direct detection rates. For the IDD detection rate, which is similar to that expected by Dine from the intermediate branch of the IIB string theory landscape [27].
| parameter | nAMSB1 |
|-----------|--------|
| $m_{3/2}$ | 135000 |
| $\tan \beta$ | 10 |
| $m_0(1, 2)$ | 13000 |
| $m_0(3)$ | 5000 |
| $A_0$ | 5400 |
| $\mu$ | 200 |
| $m_A$ | 2000 |
| $m_{\tilde{g}}$ | 3037.6 |
| $m_{\tilde{g}_L}$ | 13189.1 |
| $m_{\tilde{g}_R}$ | 13280.1 |
| $m_{\tilde{e}_R}$ | 12909.7 |
| $m_{\tilde{\tau}_1}$ | 1380.8 |
| $m_{\tilde{\tau}_2}$ | 3536.0 |
| $m_{\tilde{b}_1}$ | 3569.5 |
| $m_{\tilde{b}_2}$ | 5085.0 |
| $m_{\tilde{\tau}_1}$ | 4670.6 |
| $m_{\tilde{\tau}_2}$ | 4930.8 |
| $m_{\tilde{\nu}_r}$ | 4903.1 |
| $m_{\tilde{W}_2}$ | 398.0 |
| $m_{\tilde{W}_1}$ | 195.1 |
| $m_{\tilde{Z}_4}$ | 1225.6 |
| $m_{\tilde{Z}_3}$ | 405.9 |
| $m_{\tilde{Z}_2}$ | 209.7 |
| $m_{\tilde{Z}_1}$ | 183.9 |
| $m_h$ | 125.1 |
| $\Omega_{\tilde{Z}^0_1} h^2$ | 0.009 |
| $BF(b \to s\gamma) \times 10^4$ | 3.2 |
| $BF(B_s \to \mu^+\mu^-) \times 10^9$ | 3.8 |
| $\sigma^{SI}(\tilde{Z}_{1, p})$ (pb) | $9.8 \times 10^{-9}$ |
| $\sigma^{SD}(\bar{\tilde{Z}}_{1,p})$ (pb) | $2.4 \times 10^{-4}$ |
| $\langle \sigma v \rangle|_{v \to 0}$ (cm$^3$/sec) | $2.8 \times 10^{-25}$ |
| $\Delta_{EW}$ | 10.2 |

Table 1: Input parameters and masses in GeV units for a natural generalized anomaly mediation SUSY benchmark point with $m_t = 173.2$ GeV.
In Fig. 5, we repeat the above AMSB parameter space scans except now we include as well a scan over

- $A_0 : -20 \rightarrow +20$ TeV.

From the figure, we now see data points in accord with LHC sparticle mass constraints which populate the $\Delta_{EW} < 30$ naturalness regime whilst also allowing for $m_h \sim 125 \pm 3$ GeV. Thus, the combination of independent bulk Higgs masses and an added bulk trilinear soft term $A_0$ allows us to bring the AMSB model into accord with LHC Higgs mass measurements and naturalness requirements and dark matter constraints.

### 2.2 A nAMSB model line

In phenomenological studies of models for new physics, it is frequently useful to adopt model lines wherein new particle masses increase in a controlled manner thus allowing for collider reach calculations, decoupling, etc. We may elevate our previous benchmark model to a model line by allowing the gravitino mass to float so all sparticle masses increase with $m_{3/2}$ from the LHC limits until they become unnatural or decouple.

In Fig. 6, we show three frames resulting from a nAMSB model line versus $m_{3/2}$ starting at $m_{3/2} \approx 80$ TeV. This latter value corresponds to $m_{\tilde{g}} \sim 2$ TeV, just beyond the current LHC limits from simplified models. In frame a), we show how $\Delta_{EW}$ varies. At lower values $m_{3/2} \sim 100 - 150$ TeV, then $\Delta_{EW} \sim 10$ and the model is highly natural. As $m_{3/2}$ increases,
all soft terms increase according to Equations 6-18. As $m_{3/2}$ increases to the vicinity of 250 TeV, then $\Delta_{EW}$ has moved beyond the 30 value where fine-tuning begins to be required in the weak scale scalar potential. Thus, the regime where $m_{3/2} \lesssim 250$ TeV seems favored from a naturalness perspective. In frame b), we show the corresponding value of $m_h$ along the nAMSB model line. Its value begins at $m_h \sim 124$ GeV for $m_{3/2} \sim 80$ TeV and increases to $\sim 127$ GeV for $m_{3/2}$ as high as 370 TeV. Thus, the light Higgs mass stays within its required range (allowing for $\sim \pm 2$ GeV theory error in our $m_h$ calculation) over the entire model line. In frame c), we show various sparticle masses along the model line. The higgsinos $\tilde{W}_1^\pm$ and $\tilde{Z}_{1,2}$ remain clustered at $\sim 200$ GeV since the $\mu$ parameter remains fixed. The gluinos and stops lie in the several TeV range and as their masses increase, so too do the radiative corrections $\Sigma_u(\tilde{t}_{1,2})$ in Eq. 3. Over the range of $m_{3/2}$ consistent with naturalness, $m_{\tilde{g}}$ varies from 2 – 4 TeV while the lighter stop ranges from $m_{\tilde{t}_1} \sim 1.3 – 1.5$ TeV. Of considerable interest for collider searches is the range of the wino masses $m_{\tilde{W}_2}^\pm$ and $m_{\tilde{Z}_3}$. These vary from 300 GeV for $m_{3/2} \sim 100$ TeV to $\sim 600$ GeV for $m_{3/2} \sim 250$ GeV. This will have important ramifications for discussion of collider searches in the next Section.

