Dark Matter Prospects in Deflected Mirage Mediation

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Abstract

The recently introduced deflected mirage mediation (DMM) model is a string-motivated paradigm in which all three of the major supersymmetry-breaking transmission mechanisms are operative. We begin a systematic exploration of the parameter space of this rich model context, paying special attention to the pattern of gaugino masses which arise. In this work we focus on the dark matter phenomenology of the DMM model as such signals are the least influenced by the model-dependent scalar masses. We find that a large portion of the parameter space in which the three mediation mechanisms have a similar effective mass scale of 1 TeV or less will be probed by future direct and indirect detection experiments. Distinguishing deflected mirage mediation from the mirage model without gauge mediation will prove difficult without collider input, though we indicate how gamma ray signals may provide an opportunity for distinguishing between the two paradigms.

1 Introduction

If supersymmetry is relevant at the electroweak scale then forthcoming experimental data should reveal the presence of new states which will be studied extensively at colliders. If the lightest supersymmetric particle (LSP) is stable it will provide an excellent candidate for explaining the presence of non-baryonic dark matter in the cosmos [1]. The phenomenology of these new states – both at colliders and in cosmological observations – will be determined by their interactions with each other and with the states of the Standard Model. These are determined, in turn, by the manner in which supersymmetry breaking is communicated from a hidden sector to the observable sector [2].

Over the past ten years most studies of supersymmetric phenomenology have tended to focus on one of three general paradigms for mediation of supersymmetry breaking to the observable sector. These are modulus (or “gravity”) mediation [3,4], gauge mediation [5] and anomaly mediation [6,7]. The latter two are often motivated by bottom-up concerns such as flavor-changing neutral currents and CP-violation, but the case of modulus mediation (a special string-motivated case of general
supergravity mediation) can avoid phenomenological constraints as well, provided the modulus or moduli in question have family-universal couplings to the observable sector and the underlying string theory has certain well-motivated isometries \[8, 9\]. These paradigms are typically studied in isolation, assuming that only one such mechanism dominates the pattern of soft supersymmetry breaking. As an example one could consider the hugely influential “Snowmass Points and Slopes” benchmark models designed for studies of collider phenomenology \[10\].

This is a reasonable starting point, but recent work suggests that there are well-motivated reasons to consider pairs of these mediation mechanisms – or perhaps all three mechanisms – simultaneously. The combination of anomaly mediation and gauge mediation, known as deflected anomaly mediation, was studied soon after the anomaly-mediated scenario was first proposed \[11, 12\]. For the combination of anomaly mediation and modulus mediation to produce comparable contributions to soft supersymmetry breaking a hierarchy of mass scales must be engineered. This was first observed in certain classes of heterotic orbifold models \[13, 14, 15\] but received increased attention with the advent of KKLT-type moduli stabilization in D-brane models \[16, 17\]. A particular class of Type IIB string compactifications with fluxes \[18\] gave rise to the so-called “mirage mediation” models \[19, 20, 21\] whose phenomenology has been extensively studied \[22, 23, 24, 25, 26, 27, 28\].

Recently the two combinations described above were combined in a scheme dubbed “deflected mirage mediation” (DMM) in which all three paradigmatic transmission mechanisms are operative and comparable \[29\]. While the DMM framework does not result in fully generalized superpartner mass possibilities – particularly for the all-important gaugino sector, as we will see below – it does provide a very rich phenomenology. Like the mirage mediation case which preceded it, some basic elements of the phenomenology of DMM have been studied within the contexts of specific string-based frameworks \[30, 31, 32\]. In this paper we embark on a program to study the following general question: if all three supersymmetry-breaking mediation mechanisms are present in roughly equal amounts, how will experimental observation reveal this fact? By phrasing the question in this manner we hope to study how certain “underlying principles” (such as the mechanism of supersymmetry-breaking mediation) can be extracted directly from data without reconstructing the full low-energy Lagrangian \[34, 35\].

Here we therefore wish to study primarily the gaugino sector, which reflects fewer model-dependent properties. Indeed, the DMM scenario continues to be a two-parameter family of gaugino mass patterns, with one parameter representing deviations from universality and the other the overall mass scale. In this paper we will consider the dark matter signals of the deflected mirage media-
tion scenario, which are determined largely by the physics of the gaugino sector alone. Of course full model-independence is impossible. Such important properties as the masses of the supersymmetric Higgs fields and the Higgsino content of the LSP will depend on soft supersymmetry-breaking scalar masses which are always model-dependent. We will consider the truly model-dependent implications for hadron collider physics in a future publication [36].

In its simplest manifestation the general mediation scenario would involve three mass scales, one each for the (bulk) gravity contribution, superconformal anomaly contribution and the gauge messenger contribution. This can be re-parameterized by a single overall scale and two dimensionless ratios. Such was the parameter set adopted by Everett et al. [29]. We will review the construction of the deflected mirage mediation model in Section 2 focusing most strongly on the gaugino sector where the interplay of supersymmetry-breaking mediation mechanisms is most transparent. We will study how the addition of gauge mediation alters the allowed parameter space of the theory. We then turn to dark matter observations in Section 3, indicating how the deflection arising from gauge-charged messengers alters the phenomenology relative to the un-deflected mirage model. Our survey will include a number of current and future experiments seeking to detect the presence of relic neutralinos. We will find that regions of the parameter space which result in a wino-like or mixed Higgsino/gaugino LSP will be probed in direct detection experiments at the one ton-year scale as well as in a variety of indirect detection experiments, provided that the effective mass scales of the various mediation mechanisms are in the 1 TeV range or less. These are also the areas in which the thermal relic abundance of the LSP is no greater than the upper bound inferred from the WMAP data. Regions of the parameter space with a heavier spectrum, or with a predominantly bino-like LSP, will prove much more difficult to observe via cosmological or astroparticle observations. We perform a basic survey of cosmic ray signals for the DMM model and find that none of the parameter space gives rise to an adequate explanation for the recent ATIC data, though some points can give a reasonable fit to the PAMELA data with boost factors less than $O(100)$. However, applying such boosts to the anti-proton flux as well would make these points in conflict with a lack of signal in the anti-proton data without modification of the diffusion model. In Section 4 we consider the extent to which the observations discussed in Section 3 will be helpful in distinguishing the deflected mirage mediation model from the previously-studied mirage pattern without gauge-charged messengers. Here we will encounter a “dark matter inverse problem” in trying to distinguish between the two cases. We will demonstrate why this occurs and suggest some possible avenues for future study to resolve this problem.


2 Parameterizing Deflected Mirage Mediation

In this section we will review the construction of the soft Lagrangian in deflected mirage mediation. As this material has been presented more fully elsewhere [30, 32] we will be brief in our discussion of the underlying string-theoretical context. We begin in Section 2.1 with a basic derivation of the gaugino mass patterns in deflected mirage mediation in a model-independent manner. This will allow us to introduce some important notation and outline the parameterizations we will be using throughout the rest of the work in Section 2.2. In Section 2.3 we complete the construction by introducing the soft supersymmetry-breaking scalar masses and trilinear A-terms and explore the allowed regions in the parameter space.

2.1 Gaugino Masses in Terms of Mass Scales

In this subsection we present the construction of the gaugino mass patterns without any reference to its possible origin from string-theoretic considerations. The treatment here closely parallels that of [35] in the absence of gauge-charged messengers. For the sake of theoretical clarity we will work to one-loop order in the renormalization group equations in this section. We begin by assuming three contributions to the soft supersymmetry breaking gaugino masses of the MSSM. Let the contribution from Planck-suppressed operators be universal in form with a mass scale given by the quantity $M_0$. The contributions from the superconformal anomaly will be proportional to some (a priori different) mass scale given by $M_g$. We assume that these contributions arise at some effective high-energy scale $\mu_{uv}$ at which supersymmetry breaking is transmitted from some hidden sector to the observable sector. It is common to take this scale to be the GUT scale at which gauge couplings unify, but in string constructions one might choose a different (possibly higher scale) at which the supergravity approximation for the effective Lagrangian becomes valid. Finally, we will include a contribution from some gauge-charged messenger sector which is proportional to a third mass scale $\Lambda_{mess}$. As in the original work of [29, 30] we will assume this messenger sector comes in complete GUT representations of the Standard Model gauge group so as to preserve the successful gauge coupling unification of the MSSM. In particular we assume $N_m$ copies of $5,\bar{5}$ representations under $SU(5)$, which give rise to contributions to soft supersymmetry breaking gaugino masses at the energy scale $\mu_{mess} < \mu_{UV}$.

The full gaugino mass at the high energy boundary condition scale $\mu_{UV}$ is given by the expression [13, 37]

$$M_a (\mu_{UV}) = M_0 + g_a^2 (\mu_{UV}) \frac{b'_a}{16\pi^2} m_{3/2},$$

(1)
where \(a = 1, 2, 3\) labels the Standard Model gauge group factors \(G_a\) and \(m_{3/2}\) is the value of the gravitino mass, given by \(m_{3/2} = \left< e^{K/2} W \right>\). The quantities \(b'_a\) are the beta-function coefficients for the Standard Model gauge groups. In our conventions these are given by

\[
b_a = -(3C_a - \sum_i C^i_{a}),
\]

where \(C_a, C^i_a\) are the quadratic Casimir operators for the gauge group \(G_a\), respectively, in the adjoint representation and in the representation of the matter fields \(\Phi^i\) charged under that group. As the summation in \((2)\) is over all degrees of freedom in the theory present at the scale \(\mu_{uv} > \mu_{mess}\), it must include the contribution from the gauge-charged messenger sector. Therefore we may write

\[
b'_a = b_a + N_m
\]

with \(b_a\) the beta-function coefficients of the MSSM in the absence of these messenger fields

\[
\{b_1, b_2, b_3\} = \left\{ \frac{33}{5}, 1, -3 \right\}.
\]

Note that in the absence of the messenger sector \((N_m = 0)\) if we take \(\mu_{uv} = \mu_{GUT}\) then we would have

\[
g^2_1 (\mu_{uv}) = g^2_2 (\mu_{uv}) = g^2_3 (\mu_{uv}) = g^2_{GUT} \simeq \frac{1}{2}.
\]

In the presence of the gauge-charged messengers the unification of gauge couplings continues to occur at a scale \(\mu_{GUT} \simeq 2 \times 10^{16}\) GeV but the value of the unified gauge coupling differs from that of \((5)\) by the relation

\[
\frac{1}{g^2_a (\mu_{uv})} = \frac{1}{g^2_{GUT}} - \frac{N_m}{8\pi^2} \ln \left( \frac{\mu_{GUT}}{\mu_{mess}} \right)
\]

which is valid at one-loop order for all values of \(a\).

We now imagine evolving the expressions in \((1)\) to the scale \(\mu_{mess}\) via the one-loop renormalization group equations. The first term in \((1)\) is renormalized by a multiplicative factor at one loop

\[
M_{a}^{\text{term 1}} (\mu_{mess}) = M_0 \left[ 1 - g^2_a (\mu_{mess}) \frac{b'_a}{8\pi^2} \ln \left( \frac{\mu_{uv}}{\mu_{mess}} \right) \right],
\]

while the second term in \((1)\) can be evolved by simply replacing the gauge coupling with its value at the intermediate mass scale

\[
M_{a}^{\text{term 2}} (\mu_{mess}) = g^2_a (\mu_{mess}) \frac{b'_a}{16\pi^2} m_{3/2}.
\]
Combining (7) and (8) we have

\[ M_a (\mu_{\text{mess}}) = g_a^2 (\mu_{\text{mess}}) \frac{b_a}{16\pi^2} m_{3/2} + M_0 \left[ 1 - g_a^2 (\mu_{\text{mess}}) \frac{b_a'}{8\pi^2} \ln \left( \frac{\mu_{\text{UV}}}{\mu_{\text{mess}}} \right) \right]. \]  

(9)

Clearly, if the two parts of the expression in (9) are to be of roughly the same size it will be necessary to engineer a situation in which \( m_{3/2} \approx 16\pi^2 M_0 \).

At this intermediate scale we may now integrate out the gauge-charged messenger sector, which in the presence of supersymmetry breaking parameterized by \( m_{3/2} \) generates a threshold correction to the gaugino masses of the form

\[ \Delta M_a = -N_m g_a^2 (\mu_{\text{mess}}) \frac{b_a}{16\pi^2} \left( \Lambda_{\text{mess}} + m_{3/2} \right). \]  

(10)

Evolving the combination \( M_a (\mu_{\text{mess}}) + \Delta M_a \) from the intermediate scale \( \mu_{\text{mess}} \) to the low-energy scale \( \mu_{\text{EW}} \) (again at one-loop) results in the final expression

\[ M_a (\mu_{\text{EW}}) = g_a^2 (\mu_{\text{EW}}) \frac{b_a}{16\pi^2} m_{3/2} + \left[ 1 - g_a^2 (\mu_{\text{EW}}) \frac{b_a'}{8\pi^2} \ln \left( \frac{\mu_{\text{UV}}}{\mu_{\text{mess}}} \right) \right] \times \left[ M_0 \left( 1 - g_a^2 (\mu_{\text{mess}}) \frac{b_a'}{8\pi^2} \ln \left( \frac{\mu_{\text{UV}}}{\mu_{\text{mess}}} \right) \right) - N_m g_a^2 (\mu_{\text{mess}}) \frac{b_a}{16\pi^2} \Lambda_{\text{mess}} \right]. \]  

(11)

From (11) we immediately see that we must engineer the relations \( \Lambda_{\text{mess}} \approx m_{3/2} \approx 16\pi^2 M_0 \) if all three contributions to the gaugino masses are to be comparable in magnitude. Let us note before going forward that the preceding derivation made tacit assumptions about the relative phases between the various contributions in the final expression (11). Given a true top-down construction these phases should be determined by the same mechanism which stabilizes the auxiliary fields in the supergravity Lagrangian.

2.2 A Theory-Guided Parametrization

The masses in (11) depend explicitly on the three mass scales \( M_0, m_{3/2} \) and \( \Lambda_{\text{mess}} \) as well as the number of messenger multiplets \( N_m \). They also depend implicitly on the value of the intermediate scale \( \mu_{\text{mess}} \) through the logarithms and gauge couplings. This somewhat unwieldy expression can be simplified by relating the explicit mass scales to some overall mass scale via two dimensionless ratios. Following Everett et al. we can define

\[ \alpha_g = \frac{\Lambda_{\text{mess}}}{m_{3/2}} \]  

(12)
which will naturally be of order unity for the cases of interest to us here. To obtain another ratio of $O(1)$ it is common to define

$$\alpha_m \equiv \frac{m_{3/2}}{M_0 \ln \left( \frac{M_{pl}/m_{3/2}}{2} \right)},$$

(13)

where $M_{pl}$ is the reduced Planck mass $M_{pl} = 2.4 \times 10^{18}$ GeV. Again, for the models of interest to us here we will have $\alpha_m$ typically of order unity. Using the definitions (12) and (13) the expression in (11) can be re-expressed as

$$\frac{M_a(\mu_{EW})}{M_0} = g_a^2(\mu_{EW}) \alpha_m \frac{b_a}{16\pi^2} \ln \left( \frac{M_{pl}/m_{3/2}}{2} \right) + \left[ 1 - g_a^2(\mu_{EW}) \frac{b_a}{8\pi^2} \ln \left( \frac{\mu_{mess}}{\mu_{EW}} \right) \right]$$

$$\times \left[ 1 - g_a^2(\mu_{mess}) \frac{b_a}{8\pi^2} \ln \left( \frac{\mu_{mess}}{\mu_{EW}} \right) - \alpha_m \alpha_g N_m g_a^2(\mu_{mess}) \frac{16\pi^2}{16\pi^2} \ln \left( \frac{M_{pl}/m_{3/2}}{2} \right) \right].$$

(14)

With these conventions one can define the so-called “mirage scale” $\mu_{mir}$ at which the three soft supersymmetry breaking gaugino masses have an identical value as

$$\mu_{mir} = \mu_{GUT} \left( \frac{m_{3/2}}{M_{pl}} \right)^{\alpha_m \rho/2},$$

(15)

where the parameter $\rho$ is given by

$$\rho = \left[ 1 + \frac{N_m g_a^2(\mu_{UV})}{8\pi^2} \ln \left( \frac{\mu_{GUT}}{\mu_{mess}} \right) \right] / \left[ 1 - \alpha_m \alpha_g N_m g_a^2(\mu_{UV}) \frac{16\pi^2}{16\pi^2} \ln \left( \frac{M_{pl}/m_{3/2}}{2} \right) \right].$$

(16)

The expression in (14) is still somewhat cumbersome. It was noted by Choi [31] that the gaugino mass expressions (11) at the low scale $\mu_{EW}$ can be re-arranged into the form

$$M_a(\mu_{EW}) = M_0^{\text{eff}} \left\{ 1 + \beta_a(\mu_{EW}) \ln \left( \frac{\mu_{EW}}{\mu_{mir}} \right) \right\},$$

(17)

where we have introduced the variable

$$\beta_a(\mu) = \frac{b_a g_a^2(\mu)}{8\pi^2}$$

(18)

and have defined the effective mass scale

$$M_0^{\text{eff}} = R M_0; \quad R = 1 - \frac{N_m g_G^{2}(\mu_{UV})}{8\pi^2} \left\{ \frac{\alpha_m \alpha_g}{2} \ln \left( \frac{M_{pl}}{m_{3/2}} \right) + \ln \left( \frac{\mu_{GUT}}{\mu_{mess}} \right) \right\}.$$
$g_a^2(\mu_{\text{UV}})$ appearing in (14) via (6). We note that the dimensionless quantity $R$ is simply the inverse of the parameter $\rho$ in (16), and thus

$$\mu_{\text{mir}} = \mu_{\text{GUT}} \left( \frac{m_{3/2}}{M_{\text{pl}}} \right)^{\alpha_m/2R}.$$  

(20)

In [32] the expression in (17) was further modified by introducing two new dimensionless variables based on the quantity $R$ in (19)

$$x = \frac{1}{R + \alpha_m}, \quad y = \frac{\alpha_m}{R + \alpha_m},$$  

(21)

which can be used to write (17) in the form

$$M_a(\mu) = M_0 (1 + \beta_a(\mu) t) x^{-1} \left\{ 1 + y \left[ \frac{\beta_a(\mu) t'}{1 + \beta_a(\mu) t} - 1 \right] \right\},$$  

(22)

where we have defined two scaling variables $t$ and $t'$ via

$$t = \ln \left( \frac{\mu}{\mu_{\text{GUT}}} \right), \quad t' = \frac{1}{2} \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right).$$  

(23)

The expression in (22) is intended to be evaluated at the energy scale $\mu = \mu_{\text{EW}} < \mu_{\text{mess}}$. The free parameters in (22) are $\{M_0, x, y\}$, though (22) is still ultimately a function of only one scale and one dimensionless variable (as we will explicitly show in Section 4). Having specified $x$ and $y$ values one may obtain $\alpha_m = y/x$, and then the value of $m_{3/2}$ can be obtained from that of $M_0$ by finding the solution to the definitional equation for $\alpha_m$ in (13). Though the series of definitions in (19) and (21) may appear to sacrifice clarity for the sake of brevity, they allow an extremely powerful way to survey all possible mediation mechanisms in a simple two-dimensional plane.

2.3 Exploring the Parameter Space

As we have seen, the gaugino mass sector in deflected mirage mediation is particularly transparent. Though the three soft supersymmetry-breaking gaugino masses are sufficient to determine a large part of the dark matter phenomenology of the LSP neutralino, the scalar sector will also contribute in important, indirect ways. These include determining the $\mu$-parameter via the electroweak symmetry-breaking (EWSB) conditions and thereby the LSP wave-function, and by allowing certain important resonant annihilation or co-annihilation processes in the early universe. Therefore we cannot adequately discuss the prospects for dark matter signals in deflected mirage mediation
without specifying the rest of the soft supersymmetry breaking Lagrangian, particularly the scalar masses and trilinear couplings.

Doing so requires working within a model framework. We will here very briefly outline that framework and the relevant expressions needed to go forward. Further background can be found in the relevant references, particularly \cite{30, 32}. We imagine Type IIB string theory compactified on a Calabi-Yau orientifold. To this is added flux in the manner of KKLT, producing a region of highly warped geometry. The remainder of the compact space remains (approximately) a Calabi-Yau manifold. These background fluxes serve two purposes: the first is to stabilize most of the geometrical moduli associated with the construction. The second (and related) purpose is to produce a constant in the superpotential for the remaining modulus of sufficient size to produce stabilization with $\alpha_m \sim O(1)$. This remaining modulus is a Kähler modulus which can be stabilized by field-theoretic mechanisms such as gaugino condensation. The observable sector (i.e. the MSSM) resides on $D_3$ or $D_7$ branes (or both) in the bulk of the internal manifold. In this minimal set-up it is plausible that a conspiracy of scales can be achieved in which direct modulus-mediation (via non-vanishing $F$-terms for the Kähler modulus) and loop-induced anomaly-mediation are competitive contributions to observable sector scalar masses.

Let us denote the (chiral) Kähler modulus superfield by $T$ and its highest auxiliary component by $F_T$. If the gauge kinetic functions $f_a$ for the three Standard Model gauge groups are universally given by $f_a = T$, then a non-vanishing vacuum value for $F_T$ will generate a universal component for the gaugino masses as in (1)

$$M_0 = \langle \frac{F_T}{t + \bar{t}} \rangle,$$

(24)

where $t = T|_{\theta=0}$ is the lowest (scalar) component of the superfield $T$. The quantity in (24) will also generate scalar masses and trilinear scalar couplings. The exact form depends on how the Kähler modulus appears in the Kähler metric for the matter superfields, if at all. Such couplings are difficult to compute, but their leading behavior at large volume (large $\langle t + \bar{t} \rangle$) can be readily inferred and parameterized. We will assume a Kähler potential of the form

$$K = \sum_i (T + \bar{T})^{-n_i}\bar{\Phi}_i\Phi_i,$$

(25)

with modular weight $n_i$ for the matter superfield $\Phi_i$. These numbers depend on the location of the observable sector matter fields and may be vanishing. It was argued in \cite{39}, using the work of \cite{40}, that for intersecting $D$-brane models in the Type IIB context one should expect $n_i = 0$ if $\Phi_i$ is localized on $D_3$ branes, $n_i = 1$ if $\Phi_i$ is localized on $D_7$ branes, and $n_i = 1/2$ for matter localized at
the intersection of branes. We point out, however, that no explicit model realizing both the fluxed stabilization mechanism and the MSSM matter content has yet been constructed.

In the absence of gauge-charged messengers, the contributions of (24) and the assumption in (25) are sufficient to specify the form of the scalar sector soft supersymmetry breaking terms at a scale $\mu_{\text{UV}}$. Such a system defines the mirage mediation (or mixed modulus/anomaly mediation) scenario. To this we now wish to add $N_m$ copies of gauge-charged messengers in the representation $5, \bar{5}$ under $SU(5)$. The presence of such terms introduces new contributions to the soft scalar masses and gaugino masses at the scale $\mu_{\text{mess}}$ at which these fields are integrated out of the spectrum. The calculation of such terms follows the general procedure of gauge mediation models [5]. Soft supersymmetry breaking is transmitted to the observable sector states via loop diagrams involving the messengers. The size of these contributions is determined by the parameter $\Lambda_{\text{mess}}$ appearing in (10). In principle this mass scale can be of any size as it is a priori independent of the physics of moduli stabilization. The relative size of $\Lambda_{\text{mess}}$ to $m_{3/2}$ (and hence the parameter $\alpha_g$) will depend on the model one postulates for dynamically generating masses and splittings for the messenger sector. Circumstances for which $\alpha_g$ is of order unity are not hard to construct [29, 38, 41].

The soft terms for the scalar masses and trilinear couplings, in the presence of all three supersymmetry-breaking transmission mechanisms, are given at the initial scale $\mu_{\text{UV}} > \mu_{\text{mess}}$ by [30]

$$A_{ijk}(\mu_{\text{UV}}) = M_0 \left[ (3 - n_i - n_j - n_k) - \frac{(\gamma_i + \gamma_j + \gamma_k)}{16\pi^2} \alpha_m \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right) \right]$$

$$m_i^2(\mu_{\text{UV}}) = M_0^2 \left[ (1 - n_i) - \frac{\theta'_i}{16\pi^2} \alpha_m \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right) - \frac{\gamma'_i}{(16\pi^2)^2} \left( \alpha_m \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right) \right)^2 \right],$$

with $\alpha_m$ being the parameter introduced in (13) and the quantities $\gamma_i, \theta'_i$ and $\gamma'_i$ being functions of Standard Model gauge and Yukawa couplings. Their definitions and explicit values can be found in [30]. The primes on the quantities in (27) indicate that these parameters are computed taking into account the presence of the messenger sector as well as the MSSM matter content. At the scale $\mu = \mu_{\text{mess}}$ the messenger sector is integrated out of the theory, producing the contribution [10] for the gaugino masses

$$\Delta M_a = -N_m M_0 g_a^2(\mu_{\text{mess}}) \alpha_m (1 + \alpha_g) \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right),$$

and a contribution

$$\Delta m_i^2 = M_0^2 \sum_a (2N_m C_a) g_a^4(\mu_{\text{mess}}) \frac{\alpha_m (1 + \alpha_g) \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right)}{(16\pi^2)^2} \left[ \alpha_m (1 + \alpha_g) \ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right) \right]^2$$

(29)
for the scalar masses. These expressions can be re-expressed in terms of the parameter $R$ of (19) and the set $\{x, y\}$ of (21), but doing so is not particularly illuminating.

Unlike the case of gaugino masses, no analytic expression at low scales is available for the scalar masses and trilinear couplings – even at one loop order – except in certain approximations [32]. Therefore when we wish to consider the full model-dependent set of soft parameters we must connect energy scales using the renormalization group (RG) equations. We do so in two stages. The first connects $\mu_{\text{UV}} = \mu_{\text{GUT}}$ and $\mu_{\text{mess}}$ with boundary conditions given by (1), (26) and (27). At this scale the corrections (28) and (29) are added and the soft terms are evolved from $\mu_{\text{mess}}$ to $\mu_{\text{EW}} = M_Z$. From here we pass the low-scale soft parameters to SuSpect 2.4 [42] to check for proper electroweak symmetry breaking and to compute the physical masses of the various superpartners.