### 2.3 Locus of natural AMSB parameters

It is important to check from scans over the full generalized AMSB parameter space in Eq. 22 where exactly the natural solutions with low $\Delta_{EW}$ exist. Thus, here we implement a scan over the full parameter space and plot each parameter versus $\Delta_{EW}$. To aid the reader, we show the demarcation where $\Delta_{EW}$ exceeds 30, although it is simple to extract parameter locales for other
Figure 6: Plot of a) $\Delta_{EW}$, b) $m_h$ and c) various sparticle masses versus $m_{3/2}$ along a nAMSB model line with $m_0(1,2) = 13$ TeV, $m_0(3) = 5$ TeV, $A_0 = 5.4$ TeV, $m_A = 2$ TeV and $\mu = 200$ GeV with $\tan \beta = 10$. 

\[ \Delta_{EW} \]

\[ m_h \text{ [GeV]} \]

\[ m_{3/2} \text{ [TeV]} \]
choices of a maximal $\Delta_{\text{EW}}$ value.

In Fig. 7(a), we show $\Delta_{\text{EW}}$ versus $m_{3/2}$ from our scan. Our points are extracted from the general scan with limits given above and also from a dedicated scan over parameters where $\Delta_{\text{EW}}$ is more likely to be $\lesssim 30$: $m_{3/2} \sim 80 - 300$ GeV, $\mu : 100 - 350$ GeV and $A_0 : 0.5m_0(3) - 2m_0(3)$. All points have $122 \text{ GeV} < m_h < 128 \text{ GeV}$. From frame a), we see that to maintain naturalness, $m_{3/2}$ is roughly bounded from above by about 300 GeV (in accord with the above nAMSB model line).

In frame b), we show $\Delta_{\text{EW}}$ versus $m_0(1, 2)$. The first and second generation scalar masses enter the naturalness measure via electroweak $D$-term contributions and these terms tend to cancel for nearly degenerate matter scalars. Thus, a wide range of $m_0(1, 2)$ values extending up into the 10-20 TeV range are allowed by naturalness. Such large first/second generation matter scalar masses allow for at least a partial decoupling solution to the SUSY flavor and CP problem (which may re-arise with the addition of flavor dependent bulk soft terms).

In Fig. 7(b) we show $\Delta_{\text{EW}}$ vs. $m_0(3)$. In this case, an upper bound of $m_0(3) \lesssim 8 \text{ TeV}$ emerges. This is because for too large values of third generation matter scalars, then the $\Sigma_u^u(\tilde{t}_{1, 2})$ contributions become large thus requiring some electroweak fine-tuning.

In Fig. 8(a), we show $\Delta_{\text{EW}}$ vs. $A_0/m_0(3)$. Here we see that for $A_0 \sim 0$, then $\Delta_{\text{EW}}$ is always $\gtrsim 30$ and unnatural. For $A_0/m_0(3) \sim -2$, then $\Delta_{\text{EW}}$ drops below 30. This occurs even more sharply for $A_0/m_0(3) \sim +1$. As noted previously, the large $A_0$ values decrease the $\Sigma_u^u(\tilde{t}_{1, 2})$ contributions whilst lifting $m_h \sim 125 \text{ GeV}$[15]. Frame b) shows $\mu$ vs. $\Delta_{\text{EW}}$. Here we see a sharp demarcation for naturalness when $\mu \lesssim 350$ GeV, the lighter the better. This is also seen from direct computation from Eq. 5. In frame c), we show variation versus tan$\beta$. In this case, a wide range of tan$\beta$ is allowed by naturalness, but not the very highest values where tan$\beta \gtrsim 40$. For such high tan$\beta$, then the $\Sigma_u^u(\tilde{t}_{1, 2})$ may become large, thus requiring some fine-tuning. In frame d), we show variation with $m_A$. In this case, for $m_A \gg m_Z$, then $m_{H_u} \sim m_A$ and naturalness in Eq. 5 would require $m_A/\tan \beta \lesssim \sqrt{30(m_Z^2/2)}$. This requires $m_A$ to be bounded from above by about 7-8 TeV.