We will investigate the parameter space defined by the set $\{x, y, M_0\}$ by considering slices of the $\{x, y\}$ plane for various values of $M_0$. We do this in two different regimes. In the first case we investigate the explicit model described in this section for the case where $\tan\beta = 10$, $\mu_{\text{mess}} = 10^{10}$ GeV, $N_m = 3$ and where the modular weights of the various MSSM states are given by

$$\{n_Q, n_U, n_D, n_L, n_E, n_{H_u}, n_{H_d}\} = \{1/2, 1/2, 1/2, 1/2, 1/2, 1, 1\}.$$ (30)

This choice of modular weights is not motivated by any particular construction; we choose it to match with certain benchmark models from the literature. We will call this the “model-dependent scenario.” In the second case we will dispense with the model-dependent soft-terms by simply setting all trilinear couplings to zero, set $\mu = m_A = 1$ TeV and set all scalar masses to the value of the gluino soft mass $M_3$ or to 1 TeV, whichever is larger. We do this by hand at the low energy scale. While artificial, this provides some sense of how the phenomenology of the gaugino sector alone is influenced by the choice of $\{x, y\}$ for a given value of $M_0$. We will call this second case the “model-independent scenario.”

In Figure 1 we display the physically allowed region in the $\{x, y\}$ plane for various values of the mass scale parameter $M_0$ in the model-dependent case with modular weights given by (30). The primary constraint in the $\{x, y\}$ plane comes from the direct search limits for gauginos. The mass bounds arising from these searches are somewhat model-dependent, particularly for the gluino. We will require that the lightest neutralino have a mass $m_{\tilde{N}_1} \geq 46$ GeV and that the lightest chargino mass obeys $m_{\tilde{C}_1} \geq 103$ GeV [43]. For the gluino we will be much more conservative. Though recent searches at the Tevatron have become increasing model-independent, we do not wish to rule out too much of the parameter space of this model class prematurely. We will therefore require only that $m_{\tilde{g}} \geq 200$ GeV [44] in our figures, though a bound of $m_{\tilde{g}} \geq 300$ GeV [45] will be more
Figure 1: **Allowed Parameter Space for the “Model-Dependent Scenario.”** Contours indicate the locus of points for which $m_{\tilde{N}_1} = 46$ GeV (dashed contour), $m_{\tilde{C}_1} = 103$ GeV (solid contour) and $m_{\tilde{g}} = 200$ GeV (dash-dotted contour). The dark shaded region is the area in which the stau is the LSP, while the smaller hatched region in the center has a stop LSP for the modular weight choice of (30). The hatched region in the upper left of each plot indicates where no EWSB occurs. For larger values of the parameter $M_0$ we have indicated the area in which the gluino (or the chargino) is the LSP by the darker shaded region(s). The labeled points are the benchmark models of Table 1. The intersecting lines indicate those points for which $\alpha_{\text{em}} = 1$ (lower left to upper right) and where $R = 0$ (upper left to lower right). The intersection of these two curves designates the prediction of the simplest KKLT scenario.
Table 1: Benchmark Models. The relevant mass scales are given in the first three entries, with the parameters of Everett et al. in the second block. The final block re-casts these parameters in terms of the parameterizations of Choi et al. from (21). All models are defined with positive value of $\mu$ and $\tan\beta = 10$. Models A-B were considered in [29] while Model C was considered in [30]. Model F is a mirage model without messengers near the prediction for the basic KKLT model.

|       | Model A | Model B | Model C | Model D | Model E | Model F |
|-------|---------|---------|---------|---------|---------|---------|
| $M_0$ | 2000 GeV| 1000 GeV| 1000 GeV| 500 GeV | 500 GeV | 933 GeV |
| $m_{3/2}$ | 62.6 TeV | 32.0 TeV | 32.0 TeV | 13.2 TeV | 51.3 TeV | 31.7 TeV |
| $\mu_{\text{mess}}$ | $10^{12}$ GeV | $10^{8}$ GeV | $10^{12}$ GeV | $10^{10}$ GeV | $10^{10}$ GeV | N/A |
| $\alpha_m$ | 1       | 1       | 1       | 0.81    | 3.26    | 1       |
| $\alpha_g$ | 1       | -1/2    | -1      | 0.14    | 1.47    | 0       |
| $N_m$ | 3       | 3       | 3       | 3       | 3       | 0       |
| $x$ | 0.668   | 0.561   | 0.472   | 0.674   | 0.401   | 0.5     |
| $y$ | 0.668   | 0.561   | 0.472   | 0.543   | 1.307   | 0.5     |

Table 1: Benchmark Models. The relevant mass scales are given in the first three entries, with the parameters of Everett et al. in the second block. The final block re-casts these parameters in terms of the parameterizations of Choi et al. from (21). All models are defined with positive value of $\mu$ and $\tan\beta = 10$. Models A-B were considered in [29] while Model C was considered in [30]. Model F is a mirage model without messengers near the prediction for the basic KKLT model.

applicable over many regions of the parameters. The lightly shaded region is ruled out by these direct experimental constraints. In addition there is some parameter space for which the gluino (and for very large $M_0$, the chargino) can be the LSP, indicated by the darker shading. For this choice of modular weights the stau is the LSP for smaller values of $x$ and $y$, indicated by the dark shaded area. There is also a small region where the lightest stop can be the LSP, indicated by the hatched region near the center of each figure. Finally in the upper left corner of each panel there is a region for which $m_A^2 < 0$ for the pseudoscalar Higgs and no electroweak symmetry breaking occurs.

This choice of parameter space offers a distinct advantage: the various theoretical model limits can easily be defined in terms of the dimensionless quantities $x$ and $y$. In the $\{x, y\}$ plane cases for which $R = 0$ lie along the line for which $x + y = 1$. The line extending from the origin defined by $x = y$ is the case with $\alpha_m = 1$ for arbitrary $\alpha_g$. The intersection of these two lines is the simplest prediction of the KKLT framework which inspired the mirage meditation scenario. We have imposed these theoretical model lines on top of the allowed space in Figure 1. The gauge-mediated limit in Figure 11 is formally at the origin of the $\{x, y\}$ plane where $R^{-1} \to 0$, but this limit must be taken while keeping the product $RM_0$ finite [32] and is therefore impossible to reach in our figure. So too, the pure anomaly-mediated limit is impossible to reach since this is formally the point $\{0, 1\}$ (i.e. $\alpha_m \to \infty$) where the product $\alpha_m M_0$ is held fixed. The pure modulus (or
Figure 2: LSP Properties for the “Model-Dependent Scenario.” Contours of constant LSP mass $m_{\tilde{N}_1}$ are given for the scenario with modular weights of (90) for mass scale $M_0 = 500$ GeV (left panel) and $M_0 = 1000$ GeV (right panel). The lightly shaded region is ruled out for the reasons indicated in Figure 1. For both values of $M_0$ the LSP is primarily bino-like for $y < \sim 0.6$ and $y > \sim 1.4$. For the case of $M_0 = 500$ GeV there is also some parameter space where the LSP is a mixture of Higgsino and wino, indicated by the hatched region.

gravity) mediated limit is the point $\{1, 0\}$ in the figures. For large enough values of $M_0$ there is an expansive parameter space beyond the mirage mediation “frontier” for which a compressed spectrum of gauginos occurs. The marked points in Figure 1 represent specific parameter choices which we will single out for particular study later in the work. The parameter sets which define these cases are given in Table 1.

In Figure 2 we concentrate on the properties of the lightest neutralino for the specific cases of $M_0 = 500$ GeV (left panel) and $M_0 = 1000$ GeV (right panel). The union of all the theoretically and experimentally forbidden regions in Figure 1 is indicated by the light shading. Note that the abrupt kinks in the contours of constant $m_{\tilde{N}_1}$ are the result of level crossings where the neutralino goes from being over 95% wino-like (along $y \simeq 1$) to over 95% bino-like ($y \gtrsim 1.4$ and $y \lesssim 0.6$). The region with a wino-like LSP is indicated in Figure 2 by the nearly horizontal dashed lines. For the case of $M_0 = 500$ GeV there is also some parameter space where the LSP is a mixture of Higgsino and wino, indicated by the hatched region. The analogous region for the case of $M_0 = 1000$ GeV
would lie within the phenomenologically forbidden region.

Some of the remaining superpartner masses are given in Figure 3, again for the cases of $M_0 = 500$ GeV (left panel) and $M_0 = 1000$ GeV (right panel). As before, the lightly shaded region is excluded. The upper panels show contours of constant gluino mass (solid heavy lines) and constant (lightest) chargino mass (dashed lines) for the entire $\{x, y\}$ plane. The curved shaded region represents the area where $1.5 \leq m_A/m_{\tilde{N}_1} \leq 2.5$. In this region we expected enhanced annihilation of relic neutralinos though the pseudoscalar Higgs resonance – particularly along the curved dashed contour for which $m_A = 2m_{\tilde{N}_1}$. Note that we have restricted our attention to the region in which both $x$ and $y$ lie between zero and one. This is the range explored in the mirage family of models.

In the lower two panels we plot several superpartner masses as a function of the parameter $x$ where we take $N_m = 0$, i.e. a true “mirage” model. The phenomenology in the lower panels roughly corresponds to following the line $x + y = 1$ in the upper two panels. Note that the correspondence between the plots, while very good, is not precise. Strictly speaking, the line $x + y = 1$ is determined by the requirement that $R = 1$. But from the definition in (19) it is apparent that there are two solutions for this constraint – one is the mirage limit where $N_m = 0$. If we fix the number of messengers at a value $N_m \neq 0$ (as in the case of the top panels in Figure 3) then the line $x + y = 1$ forces the second solution in which the quantity in braces in (19) vanishes. This is necessarily a solution for which $\alpha_g$ does not vanish. Nevertheless, the phenomenology is similar to that of the mirage model near this line, so the reader can use the line $x + y = 1$ as a reasonable proxy for the mirage model limit.

Finally, we wish to point out that much of the theoretically forbidden region is an artifact of the choice (30) for modular weights. It is clear from (26) and (27) that the choice of modular weights will have $O(1)$ effects on certain key parameters. This, in turn, affects the low-scale physical masses and the derived value of the $\mu$-parameter once electroweak symmetry breaking is imposed. The amount of the $\{x, y\}$ plane available is therefore heavily dependent on these model choices, as was pointed out in [24, 25, 26]. We therefore show the allowed parameter space in Figure 4 for the “model-independent” scenario in which we set $\mu = m_A = 1$ TeV and set all scalar masses to the value of the gluino soft mass $M_3$ or to 1 TeV, whichever is larger. As $x \to 0$ all gauginos get extremely massive for fixed $M_0$ value. Thus, near the $x = 0$ axis there is always a region where $m_{\tilde{N}_1} = 1$ TeV, indicated by the dotted contour, where a scalar particle would be the LSP for our choice of scalar masses. As the gluino gets increasingly massive in the $x \to 0$ limit it induces large radiative corrections to the Higgs potential, producing regions where $m_h \leq 100$ GeV (upper left corners) or where $m_A^2 \leq 0$ (lower left corners). The exact location of these contours depends on
Figure 3: **Key Superpartner Masses for the “Model-Dependent Scenario.”** Upper panels give various masses for the \( \{x, y\} \) plane, while the lower panels plot masses along the line \( x + y = 1 \) with \( N_m = 0 \). Panels on the left take \( M_0 = 500 \) GeV while those on the right take \( M_0 = 1000 \) GeV. The lightly shaded region in all plots is phenomenologically forbidden. The darker shaded region in the upper plots shows the area where \( 1.5 \leq m_A / m_{\tilde{N}_1} \leq 2.5 \), with \( m_A = 2m_{\tilde{N}_1} \) given by the curved dashed line. In the upper plot heavy solid lines are contours of constant \( m_\phi \) while dashed lines are contours of constant \( m_{\tilde{C}_1} \). All masses are in GeV.
the choices made for $\mu$ and the soft scalar masses.

As we have seen the allowed parameter space depends on the scalar sector of the theory, which is set by the modular weights assigned to the multiplets of the MSSM. The discrete choices for modular weights can be seen as selecting (a) the relative size of the gaugino masses versus scalar masses and (b) determining the degree of non-universality in the scalar masses between the Higgs sector and the matter sector. The former is akin to the choice of $m_0$ versus $m_{1/2}$ in mSUGRA models, while the latter is related to certain non-universal extensions of these mSUGRA models. In order to capture this physics, in the next section we will consider the dark matter signals of the LSP neutralino in terms of the specific benchmark case of (30) over the space defined by $\{x, y, M_0\}$. But we will also occasionally relax this heavy constraint and randomize over all other parameters that determine the model. Together these two strategies will give a reasonable depiction of the dark matter implications of the gaugino masses in (22).
3 Survey of Dark Matter Signatures

3.1 Thermal Relic Density

The phenomenology of relic neutralino dark matter consists of two largely disconnected aspects: the cosmological density of relic neutralinos and the physics of these neutralinos here in our galaxy today. These aspects are logically distinct, but can be related in any comprehensive theory of the origin and nature of dark matter. The physics of neutralino annihilation, or neutralino interaction with terrestrial dark matter detectors, depends on not much more than the properties of the LSP itself: its wave-function and mass. It also depends on its density in our local halo and certain other astrophysical properties independent of the nature of the LSP. The cosmological relic abundance, on the other hand, can be determined by this information only in certain special cases. More often it depends on the physical masses of other particles (supersymmetric states and Higgs states) that were part of the relativistic plasma at some earlier, hotter time in the cosmos. Crucially, the calculation of the cosmological relic density depends on certain assumptions about the thermal history of the universe.

The most important inputs to this thermal relic abundance calculation are the mass of the LSP and the mass difference between it and the next lightest superpartner. The thermal relic abundance of the LSP neutralino for the theory with modular weights (30) is given in Figure 5 as a function of the dimensionless parameters $x$ and $y$ for the cases $M_0 = 500$ GeV (left panel) and $M_0 = 1000$ GeV (right panel). Here and throughout, all calculations are performed using the computer package DarkSUSY 5.0.4 [46] after computing the physical mass spectrum throughout the parameter space using SuSpect 2.4. The WMAP three-year data [47] is best fit by a relic density in the range

$$0.0855 \leq \Omega_{\chi} h^2 \leq 0.1189,$$

at the $2\sigma$ level. In our figures we will expand this region somewhat to a “WMAP preferred” region

$$0.07 \leq \Omega_{\chi} h^2 \leq 0.14,$$

which will be slightly easier to resolve than the very narrow band in (31).

In Figure 5 this “WMAP preferred” region is indicated by the narrow red shaded region. Inside the parameter space with a wino-like LSP, and within the $A$-funnel region, the relic density drops below the value of $\Omega_{\chi} h^2 = 0.025$, indicated by the yellow shading. The intermediate regime of $0.025 \leq \Omega_{\chi} h^2 \leq 0.07$ is indicated by blue shading. The green area in the plot has $0.14 \leq \Omega_{\chi} h^2 \leq 1$ while the gray area has $\Omega_{\chi} h^2 > 1$. The transition from heavily wino-like LSP to heavily bino-like
Figure 5: **Thermal Relic Density in Deflected Mirage Mediation.** Left panel takes $M_0 = 500$ GeV, right panel takes $M_0 = 1000$ GeV. The “WMAP preferred” region of $0.07 \leq \Omega_\chi h^2 \leq 0.14$ is here indicated by the narrow red shaded region. For the wino-like LSP and within the $A$-funnel region the relic density drops below the critical value of $\Omega_\chi h^2 = 0.025$, indicated by the yellow (very light) shading. The remaining regions are $0.025 \leq \Omega_\chi h^2 \leq 0.07$ (blue), $0.14 \leq \Omega_\chi h^2 \leq 1$ (green) and $\Omega_\chi h^2 > 1$ (gray).

LSP occurs near $y \simeq 1.4$ where the relic density rapidly changes from far too low to far too high. In between these extremes there exists a narrow region with $\Omega_\chi h^2 \simeq 0.10$.

The precision with which the cosmological abundance of cold dark matter can be inferred from the cosmic microwave background is remarkable. Yet despite the narrow window of (31) we cannot say that all of this cold dark matter is composed of neutralinos, nor that it was produced by standard thermal mechanisms. Thus we must be careful not to immediately exclude those parameter sets for which the thermal calculation yields $\Omega_\chi h^2$ outside the range of (31). For example, an under-abundance of relic neutralinos may simply indicate the presence of non-thermal production mechanisms for stable LSPs [48, 49, 50]. Such mechanisms are especially well-motivated in string-inspired contexts such as in the KKLT model framework [51, 52]. An over-abundance of relic neutralinos is a greater cause for concern, but here too mechanisms exist which can bring the current relic density in line with the results from WMAP [53, 54, 55]. We prefer to remain agnostic on the issue as this is not the physics of interest to us. In fact, the physics of relic neutralino detection can
Table 2: **Parameter Scan Ranges for Deflected Mirage Mediation Models.** In an effort to average over possible model-dependent effects we have produced 1000 deflected mirage mediation models by randomly choosing input parameters from within the ranges indicated. An additional random choice of modular weights for the matter and Higgs multiplets was also made to complete the model.

| Parameter | Min Value | Max Value | Parameter | Min Value | Max Value |
|-----------|-----------|-----------|-----------|-----------|-----------|
| $M_0$     | 100 GeV   | 2000 GeV  | $\mu_{\text{mess}}$ | $10^8$ GeV | $10^{14}$ GeV |
| $\alpha_m$ | 0         | 5         | $\alpha_g$ | -5        | 5         |
| $\tan \beta$ | 2         | 40        | $N_m$ | 1         | 3         |

be separated from the cosmological density of neutralinos to the extent that it is only the *local* halo density profile that is relevant. Of course the two will be related once a set of assumptions about the physics of the early universe are specified. We will, however, make one exception: for cases with $\Omega_\chi h^2 \leq 0.025$ (indicated by the yellow shading in Figure 5) it becomes difficult for the relic particle in question to account adequately for the local halo density of our galaxy \[1\]. Therefore when calculating observable quantities that depend on the relic neutralino number density $n_\chi$ present in our galaxy (or the energy density $\rho_\chi = m_\chi n_\chi$) we will rescale the assumed local density of $(\rho_\chi)_0 = 0.3$ GeV/cm$^3$ by the multiplicative factor $r_\chi = \text{Min}(1, \Omega_\chi h^2/0.025)$.

Displaying results in the $\{x, y\}$ plane defined by (21) is useful for making contact with certain theoretical limits, but requires fixing a number of important input parameters. We will therefore also show results of a scan over parameter inputs for the deflected mirage mediation paradigm. The quantities we vary are listed in Table 2. Note that the input parameters $\Lambda_{\text{mess}}$ and $m_{3/2}$ are obtained from the definitions in (12) and (13), respectively. For the modular weights we randomly chose from the possibilities

$$
\begin{align*}
n_Q = n_U = n_D = n_L = n_E & = 0 \text{ or } 1/2 \\
n_{H_u} = n_{H_d} & = 0 \text{ or } 1/2 \text{ or } 1
\end{align*}
$$

which include the cases with the widest range of allowed points at the electroweak scale \[26\]. We generated 1000 points, all of which have proper electroweak symmetry breaking and satisfy the mass bounds imposed in Section \[2.3\]. For the sake of comparison we also generated 1000 points with $\alpha_m = \alpha_g = 0$ and vanishing modular weights. These points represent unified modulus-mediated models with a phenomenology similar to traditional dilaton-dominated models \[56, 57\].

We plot the thermal relic density $\Omega_\chi h^2$ as a function of the LSP mass for both the mSUGRA-like
Figure 6: **Relic Density for Models from Table 2**  
Thermal relic density $\Omega\chi h^2$ is plotted as a function of the neutralino LSP mass. Dark circles are the ensemble of deflected mirage mediation models defined by Table 2. Larger lighter circles are an ensemble of unified modulus-mediated models with $\alpha_m = \alpha_g = 0$. Red crosses are the DMM models which fall into the “WMAP Preferred” region $0.07 < \Omega\chi h^2 < 0.14$. Models in the shaded region below $\Omega h^2_{\text{min}} = 0.025$ have the local halo density rescaled before any further observable is calculated.

In case $\alpha_m = \alpha_g = 0$ and for the ensemble of deflected mirage mediation models defined by Table 2 in Figure 6, Dark circles are the ensemble of deflected mirage mediation models, while the larger, lighter circles are the cases with $\alpha_m = \alpha_g = 0$. Red crosses are the DMM models which fall into the “WMAP preferred” region $[2]$. Note that the WMAP-preferred area for the DMM model set tends to cluster around $m_{\tilde{N}_1} \sim \mathcal{O}(1 \text{ TeV})$, as opposed to the unified models which require a much lower-mass LSP. We will designate this group of models with red crosses in the figures which follow. Finally, we note that those models in the shaded region below $\Omega h^2_{\text{min}} = 0.025$ have the local halo density rescaled by the factor $r_\chi = \text{Min}(1, \Omega\chi h^2/0.025)$ before any further observable is calculated.
Table 3: Rate Estimates for Various Experiments. The minimum threshold rate $R_{10}$ necessary to produce 10 events in a given experiment for the fiducial mass and exposure time given is tabulated in the final column. Note that for Xenon10 and CDMS II we use the experimentally quoted fiducial masses. For all other (future) experiments we assume a fiducial mass equivalent to 80% of the nominal quoted target mass.

| Ref. | Experiment Name | Fiducial Mass [kg] | Exposure Time [yr] | $R_{10}$ [counts/(kg yr)] |
|------|-----------------|--------------------|--------------------|--------------------------|
| 63   | XENON10         | 5.4                | 0.16               | 11.54                    |
| 65   | XENON100        | $170 \times 0.8$   | 1                  | $7.35 \times 10^{-2}$    |
| 66   | LUX             | $350 \times 0.8$   | 3                  | $1.19 \times 10^{-2}$    |
| 65   | XENON1T         | $1000 \times 0.8$  | 5                  | $2.50 \times 10^{-3}$    |
| 62   | CDMS II         | 3.75               | 0.29               | 9.18                     |
| 64   | SuperCDMS (SNOlab) | $27 \times 0.8$  | 3                  | 0.15                     |
| 64   | SuperCDMS (DUSEL) | $1140 \times 0.8$ | 5                  | $2.19 \times 10^{-3}$    |

3.2 Direct Detection

Given the various theoretical considerations outlined above, it is clear that terrestrial direct detection experiments would provide the firmest evidence for the existence of weakly-interacting massive dark matter. The interaction rate of relic neutralinos with target nuclei depends on the local density of the neutralino in the local halo. This number is more tightly constrained by experimental observation [1, 58] and fits to large scale structure simulations of galaxies similar to ours. The rate depends heavily on the properties of the neutralino itself, but it also requires the input of certain nuclear matrix elements that must be inferred from experiment. The values of these matrix elements are uncertain – often to a surprisingly large degree. The resulting uncertainty in the interaction cross-section can be as much as 50% in some circumstances [59, 60]. Though our interest here is in the study of broad correlations between the parameters of the generalized mediation scenario with the properties of the relic neutralino, we should keep these uncertainties in mind when discussing the results of any calculation.

We will focus here exclusively on two types of detector design: cryogenic germanium bolometers and dual-phase liquid/gas xenon detectors. These types are currently operational in at least one experiment and producing results – and future enlargements are envisioned for each. We use DarkSUSY with default settings for all nuclear form factors to obtain the differential rate of interactions per unit recoil energy on germanium and xenon.\(^2\) We do so over a range of recoil energies

\(^2\)We have corrected a normalization error in the recoil rate routines and ensured that all results are rescaled (when
relevant to the desired experiment. As in [61], we perform the integration of these rates using two possible energy ranges

\[
\begin{align*}
\text{Xe} & : \quad 5 \text{ keV} \leq E_{\text{recoil}} \leq 25 \text{ keV} \\
\text{Ge} & : \quad 10 \text{ keV} \leq E_{\text{recoil}} \leq 100 \text{ keV}.
\end{align*}
\] (34)

Note that these rate calculations include both spin-independent and spin-dependent contributions.

Bounds on these recoil rates have been set by current direct search experiments such as CDMS II [62] and Xenon10 [63], using exposures of 316.4 kg-days on xenon and 397.8 kg-days on germanium, respectively. Expected backgrounds for both experiments are low, typically on the order of (or less than) 10 expected background events per year of exposure. Background rates at future enlargements are expected to be even lower. We will therefore compute the necessary interaction rates over the energy ranges in (34) to produce 10 events for a given mass and exposure time in a number of current or planned experiments based on xenon or germanium targets. We summarize the detector details we will assume in Table 3. For the current experiments a limit can be placed by requiring \( R \leq R_{10} \), the rate for which 10 events would have been produced at each experiment for the quoted exposure. These numbers are given in the final column of Table 3 and indicated in Figure 7 by the various shaded regions. We find that none of the model points give rates which should have been detected in the initial runs of Xenon10 and CDMS II.