### 2.4 Bounds on sparticle masses in the natural AMSB model

It is desirable in any SUSY model to extract upper bounds on various sparticle masses from naturalness in order to establish a testability criterion for the model. Thus, in this Section, we implement the full scan over nAMSB parameter space (as delineated above).

In Fig. 9 we show $\Delta_{\text{EW}}$ versus $m_{\tilde{g}}$. Here, we see that $m_{\tilde{g}}$ ranges from the LHC lower limit of $\sim 2 \text{ TeV}$ up to $m_{\tilde{g}} \sim 6 \text{ TeV}$ before the model becomes unnatural (where $\Delta_{\text{EW}}$ exceeds $\sim 30$). The expected range in $m_{\tilde{g}}$ will of course have important implications for gluino searches at present and planned hadron colliders. The upper bound $m_{\tilde{g}} \lesssim 6 \text{ TeV}$ is in accord with other SUSY models: gravity mediation in NUHM2[16, 22, 42] and in mirage mediation[15]. The reason is that $m_{\tilde{g}}$ feeds into the RG evolution of top squark soft terms and a larger value of $m_{\tilde{g}}$ therefore increases the $\Sigma_u^u(\tilde{t}_{1, 2})$ values.

In Fig. 10 we show the expected range for top squark masses $m_{\tilde{t}_1}$ (frame a)) and $m_{\tilde{t}_2}$ (frame b)). In frame a), we see that $m_{\tilde{t}_1}$ ranges from its approximate LHC lower bound of $m_{\tilde{t}_1} \gtrsim 1 \text{ TeV}$ up to at most 3 TeV before the nAMSB model becomes unnatural. Meanwhile, from frame b),
Figure 7: Plot of nAMSB parameter scan in the $\Delta_{EW}$ vs. $a)$ $m_{3/2}$, $b)$ $m_0(1, 2)$ and $m_0(3)$ planes. The greater density of points for $m_{3/2} \lesssim 300$ TeV comes from the narrow scan added to the broad scan.
Figure 8: Plot of nAMSB parameter scan in the $\Delta_{\text{EW}}$ vs. a) $A_0/m_0(3)$, b) $\mu$, c) $\tan \beta$ and d) $m_A$ planes.
we see that $m_{\tilde{t}_2}$ can range up to $\sim 6$ TeV.

In Fig. 11, we plot the expected range of wino mass $m_{\tilde{W}_2}$. In this case, $m_{\tilde{W}_2}$ (which is $\simeq m_{\tilde{Z}_1}$) ranges from a lower bound $\sim 250$ GeV to an upper bound from naturalness of $m_{\tilde{W}_2} \sim 800$ GeV. In AMSB models, the weak scale wino mass is typically $m(\text{wino}) \sim m_{\tilde{g}}/8$ so that the wino mass upper bound arises due to the $m_{\tilde{g}}$ limits arising from $\Sigma_u^c(\tilde{t}_{1,2})$. The wino mass range will also have important consequences for collider signatures for nAMSB.

Lastly, we plot the phenomenologically important mass gap $m_{\tilde{Z}_2} - m_{\tilde{Z}_1}$ versus $\Delta_{EW}$ in Fig. 12. This mass gap enters higgsino pair production signatures at both LHC and at linear $e^+ e^-$ colliders. Due to the proximity of the winos to the higgsinos, the mass gap is expected to be larger than in models with unified gaugino masses. Indeed, from the figure we see that $m_{\tilde{Z}_2} - m_{\tilde{Z}_1}$ ranges from about 10 GeV all the way up to 100 GeV. This may be compared to models with gaugino mass unification where instead $m_{\tilde{Z}_2} - m_{\tilde{Z}_1}$ ranges from $\sim 10 - 25$ GeV typically\[16].

3 Consequences for collider and dark matter searches

One of many intriguing aspects of the mAMSB model is that it led to rather unique collider signatures—such as the presence of quasi-stable winos in sparticle cascade decays. In this Section, we will find very different collider signatures for the nAMSB model.
Figure 10: Plot of nAMSB parameter scan in the $\Delta_{EW}$ vs. $a) \ m_{\tilde{t}_1}$ and $b) \ m_{\tilde{t}_2}$ planes.
Figure 11: Plot of nAMSB parameter scan in the $\Delta_{EW}$ vs. $m_{\tilde{W}_2}$ plane.