For the future experiments listed in Table 3 we computed the value of \( R_{10} \) assuming a fiducial target mass equal to 80% of the quoted (nominal) mass and the exposure time listed. The results (in units of recoils per kg per year) for the \( \{x, y\} \) plane are plotted in Figure 7 for xenon and germanium targets, respectively. Results for our random model survey are given in Figure 8. The yellow (light shading) in all panels indicates parameter space that will be inaccessible to any of the future experiments in Table 3. At roughly the 100 kg-year level of exposure a fair amount of the parameter space will be detectable, particularly for low mass scales. This is indicated by the green (medium shading) in all panels, for which \( 11.54 > R_{10}^{\text{Xe}} \geq 0.0735 \) counts/kg-yr for xenon and \( 9.18 > R_{10}^{\text{Ge}} \geq 0.15 \) counts/kg-yr for germanium, respectively. This is within reach of the Xenon100 experiment after about one year of operation, or the SuperCDMS SNOLab proposal after about three years of operation. In about one ton-year of exposure – equivalent to about three years of the LUX experiment – the red (dark) shaded region in the xenon panels with \( 0.0735 > R_{10}^{\text{Xe}} \geq 0.0119 \) counts/kg-yr will be probed. Finally, after approximately five ton-years in xenon the blue shaded region with \( 0.0119 > R_{10}^{\text{Xe}} \geq 0.0025 \) counts/kg-yr can be explored, or the red shaded region with (necessary) by the thermal relic density of LSPs.
Figure 7: Neutralino Recoil Rates on Xenon (top) and Germanium (bottom). Left panels set $M_0 = 500$ GeV, right panels set $M_0 = 1000$ GeV. Phenomenologically allowed areas are enclosed by the heavy dashed lines. Colored shading indicates the reach of future direct detection experiments, as computed in Table 3. Yellow in all panels indicates parameter space that will be inaccessible to any of the future experiments in Table 3. For xenon targets, green indicates $11.54 > R_{Xe}^{10} \geq 0.0735$, red indicates $0.0735 > R_{Xe}^{10} \geq 0.0119$ and blue indicates $0.0119 > R_{Xe}^{10} \geq 0.0025$ in recoils per kg-year. For germanium targets, green indicates $9.18 > R_{Ge}^{10} \geq 0.15$ and red indicates $0.15 > R_{Ge}^{10} \geq 0.00219$ recoils per kg-year.
Reach of current and future experiments in terms of the $R_{10}$ values of Table 3 are given for xenon-based experiments (left panel) and germanium-based experiments (right panel). Dark circles are the ensemble of deflected mirage mediation models defined by Table 2. Larger lighter circles are an ensemble of unified modulus-mediated models with $\alpha_m = \alpha_g = 0$. Red crosses are the DMM models which fall into the “WMAP preferred” region. The xenon-based experiments are (a) Xenon10 (reported), (b) Xenon100 (1 year), (c) LUX (3 year) and (d) Xenon1T (5 year). The germanium-based experiments are (a) CDMS II (reported), (b) SuperCDMS (SNOLab - 3 year) and (c) SuperCDMS (DUSEL - 5 year).

$0.15 > R_{10}^{Ge} \geq 0.00219$ counts/kg-yr in germanium. Note that these regions are related to the wave-function of the neutralino LSP and to the presence of light Higgs states in the spectrum. In the $A$-funnel region where the relic density is well below the $\Omega_{\chi} h^2 = 0.025$ threshold the interaction rates have been rescaled, thereby requiring a much larger exposure for detection.

The results of the survey of randomized models are given in Figure 8. The current experimental limits are given by the horizontal lines labeled (a) in the figures. Though a few of our unified models would have been detected at these experiments, none of the deflected mirage models would have produced at least ten recoils. The five ton-year limit is indicated by the horizontal lines labeled (d) for xenon (with $R_{10}^{Xe} = 2.50 \times 10^{-3}$ counts/kg-yr) and (c) for germanium (with $R_{10}^{Ge} = 2.19 \times 10^{-3}$ counts/kg-yr). The majority of the points which satisfy the WMAP preferred bound of (32) can
Table 4: **Muon Flux Estimates for IceCube.** Flux of muons needed to produce 10 signal events at IceCube assuming various exposures.

be probed at this level of exposure.

### 3.3 Muons

The presence of relic neutralinos can also be inferred from experiments which seek to detect the products of neutralino annihilation processes. The rate for neutralino pair annihilation will be quite low except in those areas where the present-day relic density is high. For indirect detection searches to be successful, therefore, they must be sensitive to objects that (a) are able to travel unimpeded for long distances from their point of origin, and (b) are not themselves copiously produced by mundane astrophysical processes. For the second point, it may be sufficient to demand that some property of the object (such as the differential flux as a function of energy) be significantly different from estimates of the astrophysical backgrounds. Any calculation that seeks to predict these signals will be faced with uncertain theoretical inputs – particularly the density of neutralinos at the extra-terrestrial location of interest.

A good place to start, therefore, is at a location that is well-understood such as at the center of the sun or earth. Relic neutralinos can become gravitationally trapped in the core of the sun, significantly enhancing their probability for annihilation. Of the annihilation products, some fraction of the neutrinos can eventually exit the sun and be detected in experiments such as IceCube via conversion of muon neutrinos into muons. To calculate this rate we integrate the differential flux of conversion muons from solar-born as well as earth-born neutrinos over the energy range $50 \text{ GeV} \leq E_\mu \leq 300 \text{ GeV}$, assuming an angular resolution of 3 degrees. The nominal target area for IceCube is 1 km$^2$, but the effective area for detection of neutrinos via muon conversion is smaller and can be cast as a function of the muon energy. To determine the prospects for observing muons at IceCube we estimate various values of exposure in km$^2$-years needed for 10 signal events, $\Phi_{10}$, as given in Table 4. We note that the estimated background of muon events in one
Figure 9: **Upward Going Muon Rates at IceCube.** Left panel shows the IceCube reach for $M_0 = 500$ GeV. The red shaded region is detectable with 0.2 km$^2$-years of exposure. Green and blue shaded regions would require 0.5 and 1.5 km$^2$-years of exposure, respectively. The yellow region gives a visible signal after 10 km$^2$-years of exposure. The gray region has a flux below 1 muon per km$^2$-yr and is likely undetectable at IceCube. The right panel gives the muon flux at IceCube for the random models defined by Table 2. Points above the horizontal line at $\Phi_\mu = 400$ muons/km$^2$-year would likely have given a signal in the 22-string IceCube data. The ultimate reach for IceCube is approximated by the lowest horizontal line at $\Phi_\mu = 1$ muon/km$^2$-year.

A year of IceCube data taking is expected to be $\mathcal{O}(10)$. In Figure 9, we show the number of observed muons at IceCube in units of counts per km$^2$-yr, combining those from neutralino annihilation in the sun with those in the center of the earth. The left panel shows the regions in the $\{x, y\}$ plane satisfying the condition $\Phi \geq \Phi_{10}$ for the values listed in Table 2 for the case when $M_0 = 500$ GeV. For higher values of $M_0$ the (already low) muon rates drop rapidly and little of the parameter space would be visible at IceCube. This is evident in the right panel, which shows the muon rate in units of muons/km$^2$-year for the random model set of Table 2. Recent results using a 22-string detector at IceCube constrain the muon flux to be below 300-500 muons per km$^2$-yr. None of the points in the left panel of Figure 9 produce a flux this large, though a few of the randomly-generated points would have produced a signal at IceCube. We conservatively place the limit at 400 muons per km$^2$-yr, indicated by the horizontal line.
For low values of the mass scale $M_0$ there are some points which can produce a flux of 50 muons per km$^2$-yr. These points are indicated in the left panel Figure 9 by the red shaded region. The three smaller flux values of Table 9 are indicated by the green, blue and yellow regions, respectively. This is roughly congruent with the parameter space that gives rise to a mixed Higgsino/gaugino LSP and a light supersymmetric Higgs sector. For the case where $M_0 = 1000$ GeV and there is very little region with a mixed LSP we expect the flux of muons at IceCube to be well below 1 muon per km$^2$-yr throughout most of the parameter space. The last two lines in Table 4 are indicated by the lower two horizontal lines in the right panel of Figure 9. Though the DMM models tend to give a larger flux of muons at IceCube than unified models with an equivalent LSP mass, the prospects for a visible signal at IceCube remain limited for most of the parameter space.

3.4 Photons

At distances further from the earth one can search for photons in the gamma ray energy regime, which travel largely unimpeded from their source. This allows gamma ray observatories to concentrate on areas of the sky likely to have a high relic neutralino density – such as the center of our galaxy. Photons can be produced as part of decay chains or directly through loop-induced diagrams. The former contribute to a continuous general spectrum of photons and are thus more difficult to distinguish from astrophysical backgrounds. Direct production of photon pairs – or production of a single photon in association with a Z-boson – offers the possibility of a monochromatic spectrum that can more easily be distinguished from the background at the cost of a reduction in rate relative to the continuous photon rate. We will consider searches for both phenomena in this subsection.

The calculation of the flux of gamma rays observed from the direction of the galactic center depends on the microscopic physics of the neutralino and its interactions, but also very strongly on the macroscopic physics of the halo profile assumed for the galaxy. The latter is conveniently summarized by a single parameter $\bar{J}(\Delta \Omega)$ where $\Delta \Omega$ represents the solid angle resolution of the observatory. Here we will specifically consider the Fermi/GLAST experiment for which we assume an angular resolution of $\Delta \Omega = 10^{-5}$ sr. We will consider the commonly adopted profile of Navarro, Frank and White (NFW) [72] which gives a value of $\bar{J}(\Delta \Omega) = 1.2644 \times 10^4$, as well as a modified NFW profile that takes into account the effect of baryons (a so-called “adiabatic compression” (AC) model) [73, 74] for which we have $\bar{J}(\Delta \Omega) = 1.0237 \times 10^6$. The large variation in calculated rates for different halo profile models should always be borne in mind when discussing “predictions”
Table 5: Reach Estimates in the Continuous Gamma Ray Flux for the Fermi/GLAST Experiment.

The quantity \( \Phi_{100} \) is the flux needed for 100 signal events at Fermi/GLAST assuming various exposure times and halo profiles.

| Exposure [m^2 yr] | Halo Profile | \( \Phi_{100} \) [counts/(cm^2 s)] |
|-------------------|--------------|-----------------------------------|
| 1                 | NFW          | 5.79 \times 10^{-10}              |
| 5                 | NFW          | 1.16 \times 10^{-10}              |
| 1                 | NFW+AC       | 7.15 \times 10^{-12}              |
| 5                 | NFW+AC       | 1.43 \times 10^{-12}              |

for gamma ray signals of any particular model. Nevertheless, we can study the dependence of the calculated rate on the underlying parameters and scale the normalization of that rate to any particular halo profile via the ratio of \( \bar{J}(\Delta \Omega) \) values, as needed.

For the diffuse signals we use DarkSUSY to compute the differential photon flux from neutralino annihilation in the halo in units of photons/cm^2/s/GeV. We do so over the photon energy range relevant for the Fermi/GLAST experiment of 1 GeV \( \leq E_\gamma \leq 200 \) GeV, with a step size of 1 GeV. Note that the upper limit on the Fermi detector sensitivity is typically well below the mass of the neutralino LSP for much of the parameter space we are studying. We then numerically integrate these values to obtain an overall photon flux in units of photons/cm^2/sec. Using the effective area of 1 m^2 for the Fermi detector we compute the minimum flux \( \Phi_{100} \) necessary to produce an event rate of 100 signal photons from the galactic center in a given length of exposure. As an example, using the NFW profile one obtains a value of \( \Phi_{100} = 5.79 \times 10^{-10} \) photons/cm^2/sec for one year of running, where we assume data collection for 200 days per year. If we require 100 signal photons over a five year mission this becomes \( \Phi_{100} = 1.16 \times 10^{-10} \) photons/cm^2/sec. To take into account the possibility of a more generous halo profile such as the NFW profile with adiabatic compression, we can simply rescale this reach by the ratio of \( \bar{J}(\Delta \Omega) \) for the two profiles. In this case this implies a reach to lower fluxes by a factor of approximately 80. The corresponding \( \Phi_{100} \) values are listed in Table 5.

Translating this into a discovery reach for Fermi/GLAST requires making an assumption about the background sources for gamma rays in the 1 GeV \( \leq E_\gamma \leq 200 \) GeV energy regime. We will use the estimate of [58] which gives a background flux of \( \Phi_{\text{bkgnd}} = 5.06 \times 10^{-10} \) photons/cm^2/sec, or \( \mathcal{O}(100) \) photons/year for the aperture and energy range of Fermi/GLAST. This coincides with sensitivity measures on the photon flux at Fermi/GLAST often quoted in the literature [75, 76]. In
Figure 10: **Reach of the Fermi/GLAST Experiment for Integrated Gamma Ray Flux.** Left panel is for $M_0 = 500$ GeV, right panel is for $M_0 = 1000$ GeV. Phenomenologically allowed areas are enclosed by the heavy dashed lines. For $M_0 = 500$ GeV the green shaded region has $5.79 \times 10^{-10} > \Phi_{100} \geq 1.16 \times 10^{-10}$ photons/cm$^2$/sec and therefore may be visible at Fermi/GLAST after five years of exposure with the NFW profile. If the fluxes are rescaled to the NFW+AC profile we obtain the reach shaded in blue (one year of data-taking) and yellow (five years of data-taking). The gray shaded area has an effective flux below $\Phi_\gamma = 0.014 \times 10^{-10}$ and is unlikely to yield a visible signal for diffuse gamma rays at Fermi/GLAST.

In addition to looking for excesses in the overall rate, the Fermi experiment will also be able to use the shape of the observed spectrum to augment signal extraction. Furthermore, there is reason to believe that point sources near the galactic center can be identified and subtracted from the signal to further enhance signal significance [77]. We will therefore optimistically take the calculated flux $\Phi_{100}$ as a measure of the reach of the Fermi experiment in the deflected mirage mediation model space.

The result of this reach analysis is given in Figure 10 for the cases of $M_0 = 500$ GeV (left panel) and $M_0 = 1000$ GeV (right panel). Detecting a gamma ray signal in the deflected mirage mediation model will be challenging, even with a relatively favorable set of assumptions about the halo profile. For the low mass case of $M_0 = 500$ GeV there is a small region near the $A$-funnel (green shading) for which a signal is possible after five years of data-taking at Fermi/GLAST with the NFW profile. This coincides roughly with the region with a mixed Higgsino/gaugino LSP. No such region exists
in the higher mass $M_0 = 1000$ GeV case. The blue and yellow regions in both panels are the reach at Fermi/GLAST after one year and five years, respectively, if the flux calculation was rescaled to the NFW + AC profile. Roughly speaking, with this set of halo assumptions it should be possible to detect a diffuse gamma ray signal over most of the parameter space involving a wino-like or mixed Higgsino/gaugino LSP. The gray area in the figures corresponds to the area with a bino-like LSP and typically has an effective flux below $\Phi_\gamma = 0.014 \times 10^{-10}$ photons/cm$^2$/sec. These regions are unlikely to yield visible signals without additional favorable assumptions about the halo model. We note, however, that these rates would increase considerably if we were to assume a local halo density normalized to $\rho_0 = 0.3$ GeV/cm$^3$ regardless of the calculated thermal relic density.

The best prospects for diffuse gamma ray signals occur near regions that are ruled-out by
Table 6: **Reach Estimates in the Monochromatic Gamma Ray Flux for a Generic ACT Experiment.**
Flux needed for 10 signal events at a generic ACT assuming various exposures and halo profiles. Note that we are here taking one year of data-taking to mean 365 days.

| Exposure [m² yr] | Halo Profile | Φ₁₀ [counts/(cm² s)] |
|-----------------|--------------|---------------------|
| 1000            | NFW          | 31.7 × 10⁻¹⁵        |
| 100             | NFW+AC       | 3.92 × 10⁻¹⁵        |
| 500             | NFW+AC       | 7.83 × 10⁻¹⁶        |
| 1000            | NFW+AC       | 3.92 × 10⁻¹⁶        |

certain scalars becoming the lightest supersymmetric particle. This, in turn, is a property of the modular weights assumed for the scalar sector of the theory. In Figure 11 we show the results for the integrated gamma ray flux in the range $1 \text{ GeV} \leq E_\gamma \leq 200 \text{ GeV}$ for the random collection of models defined by Table 2 and modular weights (33). Choosing different modular weights does not produce a large number of models with signals visible at Fermi/GLAST using the NFW profile. The one-year and five-year $\Phi_{100}$ reach values are indicated by the solid horizontal lines. More encouragingly, the models that tended to satisfy the WMAP relic density constraints predicted values of $\Phi_{100}$ that could potentially give observable signals at Fermi/GLAST if the NFW + AC profile were applicable.

For the case of monochromatic signals, $\chi\chi \rightarrow \gamma\gamma$ or $\chi\chi \rightarrow Z\gamma$, we obtain the flux at a single energy, $E_\gamma = m_\chi$ or $E_\gamma = m_\chi - \frac{m_\gamma^2}{4m_\chi}$, respectively. As we have seen previously, for the deflected mirage mediation scenario the neutralino LSP can often be quite massive, particularly for the cases that satisfy the WMAP thermal relic density constraint. These monochromatic gamma ray signals are therefore likely to be at energies beyond the reach of the orbiting Fermi/GLAST satellite experiment. On the other hand, ground based atmospheric Cherenkov telescopes (ACTs) such as CANGAROO [78], HESS [79], MAGIC [80] and VERITAS [81] have thresholds for photon detection in the 100 GeV range and can detect energetic photons up to $\mathcal{O}(10 \text{ TeV})$. Though data-taking is restricted to dark, cloudless nights the effective area $A_{\text{eff}}$ of the target detector is quite large.

Unlike the case for continuum photons, the monochromatic signal has virtually no background, particularly at high photon energies. We therefore calculate the minimum flux $\Phi_{10}$ to produce 10 photons for a given exposure area $\times$ time, within the window defined by the energy resolution of a typical ACT. We take this energy resolution to be $\sigma_E/E = 0.15$. In Figure 12 we present the reach for a generic ACT with $A_{\text{eff}} \simeq 3 \times 10^8 \text{ cm}^2$ using the estimates for $\Phi_{10}$ given in Table 6.
Figure 12: Reach of the Fermi/GLAST Experiment for Monochromatic Gamma Ray Flux. Left panel is for $M_0 = 500$ GeV, right panel is for $M_0 = 1000$ GeV. Red shaded regions near the edge of the allowed area in the two plots have $\Phi_{10} \geq 31.7 \times 10^{-15}$ photons/cm$^2$/sec. Rescaled to the NFW+AC profile we obtain the reach shaded in green (100 m$^2$-yr exposure), blue (500 m$^2$-yr exposure) and yellow (1000 m$^2$-yr exposure). Note that some of the indicated regions of the parameter space give photon energies below the lower threshold of $E_\gamma = 100$ GeV and are therefore invisible to our generic ACT. The gray shaded region has a monochromatic flux $\Phi_{10} \leq 0.39 \times 10^{-15}$ photons/cm$^2$/sec and is unlikely to give a visible signal without a more favorable set of halo assumptions.

we have not included models with a neutralino LSP mass below 100 GeV, as the ACTs do not typically take data below this energy. As in Table 5 we convert the quantity $\Phi_{10}$ into an estimate of the reach of our generic ACT over the model space for the NFW profile as well as for the NFW + AC profile. Here we combine the photons from both the $\gamma\gamma$ and $\gamma Z$ signals as the 15% energy resolution will not be sufficient to resolve the two lines for LSP masses much in excess of 200 GeV.

As in the case with the continuous gamma ray spectrum, detecting a signal from neutralino annihilation will be difficult without help from the halo profile. The red areas in the two plots have $\Phi_{10} \geq 31.7 \times 10^{-15}$ photons/cm$^2$/sec and are potentially visible with the NFW profile given 1000 m$^2$-years of exposure. Note that the heavier mass cases ($M_0 = 1000$ GeV) are easier to detect due to the sensitivity of ACTs to very energetic photons. If the monochromatic fluxes are scaled to the values that one would obtain if we had used the NFW + AC profile then the area in green, blue and yellow shading (corresponding to $31.7 \times 10^{-15} > \Phi_{10} \geq 3.91 \times 10^{-15}$ photons/cm$^2$/sec,
Figure 13: **Monochromatic Gamma Ray Flux for Models in Table 2**  Dark circles are the ensemble of deflected mirage mediation models defined by Table 2. Larger lighter circles are an ensemble of unified modulus-mediated models with $\alpha_m = \alpha_g = 0$. Red crosses are the DMM models which fall into the “WMAP preferred” region (32). The solid line indicates a generic ACT reach ($\Phi_{10}$) after 1000 m$^2$-years of exposure. The dashed lines indicate the reach after rescaling all fluxes using the NFW + AC profile for 100, 500 and 1000 m$^2$-years of exposure.

$$3.91 \times 10^{-15} > \Phi_{10} \geq 0.78 \times 10^{-15} \text{ photons/cm}^2/\text{sec} \quad \text{and} \quad 0.78 \times 10^{-15} > \Phi_{10} \geq 0.39 \times 10^{-15} \text{ photons/cm}^2/\text{sec},$$

respectively) would be visible after 1000 m$^2$-years of exposure. This roughly corresponds to the area of wino and mixed Higgsino/gaugino LSP.

In Figure 13 we show the results for the monochromatic gamma ray flux for the random collection of models defined by Table 2 and modular weights (33). The reach $\Phi_{10}$ for a generic ACT after 1000 m$^2$-years of exposure is given by the solid horizontal line. Dashed horizontal lines give the reach after 100, 500 and 1000 m$^2$-years of exposure, respectively, when all values of the flux are scaled up using the NFW + AC profile. The rates of annihilation through loop diagrams into mono-energetic photons are sensitive to the wave-function of the LSP. Hence the deflected mirage mediation models break into groups in Figure 13 determined by the makeup of the LSP. The unified
models, on the other hand, are nearly universal in having a bino-like LSP. We will return to this feature in our discussion in Section 4. As in Figure 11 we see that the WMAP preferred models are visible after a reasonable amount of ACT exposure provided the halo profile has a sufficiently high concentration at the galactic center.

3.5 Anti-matter in Cosmic Rays

Dark matter annihilations will also produce charged particles, of which positrons and anti-protons are the most promising for giving signals above astrophysical backgrounds. There has been experimental evidence for an excess of positrons in the 10-100 GeV range for some time [82, 83]. The recent data from the PAMELA experiment seems to confirm the previous results [84]. It is by now well appreciated that standard MSSM models can fit the signal from PAMELA only with some difficulty, and are even more challenged by the possible positron excess at even higher energies reported by the ATIC experiment [85].

For that reason we will here only concern ourselves with the positron and anti-proton energy ranges relevant for PAMELA, roughly a few GeV up to 100 GeV.

The excess over background was reported at higher energies, so we therefore consider the ten highest energy bins from the reported signal. We will take as the central value of these bins the values

\[ E_i = \{5.55, 6.75, 8.25, 10.1, 13.05, 17.45, 23.85, 34.75, 53.2, 82.15\} \] (35)

for the energies (in GeV) of the electron or positron cosmic ray. Using DarkSUSY 5.0.4 (with all galactic propagation parameters set to their default values) we calculate the differential positron flux from neutralino annihilation with the NFW halo profile in units of positrons/GeV/cm\(^2\)/sec/sr and obtain an average flux in each bin, denoted \(\Phi_{\text{sig}}^{+e}\). We then form the ratio of positron flux to positron and electron flux as a function of the “boost factor” \(B\) via the relation

\[ f_i(B) = \frac{\Phi_{\text{blnd}}^{+e} + B\Phi_{\text{sig}}^{+e}}{\Phi_{\text{blnd}}^{+e} + B\Phi_{\text{sig}}^{+e} + \Phi_{\text{prim}}^{e-} + \Phi_{\text{sec}}^{e-}}, \] (36)

where \(\Phi_{\text{blnd}}^{+e}\) is the background positron flux and \(\Phi_{\text{prim}}^{e-}\), \(\Phi_{\text{sec}}^{e-}\) are the contributions to the electron flux from primary and secondary sources, respectively. Following [89] [90] we use analytic formulæ to approximate these contributions from astrophysical processes

\[ \Phi_{\text{blnd}}^{+e} = \frac{4.5E^{0.7}}{1 + 650E^{2.3} + 1500E^{1.2}} \] (37)

\(^3\)As evidence for this statement one can consider the analysis performed in [80] [87] [88], among many other possible studies.
\[ \Phi_{e^{-}}^{\text{prim}} = \frac{0.16E^{-1.1}}{1 + 11E^{0.9} + 3.2E^{2.15}} \]
\[ \Phi_{e^{-}}^{\text{prim}} = \frac{0.7E^{0.7}}{1 + 110E^{1.5} + 600E^{2.9} + 580E^{4.2}} \]

where \( E \) represents the particle energy (normalized to units of GeV).

To quantify how well a particular point in the parameter space of the DMM model can fit the PAMELA data, we use the function \( f(B) \) evaluated at each of the energy values in (35). The resulting quantities \( f_i(B) \) are then compared to the observed PAMELA data points \( f_i^{\text{PAM}} \) using a chi-squared like function of the boost factor

\[ \chi^2(B) = \sum_i \left( \frac{f_i^{\text{PAM}} - f_i(B)}{\sigma_i} \right)^2, \]

where the error values, \( \sigma_i \), are taken to be the average value of the full vertical error bar from the reported signal. We then determine the value of the boost factor \( B \) for which \( d\chi^2/db = 0 \) and (40) is minimized.

The results for the scan in the \( \{x, y\} \) plane for \( M_0 = 500 \) GeV (left panels) and \( M_0 = 1000 \) GeV (right panels) are shown in Figure 14. The values of the boost factors that we obtain from minimizing \( \chi^2(B) \) range from the relatively innocuous \( O(10) \) to the preposterous \( O(10^9) \). The largest boost factors occur for the heavy bino-like LSP in the red shaded regions of the upper panels. The smallest boost factors are denoted by the yellow and purple shaded regions and occur for the wino-like LSP region (both panels) and the mixed Higgsino/neutralino LSP region (\( M_0 = 500 \) GeV). Note that for these cases a large part of what is here termed a “boost factor” is in fact simply the result of a low thermal relic density. Invoking a non-thermal mechanism for explaining the present cosmological density of such wino-like LSPs would therefore result in \( O(1) \) boost factors for many of the models in this part of the parameter space. This is consistent with previous results suggesting that a low mass wino-like LSP is the best candidate for an MSSM explanation for the PAMELA data [91].

The heavy bino-like LSP region fit the PAMELA data extremely well (purple and yellow shaded regions in the lower panels of Figure 14), but this is merely an artefact of the LSP being very massive – the PAMELA signal is then interpreted as a (massively-boosted) tail with a peak at a much higher energy level. More realistic is the wino-like LSP region. Though the boost factors were relatively small here, the calculated flux was not a particularly good fit to the shape of the resulting positron spectrum. Typical values of the minimized \( \chi^2(B) \) were near unity (the blue shaded region in the lower panels of Figure 14), but the \( f_i(B) \) values tended to fall with increasing energy like
Figure 14: Best-Fit Boost Factors and Minimized $\chi^2(B)$ Values for Fitting Anomalous Positron Data. Left panels set $M_0 = 500$ GeV, right panels set $M_0 = 1000$ GeV. Phenomenologically allowed areas are enclosed by the heavy dashed lines. In the upper panels the regions are separated according to best-fit boost values as follows: $B \geq 10^8$ red, $10^8 > B \geq 10^6$ green, $10^6 > B \geq 10^4$ blue, $10^4 > B \geq 10^2$ yellow, and $B < 10^2$ in purple. In the lower panels the regions are separated according to the minimized $\chi^2_{\text{min}}(B)$ values as follows: $\chi^2_{\text{min}}(B) \geq 2$ red, $2 > \chi^2_{\text{min}}(B) \geq 1.5$ green, $1.5 > \chi^2_{\text{min}}(B) \geq 1$ blue, $1 > \chi^2_{\text{min}}(B) \geq 0.5$ yellow, and $\chi^2_{\text{min}}(B) < 0.5$ in purple.
the background, as opposed to showing an upward trend as seen by the PAMELA experiment. Nevertheless, when certain reasonable modifications to the cosmic ray propagation parameters in DarkSUSY are made, the shape of the spectrum can be made to fit the observations [92].