Figure 12: Plot of nAMSB parameter scan in the $\Delta_{EW}$ vs. $m_{\tilde{Z}_2} - m_{\tilde{Z}_1}$ plane.
3.1 LHC

3.1.1 Gluino pair production

At the CERN LHC, an important SUSY search channel comes from gluino pair production. In nAMSB, almost always $m_{\tilde{g}} > m_{\tilde{t}_1}$ so that $\tilde{g} \rightarrow \tilde{t}_1\tilde{t}_1^*$ followed by $\tilde{t}_1 \rightarrow t\tilde{Z}_{1,2,3}$ or $b\tilde{W}_{1,2}$. Thus, gluino pair production events are expected to be rich in both $t$ and $b$ quarks arising from gluino cascade decays. Recently, the reach for various LHC luminosity upgrades has been estimated for natural SUSY models. It is found in Ref. [43] that HL-LHC with $\sim 3 \text{ ab}^{-1}$ of integrated luminosity has a $5\sigma$ reach in $m_{\tilde{g}}$ to about $m_{\tilde{g}} \sim 2.8 \text{ TeV}$. From Fig. 9, we see that this covers only a small portion of nAMSB parameter space.

Meanwhile, the reach of HE-LHC has also been estimated. Using $\sqrt{s} = 33 \text{ TeV}$ and $1 \text{ ab}^{-1}$ integrated luminosity, it is found that HE-LHC reach extends to about $5 \text{ TeV}$[44, 45], thus covering essentially all of nAMSB parameter space. Updated run parameters for HE-LHC have recently been proposed as $\sqrt{s} = 27 \text{ TeV}$ but $L = 10 - 15 \text{ ab}^{-1}$. The HE-LHC reach using the lower energy/higher luminosity parameter is likely comparable to our quoted numbers.

3.1.2 Top squark pair production

Top squark pair production $pp \rightarrow \tilde{t}_1\tilde{t}_1^*$ is another important LHC search channel. In nAMSB, we expect $m_{\tilde{t}_1} : 1 - 3 \text{ TeV}$. This is to be compared to the $5\sigma$ HL-LHC reach to $m_{\tilde{t}_1} \sim 1.2 \text{ TeV}$[46]. Thus, in this channel again HL-LHC will be able to cover only a small portion of mAMSB parameter space. The $5\sigma$ HE-LHC reach extends to $m_{\tilde{t}_1} \sim 3.2 \text{ TeV}$[45]. Thus, HE-LHC should be able to cover essentially all mAMSB parameter space via top squark pair searches.

3.1.3 Higgsino pair production

From Eq. 5, we find that for $\Delta_{EW} < 30$, then $m_{\tilde{Z}_{1,2}}, m_{\tilde{W}_3} \sim \mu \lesssim 350 \text{ GeV}$. Thus, higgsino pair production reactions occur at potentially observable rates[47] at LHC. Typically, most of the energy from higgsino pair production goes into making up the two $\tilde{Z}_1$ particle’s rest mass, so the visible energy release is small, making higgsino pair production reactions challenging to see[24, 48]. A way forward has been proposed in References [49] where one produces $\tilde{Z}_1\tilde{Z}_2$ in association with hard initial state jet radiation. Then one may trigger on the hard jet (or $E_T$) and within such events search for low mass, soft opposite-sign dileptons arising from $\tilde{Z}_2 \rightarrow \tilde{Z}_1\ell^+\ell^-$ decay. Recent search results from CMS have been presented[50] and results from Atlas are imminent[51]. With HL-LHC, this channel may well be able to explore the entire parameter space. A distinctive feature of the nAMSB model is that the $\tilde{Z}_2 - \tilde{Z}_1$ mass gap is expected to be substantially larger than in models with gaugino mass unification or in mirage mediation[52] due to the smaller higgsino-wino mass gap.

3.1.4 Wino pair production

In SUSY models with light higgsinos, a compelling new signature has emerged[53]: wino pair production followed by decay to same-sign dibosons (SSdB): $pp \rightarrow \tilde{W}_2^\pm \tilde{Z}_3$ with $\tilde{W}_2 \rightarrow W\tilde{Z}_{1,2}$ and $\tilde{Z}_3 \rightarrow W^\pm\tilde{W}_1^\mp$. The higgsinos at the end of the decay chain are again quasi-visible so one
really expects half the time a $W^+W^± + E_T$ signal which has very low SM backgrounds arising mainly from $t\bar{t}W$ and other processes. Signal and background have been estimated in Ref’s \[53, 47, 54\]. It is found that the reach of HL-LHC extends to about $m_{\tilde{W}} \lesssim 1$ TeV. Thus, in this channel as well we expect HL-LHC to completely cover the nAMSB parameter space. If such a signal doesn’t emerge at HL-LHC, then the mAMSB model will be ruled out. If a signal does emerge, then in Ref. \[54\] several suggestions have been proposed to extract a measurement of the wino masses: via counting, via distributions and via $++$ to $−−$ charge asymmetry. The importance of this channel for the nAMSB model derives from the expected weak scale gaugino mass ratio in AMSB models $M_1 : M_2 : M_3 \sim 0.4 : 0.13 : 1$ where winos are expected to be far lighter than gluinos (or binos).