The PAMELA experiment has also reported data on anti-protons in cosmic rays which were consistent with expectations from background astrophysical sources as well as previous experimental data [93]. To compare the deflected mirage mediation parameter space to this data we compute the differential flux of anti-protons using DarkSUSY with the NFW profile and default diffusion parameters for the ten highest bins reported in [93], corresponding to anti-proton mean kinetic energies of 5.85, 6.98, 8.37, 10.1, 12.3, 15.3, 19.5, 25.9, 37.3 and 61.2 GeV. We use the estimations of Cirelli et al. [94] for the background flux from astrophysical processes. We consider a model point to be in conflict with the PAMELA data if the predicted signal flux in any one of the ten bins were twice as high as the reported experimental observation. None of the parameter space gives rise to an anti-proton flux of this magnitude. If, on the other hand, we were to naively apply the same boost factor to the anti-proton flux as we do to the positron flux then the majority of these models would be in conflict with the PAMELA anti-proton data. While anti-protons and positrons will propagate through the galaxy differently [95], it is reasonable to assume that the fluxes will involve similar astrophysical backgrounds. Thus the two sets of data appear to be in conflict throughout the DMM parameter space unless modifications to the diffusion model from the DarkSUSY default values are made.

### 3.6 Benchmarks

Before concluding this section we will look back at the dark matter signatures for the specific benchmarks outlined in Table 1 in Section 2.3. To remind the reader, models A, B and C are specific examples presented in the original papers on deflected mirage mediation [29, 30]. Models D and E are chosen from the results of our parameter space survey, while model F is a point without gauge-charged messengers with \( \alpha_m = 1 \) as in the KKLT model. The properties of the neutralino LSP, the masses of other key particles, and the values for the thermal relic density and other dark matter observables are given in Table 7.

Properties of the LSP are given in the first block of entries in Table 7. The wave-function of the LSP is parameterized as

\[
\tilde{N}_1 = N_{11}\tilde{B} + N_{12}\tilde{W} + N_{13}\tilde{H}_d^0 + N_{14}\tilde{H}_u^0,
\]

which is normalized to \( N_{11}^2 + N_{12}^2 + N_{13}^2 + N_{14}^2 = 1 \). The properties of the NLSP and key Higgs
Table 7: Characteristics of Benchmark Models From Table 1

|                  | Model A | Model B | Model C | Model D | Model E | Model F |
|------------------|---------|---------|---------|---------|---------|---------|
| $m_{\tilde{N}_1}$ | 1009    | 626.6   | 772.8   | 257.2   | 193.2   | 710.8   |
| $f_B$            | 0.912   | 0.989   | 0.983   | 0.914   | 0.000   | 0.932   |
| $f_W$            | 0.071   | 0.001   | 0.001   | 0.030   | 0.997   | 0.004   |
| $f_H$            | 0.017   | 0.009   | 0.017   | 0.057   | 0.003   | 0.064   |
| NLSP             | $C_1$   | $C_1$   | $\tilde{\tau}$ | $C_1$   | $C_1$   | $\tilde{\tau}$ |
| $m_{NLSP}$       | 1026    | 743.8   | 781.3   | 302.6   | 193.3   | 762.8   |
| $m_B$            | 119.5   | 120.6   | 121.5   | 112.1   | 119.8   | 119.9   |
| $m_A$            | 1559    | 1043    | 1136    | 473.5   | 1518    | 1136    |
| $\Omega h^2$    | 0.117   | 0.891   | 0.517   | 0.088   | 0.001   | 0.537   |
| $R_{Xe}$         | 0.046   | 0.022   | 0.032   | 1.182   | 7 $\times$ 10^{-5} | 0.178   |
| $R_{Go}$         | 0.037   | 0.018   | 0.025   | 0.881   | 6 $\times$ 10^{-5} | 0.140   |
| $\Phi_{\gamma\gamma}$ | $1.778 \times 10^{-17}$ | $3.517 \times 10^{-17}$ | $2.734 \times 10^{-17}$ | $4.128 \times 10^{-16}$ | $2.071 \times 10^{-15}$ | $2.267 \times 10^{-17}$ |
| $\Phi_{\gamma Z}$ | $9.169 \times 10^{-17}$ | $3.244 \times 10^{-18}$ | $3.326 \times 10^{-18}$ | $8.868 \times 10^{-16}$ | $6.399 \times 10^{-15}$ | $1.157 \times 10^{-17}$ |
| $\Phi_{tot}$    | $1.846 \times 10^{-12}$ | $2.260 \times 10^{-12}$ | $2.041 \times 10^{-12}$ | $1.449 \times 10^{-10}$ | $1.463 \times 10^{-11}$ | $4.319 \times 10^{-12}$ |
| $\Phi_p$        | 0.738   | 0.194   | 0.495   | 11.766  | 0.016   | 3.616   |
| $\Phi_D$        | $8.096 \times 10^{-15}$ | $2.155 \times 10^{-14}$ | $1.195 \times 10^{-14}$ | $1.631 \times 10^{-12}$ | $2.824 \times 10^{-14}$ | $2.560 \times 10^{-14}$ |
| $\Phi_\mu$      | $7.438 \times 10^{-11}$ | $9.979 \times 10^{-11}$ | $8.923 \times 10^{-11}$ | $6.523 \times 10^{-9}$ | $8.617 \times 10^{-10}$ | $1.887 \times 10^{-10}$ |
| $B$              | 4774    | 5156    | 5145    | 102.9   | 1147    | 2510    |

All masses are given in GeV and we denote wave-function components as $f_B = |N_{11}|^2$, $f_W = |N_{12}|^2$ and $f_H = |N_{13}|^2 + |N_{14}|^2$. Direct detection rates $R_i$ have units of recoils/kg-year. Photon fluxes are given in units of photons/cm²/sec and $\Phi_{\gamma\gamma}$ represents the diffuse gamma flux integrated from 1 to 200 GeV. The muon flux is given in units of muons/km²/year. $\Phi_D$ is given in units of anti-deuterons/(GeV s cm² sr) and is computed at 0.25 GeV. $\Phi_\mu$ is given in units of anti-protons/(GeV s cm² sr) and is computed at 10 GeV.

masses are given in the next block of entries. These can have a significant impact on the thermal relic density calculation, which is provided in the next group of entries. All of the benchmark models have a thermal relic density that is below unity, though only models A and D agree well with the WMAP results of (31). In computing the relic density values all the models receive some reduction in abundance due to coannihilation. This is particularly true for points A, C and E. For this last case the LSP is mostly wino-like and the computed relic density is well below the WMAP bound. It is also below the threshold value of $\Omega h^2_{\text{min}} = 0.025$ and thus all subsequent dark matter signals have been rescaled to take this into account. Model D is an example of a point which resides in the $A$-funnel region where $2m_{\tilde{N}_1} \simeq m_A$ which also has effects on other dark matter observables.

Apart from the wino-like model E, all of these cases would be probed by direct detection.
experiments at the one-ton level, and most would give a sizable signal for detectors in the 100-300 kilogram range. As previously mentioned, one should keep in mind that these rate values do come with large theoretical uncertainties associated with the nuclear matrix element factors which go into the calculation. Models D and F have the largest Higgsino content and hence the largest direct detection rates, with model D giving the most sizable rate due to its low mass LSP. We point out that the very low rate for model E is in part due to the rescaling of the local halo density we performed. If we had simply assumed a local halo density normalized to 0.3 GeV/cm$^3$ these rates would be increased by a factor of 25.

Gamma ray signals for the benchmarks are shown using the NFW halo profile. For the continuous gamma ray signal we estimate that a flux on the order of $1-5 \times 10^{-10}$ photons/cm$^2$/sec in the Fermi/GLAST photon energy range of 1 to 200 GeV is needed to observe a signal over the astrophysical background. Only model D gives rise to such a large flux of gamma rays. Each of the others, however, would be observable above background if the NFW profile with adiabatic compression were employed. The wino-like LSP of model E gives a large signal for the two monochromatic gamma ray signals. Such a (combined) signal is just above the threshold for most ACT experiments and would give a visible signal with the NFW profile for 5000 m$^2$-years of exposure – not an unreasonable amount given the size of most atmospheric Cherenkov detectors. For the NFW + AC profile it is likely that only models A, D and E will give rise to a detectable monochromatic gamma ray signal. Expectations for observation of these benchmarks at IceCube via the flux of muons from neutralino annihilation in the sun are also a bit mixed. Models D and F have the highest muon fluxes and ought to produce muons at rates above background for reasonable estimates of the effective area of the IceCube detector at these energies. The remaining models are unlikely to produce a signal even after 10 km$^2$-years of data-taking at IceCube.

Anti-matter prospects fare slightly worse across our benchmark models. In all cases the naive estimate of anti-proton fluxes is well below the background expectation of $\mathcal{O}(10^{-7}) \bar{p}/(\text{GeV s cm}^2 \text{ sr})$ at our reference point of 10 GeV of kinetic energy for the anti-proton. If one wishes to explain the PAMELA positron data using these models, however, one is confronted with rather large boost factors for all models save model D. Naively applying such boost factors to the anti-proton flux would place all of these models in tension with the anti-proton data, giving a signal flux of roughly an order of magnitude above the background. For completeness we have also included the flux of anti-deuterons computed at a reference point of 0.25 GeV for the $\bar{D}$ kinetic energy which is relevant for the planned GAPS experiment \cite{96}. Only model D again is likely to give a sizable signal in this indirect detection channel.
4 Mirage Mediation versus Deflection Mirage Mediation

In the previous section we looked at how a number of astroparticle/cosmological signals depend on the parameters of the deflected mirage mediation paradigm. We did this in terms of the theory motivated \( \{x, y\} \) parameter space where we saw that the detection prospects are good for general values of the parameter \( x \) when \( 0.4 \lesssim y \lesssim 1.4 \) for the modular weight choice \([30]\). We also averaged over the other more model-dependent parameters in the theory, presenting our results in the form of scatter plots, using a unified mSUGRA-like theory as a foil with which to compare the predictions of the DMM model. That DMM and unified models give different “footprints” should not be a surprise, since the importance of non-universality in gaugino masses for dark matter observables has been emphasized for some time \([97, 98]\). More interesting is the question of whether observations in the dark matter arena are capable of distinguished DMM from the (un-deflected) mirage pattern which preceded it.

Looking back at the figures from the previous section we might conclude that the answer is certainly affirmative. After all, when we move away from the mirage prediction \( x + y = 1 \) we see strong variation in both the mass and composition of the LSP, resulting in dramatic changes in the thermal relic density of the neutralino. Moving away from the \( x + y = 1 \) limit can increase recoil rates at direct detection experiments and change the strength of signals at indirect detection experiments. But this is somewhat deceptive, since in those figures a number of important parameters – particularly the overall scale \( M_0 \) – were held constant. When we average over these values the distinction is far less clear. To exemplify this, consider Figure 15 in which we look at two of the more commonly considered observables in the literature: recoil rates on xenon targets and the integrated photon flux above 1 GeV in energy from the direction of the galactic center. In this figure we have fixed the value of \( M_0 = 500 \) GeV. The dark points continue to represent a random collection of 1000 deflected mirage mediation models with parameters taken from the ranges given in Table 2. The lighter shaded points are now no longer a unified model, but instead represent an additional 1000 points with the same ranges of Table 2 except that we fix \( N_m = 0 \). These are, therefore, mirage models with a similar overall mass scale. Though the range in LSP mass is more restricted in the un-deflected case, the signature range expected is nearly identical. These examples could be replaced with many of the signals we considered in Section 3. Such “inclusive” style measurements in the cosmological arena will not in themselves be able to determine the presence of the gauge mediation component in the underlying theory.

The reason for this degeneracy is readily found by considering another way to parameterize the
Figure 15: Comparison of Deflected Mirage Mediation with the Case of No Messengers. A random sample of models based on the ranges of Table 2 with \( M_0 \) fixed at 500 GeV for the deflected mirage case (dark points) versus the same ranges but with \( N_m = 0 \) (light points). The left plot gives the recoil rate per kg-year on xenon targets with sensitivities for Xenon100 after one year (solid line), LUX after three years (dashed line), and Xenon1T after five years (dotted line). The right figure gives the photon flux for all energies above 1 GeV from the galactic center. The solid line is the reach of Fermi/GLAST after five years with the NFW profile while the dashed line is the five year reach with the NFW + AC profile.

gaugino masses in the DMM model. We can translate each of the instances of the running gauge coupling \( g_a^2(\mu) \) in (11) to the high scale \( \mu_{\text{UV}} = \mu_{\text{GUT}} \) via the relation

\[
g_a^2(\mu_{\text{UV}}) = \frac{g_a^2(\mu)}{1 - \frac{g_a^2(\mu)}{8\pi^2} \ln \left( \frac{\mu_{\text{GUT}}}{\mu_{\text{mess}}} \right)^b \frac{\mu_{\text{GUT}}}{\mu_{\text{mess}}} N_m}.
\]

The gaugino masses can now be written

\[
\frac{M_a(\mu_{\text{EW}})}{g_a^2(\mu_{\text{EW}})} = \frac{M_0}{g_a^2(\mu_{\text{UV}})} \left[ 1 + \alpha m \frac{g_a^2(\mu_{\text{UV}})}{16\pi^2} (b_a - \alpha g N_m) \ln \left( \frac{M_{\text{Pl}}}{m_{3/2}} \right) \right],
\]

where the first term in the square brackets is precisely the case of gaugino mass unification at the high scale, as in the minimal supergravity paradigm. As in [21, 35] we see that departures from
Parameters

|     | \(\tilde{\alpha}\) | \(M_0\) [GeV] | \(M_0\) [GeV] | \(\alpha_m\) | \(N_m \times \alpha_g\) |
|-----|------------------|---------------|---------------|--------------|------------------|
| Mirage | 6.52         | 587           | 1976          | 1.5          | 0                |
| DMM    | 6.52          | 599           | 889           | 3.4          | -4.3             |

Physical Quantities

|     | H%     | \(M_X\) [GeV] | \(\Omega_X h^2\) | \(\sigma^p_{SI}\) [cm\(^2\)] | \(\Phi_{\tilde{\chi}^0} \geq 1\) GeV [cm\(^{-2}\)s\(^{-1}\)] |
|-----|--------|---------------|-----------------|-----------------|-----------------------------|
| Mirage | 98.1%  | 1247          | 0.13            | 0.63 \times 10^{-45} | 6.51 \times 10^{-12}        |
| DMM    | 99.5%  | 1040          | 0.10            | 2.86 \times 10^{-45} | 2.41 \times 10^{-12}        |

Table 8: Two Degenerate Models from the Mirage Family. The mirage model and the DMM model have nearly identical values for the gaugino mass-determining parameters of equation (44). Both have nearly Higgsino-like LSPs of approximately a TeV in mass. Both give rise to a thermal relic abundance of LSP neutralinos in line with the WMAP preferred region of (32) and have similar predictions for the observables in Figure 15.

universality are governed by the parameter \(\alpha_m\), but the slopes of these departure trajectories are altered relative to the case of simple mirage mediation. We can make this more explicit by defining new variables

\[
\tilde{\alpha} \equiv \left[ \frac{\alpha_m g_a^2 (\mu_{\text{EW}})}{16\pi^2} \ln \left( \frac{M_{\text{Pl}}}{m_3/2} \right) \right]^{-1} - N_m \alpha_g ; \quad \tilde{M}_0 \equiv \alpha_m \frac{M_0}{16\pi^2} \ln \left( \frac{M_{\text{Pl}}}{m_3/2} \right) = \frac{m_3/2}{16\pi^2} \quad (44)
\]

such that (43) becomes

\[
\frac{M_a(\mu_{\text{EW}})}{g_a^2(\mu_{\text{EW}})} = \tilde{M}_0 (\tilde{\alpha} + b_a) . \quad (45)
\]

From (45) it is immediately apparent that despite the added complexity of the gauge-charged messenger sector the gaugino masses continue to depend on only a single dimensionless parameter in addition to an overall mass scale, here given by \(\tilde{M}_0\). In other words, for every mirage model with \(N_m = 0\) there exists a family of deflected mirage models with the same set of soft supersymmetry breaking gaugino masses at the electroweak scale. Any differences in phenomenology will therefore be due to the differences in scalar masses and the resultant EWSB parameters.

The parametrization in (45) makes it easy to construct such look-alike “degenerate” points. Consider the two example points in Table 8. The (un-deflected) mirage model and the DMM model have nearly identical values for the key parameters \(\tilde{\alpha}\) and \(\tilde{M}_0\) of (44) despite having very different underlying physics inputs. The mirage model has modular weights \(n = 1/2\) for all fields.

\[\text{Note that this variable is equivalent to} \ M_s \text{ defined in} [22] \text{ and utilized by} [23, 26].\]
and $\tan \beta = 2.5$ while the deflected mirage model has vanishing modular weights and $\tan \beta = 27.8$. Both are in accordance with the WMAP “preferred” relic abundance in equation (32) and have similar predictions for the two observables plotted in Figure 15. In particular, both would give an observable signal in one ton-year of exposure on a xenon target but neither would be detectable above background at Fermi/GLAST with the NFW halo profile (though both would be very nearly detectable if the NFW + AC profile was used). Their predictions for other observables are also similar in magnitude.

Though many such degenerate pairs can be constructed, it is not the case that there are no signals associated with relic neutralinos that are sensitive to the presence of gauge mediation in the underlying model. We will here merely point out two such examples. To do so we will change variables once again to a dimensionless quantity that is directly proportional to the number of gauge-charged messengers in the theory. Writing out (17) in terms of the explicit mass scales in the theory as

$$M_a(\mu) = M_0 \left\{ 1 + \beta_a(\mu) \left[ t - \frac{N_m}{b_{a}} \ln \left( \frac{\mu_{\text{GUT}}}{\mu_{\text{mess}}} \right) \right] - \frac{\beta_a(\mu) N_m}{2 b_{a}} \Lambda_{\text{mess}} + \frac{\beta_a(\mu)}{2} m_{3/2} \right\},$$

we may apply the definitions (12) and (13) to write (46) as

$$M_a(\mu) = M_0 \left[ 1 + \beta_a(\mu)t \right] + m_{3/2} \frac{\beta_a(\mu)}{2} \left[ 1 - \frac{\alpha'}{b_{a}} \right],$$

where we have introduced a new dimensionless parameter $\alpha'$ defined as

$$\alpha' = N_m \left[ \alpha_g + \frac{2}{\alpha_m} \frac{\ln \left( \frac{\mu_{\text{GUT}}}{\mu_{\text{mess}}} \right)}{\ln \left( \frac{M_{\text{pl}}}{m_{3/2}} \right)} \right].$$

Although specifying the value of $\alpha'$ does not uniquely determine the underlying model, all mirage models will have $\alpha' = 0$, which makes this a convenient parametrization for our purpose here. When one plots inclusive counting signatures as a function of $\alpha'$ – such as the number of recoils per kg-year on xenon targets or the total integrated flux of photons coming from the galactic center – the results are much as in Figure 15. But certain exclusive measurements reveal some correlation. For example, one can take advantage of the Fermi/GLAST experiment to measure the spectral profile of detected gamma rays. The shape of this spectrum carries information about the mass and wave-function content of the annihilating neutralino LSP. We can attempt to take this information into account by using DarkSUSY to compute the differential photon flux in 1 GeV increments over the energy range $1 \text{ GeV} \leq E_\gamma \leq 200 \text{ GeV}$. From this information we create an interpolating function...
which is then integrated over six energy bins: 1 - 10 GeV, 10 - 30 GeV, 30 - 60 GeV, 60 - 100 GeV, 100 - 150 GeV, and 150 - 200 GeV.

In Figure 16 we take the ratio of the total flux over the three highest energy bins to the total flux over the three lowest energy bins and plot it as a function of the parameter $\alpha'$. Over the range $-10 \leq \alpha' \leq 10$ (which is the regime of moderate $\alpha_g$ values) the hardness of the photon spectrum is reasonably well correlated with the value of $\alpha' \simeq N_m \times \alpha_g$. Note that over this range of parameters we generally expect the gamma rays from neutralino annihilation to have a harder spectrum than those coming from background astrophysics sources, as indicated by the heavy horizontal line. The collection of models in Figure 16 was generated using a fixed mass scale of $M_0 = 500$ GeV. The degree of correlation increases if we choose higher values for this parameter.

More generally, the presence of the gauge-charged messengers allows the theory space to explore neutralino LSPs with differing wave-function composition from that of the mirage case with $N_m = 0$. 

Figure 16: Hardness of the Gamma Ray Spectrum from the Galactic Center as a Function of $\alpha'$. Flux of photons in the energy range $60 \text{ GeV} \leq E_\gamma \leq 200 \text{ GeV}$ divided by the flux in the range $1 \text{ GeV} \leq E_\gamma < 60 \text{ GeV}$ is plotted for the signal photons versus the parameter $\alpha'$ of equation (48). The same ratio for the background is given by the heavy horizontal line. Note the logarithmic scale of the vertical axis.
Figure 17: **Combined Monochromatic Line Flux as a Function of Neutralino Mass.** The solid line represents the reach at a generic ACT for 500 m$^2$-years of exposure with an NFW profile. The dashed line is the reach for the same exposure assuming an NFW + AC profile.

This is already apparent from Figure [13]. Such effects are known to have a large impact in the rate for loop-induced annihilation into monochromatic photons. In Figure 17 we give the combined flux at both the $\gamma\gamma$ and $\gamma Z$ energies for our random sample of DMM models with $M_0 = 500$ GeV (dark points) and our random sample of mirage model (light points) at the same $M_0$ value. When the overall mass scale $M_0$ is held fixed the points tend to cluster as a function of the neutralino wave-function composition. In the mirage mediation model without gauge messengers the majority of models have a bino-like LSP with a smaller grouping of Higgsino-like LSPs. In fact, the maximal value of the wino-component of the LSP was typically 20%, and these occurred mostly in cases with vanishing modular weights for the matter multiplets and $n = 1$ for the Higgs multiplets. By contrast the DMM models are able to achieve a wino-like LSP across all modular weight choices. Assuming an NFW halo profile only those cases where the LSP is nearly 100% wino-like have a chance to be observed in a typical ACT experiment. Assuming an adiabatically-compressed profile allows most of the wino-like and Higgsino-like cases to be detected in 500 m$^2$-years of exposure.
5 Conclusion

Deflected mirage mediation, like the simpler mirage pattern which preceded it, are motivated by realistic string compactifications which stabilize all geometrical moduli while simultaneously addressing the vacuum energy problem of supersymmetry breaking. As such they are important cases for phenomenological investigation in their own right and have justly received much attention. More generally, if we are indeed on the eve of a new supersymmetric data era then it is of the utmost importance to begin the study of how this wealth of data can be used to understand the underlying principles of supersymmetry breaking. The deflected mirage paradigm provides a framework that is remarkably rich, allowing a smooth interpolation between supergravity (modulus) mediation, anomaly mediation, and gauge mediation. It is an important empirical question to ask how well we can isolate these various contributions to the transmission of supersymmetry breaking to the observable sector if all three are present at roughly the same mass scale.

Ultimately we would like to answer this question in as model-independent a manner as possible, a point emphasized by Choi and Nilles [21]. As those authors point out, gaugino masses are a much cleaner window on the nature of supersymmetry breaking and transmission than the rest of the soft supersymmetry-breaking Lagrangian. This remains true in the deflected mirage paradigm as well. For this reason we have chosen to begin answering the question of the previous paragraph with dark matter observations which depend strongly on the gaugino sector and relatively less strongly on the scalar sector. Our results indicate that loosening the restrictions imposed by the original mirage pattern allows for an expanded parameter space in which the LSP can be predominantly wino-like, or mixed Higgsino/gaugino. This regime roughly corresponds to circumstances in which $0 \lesssim N_m \times \alpha_g \lesssim 10$, or where $N_m \Lambda_{mess} = cm^{3/2}$ with $c = O(1)$. In such cases a number of direct and indirect detection experiments running now or scheduled to begin data taking in the near future have excellent prospects for detecting the relic neutralino. This is also the area in which a wino-like LSP with roughly the right properties to explain the PAMELA positron data can be found. For this statement to be borne out, however, it is still necessary to make modifications to the standard diffusion parameters and to imagine non-thermal production mechanisms to explain the present number density. It would be worthwhile to further investigate the utility of DMM-like models at fitting the PAMELA positron data. If the data does indeed seem to point toward a wino-like LSP, it would be interesting to see how a mixed-mediation wino arising from the DMM scenario would fare when a more sophisticated analysis is performed.

Though the predictions from the deflected mirage mediation model depend strongly on the pres-
ence and magnitude of the gauge-mediation sector, it is not necessarily the case that observations in the dark matter arena are in and of themselves sufficiently powerful to determine this fact. That dark matter observations should suffer from such an “inverse problem” is perhaps not surprising. In fact, the gaugino sector of the DMM scenario is still a two-parameter family of models and is thus not fully general. Any trio of gaugino masses in the DMM theory space can be mapped to an equivalent trio on an un-deflected model. Thus dark matter signals will distinguish between the two paradigms only via the scalar sector, which serves to determine the LSP properties only via the electroweak symmetry-breaking conditions. Using this fact to detect the presence of the gauge messenger sector will prove difficult, but some handle can be obtained via gamma ray signals. To make optimal use of this data we would prefer to have some knowledge of the relevant mass scales for the gaugino sector. For this it is crucial that the collider signatures for the deflected mirage mediation model be computed as a function of the contribution of the three mediation mechanisms.

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