### 3.2 Linear electron-positron colliders

Since (simple) natural SUSY models require the presence of light higgsinos (via Eq. 5 and Fig. 8), then the proposed International Linear $e^+e^−$ Collider, or ILC, is expected to become a higgsino factory for $\sqrt{s} > 2m(higgsino)$ \[55\]. The main production reactions are $e^+e^- \to \tilde{W}_1^±\tilde{W}_1^−$ and $\tilde{Z}_1\tilde{Z}_2$. In spite of the low energy release expected from these reactions, the clean operating environment and low SM backgrounds should allow the higgsino pair production events to be easily visible. These features, along with kinematic restrictions on the events, should allow for precision mass measurements of $\tilde{W}_1^+$, $\tilde{W}_1^−$, $\tilde{Z}_1$, and $\tilde{Z}_2$. If ILC is built with extendable energy ranging up to $\sqrt{s} \sim 1$ TeV, then there is a strong chance that direct wino production could also be detected via the $e^+e^- \to \tilde{W}_1^±\tilde{W}_2^∓$ channel in nAMSB.

Since the mass gaps $m_{\tilde{W}_1} - m_{\tilde{Z}_1}$ and $m_{\tilde{Z}_2} - m_{\tilde{Z}_1}$ depend sensitively on the higgsino-gaugino mixing, it has been shown in Ref’s \[55, 56, 57\] that the electroweak gaugino masses can also be extracted to percent level accuracy. Once the EW gaugino masses are known to sufficient accuracy, then they may be run via RGE’s to higher energies to test whether or not they unify. In the case of nAMSB, where $M_2 \ll M_1$ is expected, the ILC would be able to quickly show that anomaly-mediation is the likely underlying SUSY model.

### 3.3 Dark matter: WIMPs and axions

Dark matter in nAMSB is expected to be a higgsino-like WIMP plus SUSY DFSZ axion admixture, as with other natural SUSY models. As seen in Table 1, the $\tilde{Z}_1$ are thermally underproduced in the early universe although non-thermal processes such as axino and/or saxion production and decay in the early universe may augment these rates. The remaining abundance is expected to be comprised of axions. In the case where thermal WIMP production dominates, then indeed the bulk of dark matter would be axions. Precise estimates of the dark matter abundance require the solution of eight coupled Boltzmann equations which track the radiation density and number densities for WIMPs, axions, axinos, thermal- and coherent oscillation-production of saxions and gravitino production \[58\].

To assess WIMP detection prospects, one must account for the diminished abundance of WIMPs that is quantified by $\xi \equiv \Omega_{\tilde{Z}} h^2/0.12$ and where in Table 1 we would expect $\xi$ as low as 0.075. The spin-independent neutralino-proton scattering cross section from Isatools is shown in Table 1. Mutiplying by $\xi$ and comparing to recent exclusion limits, it is found the benchmark
point to be slightly excluded by recent LUX limits. But for the case of the nAMSB model, we expect typically higher $\sigma_{SI}(\tilde{Z}_1p)$ rates because the WIMP-WIMP-$h$ coupling, which enters the SI detection rate, is a product of gaugino times higgsino component. The typically reduced wino mass in nAMSB (as compared to models with gaugino mass unification) raises up the scattering rate somewhat. Detailed WIMP scattering calculations in this model will be needed for a complete assessment of detectability. The $\sigma_{SD}$ rate and indirect detection rate (in terms of $\langle \sigma v \rangle|_{v \to 0}$) are also given. Multiplying by $\xi$ and $\xi^2$ respectively, these two rates are still below current bounds as shown in Ref. [29].

While our benchmark point is nominally excluded, even with inclusion of the $\xi$ factor, we remark that further entropy dumping in the early universe could possibly lower the WIMP abundance even further from its thermal value [40]. A perhaps more compelling scenario is that the nAMSB model may provide a viable niche for light axino dark matter. In usual gravity-mediation, the axino (and saxion) are expected to gain masses of order $\sim m_3/2$ [59, 60]. In nAMSB, we would expect the saxion to gain a bulk soft mass $m_s \sim 1 \text{ TeV}$ but the axino mass could be suppressed leading to an unstable lightest neutralino which suffers late decay to e.g. $\tilde{a} + \gamma, Z, h$. In such a case, dark matter would be an axion/axino admixture.

Meanwhile, detection of the SUSY DFSZ axion has been shown to be more difficult than in the non-SUSY models due to the circulation of higgsinos in the $a - \gamma - \gamma$ triangle coupling [61]. Thus, we do not expect detection of the associated axion any time soon unless the presence of exotic matter in the $a - \gamma - \gamma$ coupling leads to an increased axion detection rate for microwave cavity experiments.

4 Concluding Remarks

In this paper, we have proposed a new anomaly-mediation paradigm model which evades the problems of 1. too low a value of $m_h$, 2. unnaturalness and 3. winolike LSPs which may be excluded by lack of IDD of dark matter. Our new model, dubbed natural anomaly-mediated SUSY breaking or nAMSB, merely incorporates the inclusion of non-universal bulk scalar masses and a bulk trilinear term $A_0$. The former allows for small $\mu$ as required by naturalness and leads instead to a higgsino-like WIMP as LSP. The inclusion of a bulk $A_0$ term allows for large stop mixing which lifts $m_h$ up to $\sim 125 \text{ GeV}$ whilst decreasing the top-squark radiative corrections to the scalar potential $\Sigma_u(\tilde{t}_{1,2})$. In fact, these revision were suggested by the model’s creators [2].

We computed the sparticle mass spectrum in nAMSB. While weak scale gaugino masses are still related as $M_1 : M_2 : M_3 \sim 0.4 : 0.13 : 1$ leading to wino as the lightest gaugino, the lightest charginos and neutralinos are instead mainly higgsino-like (but with a non-negligible wino component). These modifications bring the model into line with Higgs mass, naturalness and dark matter constraints. But they also greatly modify the collider and dark matter signatures which are expected from anomaly-mediation. Instead of quasi-stable charged winos leading to terminating tracks in collider experiments, now there are more rapidly decaying higgsinos at the bottom of the spectra. We computed upper bounds on gluino and top squark masses in nAMSB and found these to be possibly well beyond reach of HL-LHC although they should be accessible to HE-LHC. However, since higgsinos are required to be not too far from the
100 GeV scale, then the $\ell^+\ell^- j + E_T$ signature should likely be accessible to HL-LHC albeit with larger chargino and neutralino mass gaps than in models with unified gauginos. Also, the SSdB signature from wino pair production should be detectable over the entire natural range of wino masses in nAMSB leading to a conclusive test of this model. An ILC operating with $\sqrt{s} > 2m(higgsino)$ could also discover SUSY and unravel the underlying mediation mechanism via precision higgsino pair production measurements.

Dark matter is expected to consist of a higgsino-like WIMP plus axion admixture. Prospects for WIMP detection should be better than in natural models with gaugino mass unification due to the presence of rather light winos which enhance the SI DD scattering rates. Axions may remain difficult to detect. A further alternative is that the nAMSB model may provide a viable home for mixed axion/axino dark matter.

**Acknowledgments**

This work was supported in part by the US Department of Energy, Office of High Energy Physics.

**References**

[1] J. Polchinski, Phys. Rev. Lett. 75 (1995) 4724.
[2] L. Randall and R. Sundrum, Nucl. Phys. B 557 (1999) 79.
[3] I. Jack, D. R. T. Jones and A. Pickering, Phys. Lett. B 426 (1998) 73; L. V. Avdeev, D. I. Kazakov and I. N. Kondrashuk, Nucl. Phys. B 510 (1998) 289; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP 9812 (1998) 027; J. A. Bagger, T. Moroi and E. Poppitz, JHEP 0004 (2000) 009; A. Pomarol and R. Rattazzi, JHEP 9905 (1999) 013; M. K. Gaillard and B. D. Nelson, Nucl. Phys. B 588 (2000) 197; P. Binetruy, M. K. Gaillard and B. D. Nelson, Nucl. Phys. B 604 (2001) 32; M. Dine and N. Seiberg, JHEP 0703 (2007) 040; S. P. de Alwis, Phys. Rev. D 77 (2008) 105020; M. Dine and P. Draper, JHEP 1402 (2014) 069; F. D’Eramo, J. Thaler and Z. Thomas, JHEP 1206 (2012) 151.
[4] J. L. Feng, T. Moroi, L. Randall, M. Strassler and S. f. Su, Phys. Rev. Lett. 83 (1999) 1731.
[5] T. Gherghetta, G. F. Giudice and J. D. Wells, Nucl. Phys. B 559 (1999) 27.
[6] J. L. Feng and T. Moroi, Phys. Rev. D 61 (2000) 095004.
[7] Z. Chacko, M. A. Luty, I. Maksymyk and E. Ponton, JHEP 0004 (2000) 001; E. Katz, Y. Shadmi and Y. Shirman, JHEP 9908 (1999) 015; I. Jack and D. R. T. Jones, Phys. Lett. B 482 (2000) 167; M. Carena, K. Huitu and T. Kobayashi, Nucl. Phys. B 592 (2001) 164; R. Dermisek, H. Verlinde and L. T. Wang, Phys. Rev. Lett. 100 (2008) 131804; M. Hindmarsh and D. R. T. Jones, Phys. Rev. D 87 (2013) 075022.
[8] H. Baer, J. K. Mizukoshi and X. Tata, Phys. Lett. B 488 (2000) 367.
[9] A. J. Barr, C. G. Lester, M. A. Parker, B. C. Allanach and P. Richardson, JHEP 0303 (2003) 045; B. C. Allanach, T. J. Khoo and K. Sakurai, JHEP 1111 (2011) 132.

[10] G. Aad et al. [ATLAS Collaboration], JHEP 1301 (2013) 131.

[11] T. Moroi and L. Randall, Nucl. Phys. B 570 (2000) 455.

[12] A. Arbey, M. Battaglia, A. Djouadi, F. Mahmoudi and J. Quevillon, Phys. Lett. B 708 (2012) 162; A. Arbey, A. Deandrea and A. Tarhini, JHEP 1105 (2011) 078.

[13] H. Baer, V. Barger and A. Mustafayev, JHEP 1205 (2012) 091

[14] A. Arbey, A. Deandrea, F. Mahmoudi and A. Tarhini, Phys. Rev. D 87 (2013) no.11, 115020.

[15] H. Baer, V. Barger, P. Huang, A. Mustafayev and X. Tata, Phys. Rev. Lett. 109 (2012) 161802.

[16] H. Baer, V. Barger, P. Huang, D. Mickelson, A. Mustafayev and X. Tata, Phys. Rev. D 87 (2013) 11, 115028.

[17] J. Ellis, K. Enqvist, D. Nanopoulos and F. Zwirner, Mod. Phys. Lett. A 1 (1986) 57; R. Barbieri and G. Giudice, Nucl. Phys. B 306 (1988) 63.

[18] H. Baer, V. Barger, D. Mickelson and M. Padeffke-Kirkland, Phys. Rev. D 89 (2014) 115019.

[19] R. Kitano and Y. Nomura, Phys. Rev. D 73 (2006) 095004; M. Papucci, J. T. Ruderman and A. Weiler, JHEP 1209 (2012) 035; C. Brust, A. Katz, S. Lawrence and R. Sundrum, JHEP 1203 (2012) 103.

[20] H. Baer, V. Barger and D. Mickelson, Phys. Rev. D 88 (2013) no.9, 095013.

[21] A. Mustafayev and X. Tata, Indian J. Phys. 88 (2014) 991.

[22] H. Baer, V. Barger and M. Savoy, Phys. Rev. D 93 (2016) no.3, 035016.

[23] K. L. Chan, U. Chattopadhyay and P. Nath, Phys. Rev. D 58 (1998) 096004;

[24] H. Baer, V. Barger and P. Huang, JHEP 1111 (2011) 031.

[25] ISAJET 7.85, by H. Baer, F. Paige, S. Protopopescu and X. Tata, hep-ph/0312045; Isasugra, by H. Baer, C. H. Chen, R. B. Munroe, F. E. Paige and X. Tata, Phys. Rev. D 51 (1995) 1046.

[26] D. Matalliotakis and H. P. Nilles, Nucl. Phys. B 435 (1995) 115; M. Olechowski and S. Pokorski, Phys. Lett. B 344 (1995) 201; P. Nath and R. L. Arnowitt, Phys. Rev. D 56 (1997) 2820; J. Ellis, K. Olive and Y. Santoso, Phys. Lett. B539 (2002) 107; J. Ellis, T. Falk, K. Olive and Y. Santoso, Nucl. Phys. B652 (2003) 259; H. Baer, A. Mustafayev, S. Profumo, A. Belyaev and X. Tata, JHEP0507 (2005) 065.

[27] M. Dine, JHEP 0601 (2006) 162.

[28] H. Baer, C. Balazs and A. Belyaev, JHEP 0203 (2002) 042.

[29] H. Baer, V. Barger and H. Serce, Phys. Rev. D 94 (2016) no.11, 115019.
[30] M. L. Ahnen et al. [MAGIC and Fermi-LAT Collaborations], JCAP 1602 (2016) no.02, 039 doi:10.1088/1475-7516/2016/02/039 [arXiv:1601.06590 [astro-ph.HE]].

[31] H. Abdallah et al. [HESS Collaboration], arXiv:1607.08142 [astro-ph.HE].

[32] T. Cohen, M. Lisanti, A. Pierce and T. R. Slatyer, JCAP 1310 (2013) 061 doi:10.1088/1475-7516/2013/10/061 [arXiv:1307.4082].

[33] J. Fan and M. Reece, JHEP 1310 (2013) 124 doi:10.1007/JHEP10(2013)124 [arXiv:1307.4400 [hep-ph]].

[34] K. J. Bae, H. Baer, A. Lessa and H. Serce, Front. in Phys. 3 (2015) 49.

[35] B. C. Allanach et al., Eur. Phys. J. C 25 (2002) 113.

[36] The ATLAS collaboration [ATLAS Collaboration], ATLAS-CONF-2017-022; T. Sakuma [CMS Collaboration], PoS LHCP 2016 (2017) 145 [arXiv:1609.07445 [hep-ex]].

[37] J. E. Kim and H. P. Nilles, Phys. Lett. B 138, 150 (1984).

[38] M. Dine, W. Fischler and M. Srednicki, Phys. Lett. B 104 (1981) 199; A. R. Zhitnitsky, Sov. J. Nucl. Phys. 31 (1980) 260 [Yad. Fiz. 31 (1980) 497].

[39] K. J. Bae, H. Baer and E. J. Chun, Phys. Rev. D 89 (2014) 031701; K. J. Bae, H. Baer and E. J. Chun, JCAP 1312 (2013) 028.

[40] K. J. Bae, H. Baer and A. Lessa, JCAP 1304 (2013) 041.

[41] H. Baer, V. Barger, M. Padelford-Kirkland and X. Tata, Phys. Rev. D 89 (2014) no.3, 037701.

[42] H. Baer, V. Barger and M. Savoy, Phys. Rev. D 93 (2016) no.7, 075001.

[43] H. Baer, V. Barger, J. S. Gainer, P. Huang, M. Savoy, D. Sengupta and X. Tata, Eur. Phys. J. C 77 (2017) no.7, 499.

[44] H. Baer, V. Barger, J. S. Gainer, P. Huang, M. Savoy, H. Serce and X. Tata, Phys. Lett. B 774 (2017) 451.

[45] H. Baer, V. Barger, J. S. Gainer, H. Serce and X. Tata, Phys. Rev. D 96 (2017) no.11, 115008.

[46] See, e.g. ATLAS Phys. PUB 2013-011; CMS Note-13-002.

[47] H. Baer, V. Barger, P. Huang, D. Mickelson, A. Mustafayev, W. Sreethawong and X. Tata, JHEP 1312 (2013) 013 Erratum: [JHEP 1506 (2015) 053].

[48] H. Baer, A. Mustafayev and X. Tata, Phys. Rev. D 89 (2014) no.5, 055007.

[49] Z. Han, G. D. Krubs, A. Martin and A. Menon, Phys. Rev. D 89 (2014) no.7, 075007; H. Baer, A. Mustafayev and X. Tata, Phys. Rev. D 90 (2014) no.11, 115007; C. Han, D. Kim, S. Munir and M. Park, JHEP 1504 (2015) 132.

[50] CMS Collaboration [CMS Collaboration], CMS-PAS-SUS-16-048.
[51] Talk by B. Hooberman, SUSY 2017 meeting, Mumbai, India, Dec. 13, 2017.

[52] H. Baer, V. Barger, H. Serce and X. Tata, Phys. Rev. D 94 (2016) no.11, 115017.

[53] H. Baer, V. Barger, P. Huang, D. Mickelson, A. Mustafayev, W. Sreethawong and X. Tata, Phys. Rev. Lett. 110 (2013) no.15, 151801.

[54] H. Baer, V. Barger, J. S. Gainer, M. Savoy, D. Sengupta and X. Tata, arXiv:1710.09103 [hep-ph].

[55] H. Baer, V. Barger, D. Mickelson, A. Mustafayev and X. Tata, JHEP 1406 (2014) 172.

[56] H. Baer, M. Berggren, K. Fujii, S. L. Lehtinen, J. List, T. Tanabe and J. Yan, PoS ICHEP 2016 (2016) 156; S. L. Lehtinen, H. Baer, M. Berggren, K. Fujii, J. List, T. Tanabe and J. Yan, arXiv:1710.02406 [hep-ph].

[57] K. Fujii et al., arXiv:1702.05333 [hep-ph].

[58] K. J. Bae, H. Baer, A. Lessa and H. Serce, JCAP 1410 (2014) no.10, 082.

[59] E. J. Chun and A. Lukas, Phys. Lett. B 357 (1995) 43.

[60] J. E. Kim and M. S. Seo, Nucl. Phys. B 864 (2012) 296.

[61] K. J. Bae, H. Baer and H. Serce, JCAP 1706 (2017) no.06, 024.