A WINO-LIKE LSP WORLD:
Theoretical and Phenomenological Motivations

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Central to the identification of the lightest superpartner (LSP), and the accompanying spectra of heavier supersymmetric particles, is to understand the particle content of the LSP mass eigenstate. There are strong arguments, both theoretical and phenomenological, that point to a LSP with a significant wino component that can contribute substantially to the relic density of dark matter and even accommodate the entire relic abundance. Data from the Large Hadron Collider (LHC) is expected to test signals of such models. Early signatures from the cascades of colored superpartners may point to the presence and form of softly broken supersymmetry, and of electroweak symmetry breaking.

1.1. Annihilating Dark Matter in the Halo

From the standpoint of annihilating dark matter, among the mechanisms that can simultaneously explain the PAMELA\(^1\) and WMAP\(^2\) relic density data, some are based on a wino-like LSP. These include a non-thermally produced wino LSP\(^3\) with the relic abundance explained via moduli decay\(^4\) which arises in a concrete framework of fluxless string compactifications,\(^6,7\) or a thermal wino-like LSP with a weakly interacting co-annihilating hidden sector with mass terms for the Majoranas arising from Stueckelberg fields and/or Kinetic Mixing.\(^8\) Here we will focus on the collider and cosmological implications of the wino dominated dark matter. Other dark matter annihilation mechanisms capable of explaining both data sets are a Sommerfeld enhancement in the halo,\(^9\) a Breit-Wigner enhancement of dark matter annihilations in the halo,\(^10\) and Leptonic Asymmetric Dark Matter,\(^11\) and we refer the reader to the original papers for details on these proposals.
An early approach that gave a string motivation for a wino LSP was the proposal that the mechanism of anomaly mediation might be the dominant mechanism of supersymmetry breaking\textsuperscript{14,15} (for early foundational work see\textsuperscript{16}). If a nearly pure wino LSP is to account for the entire relic density it has become more clear that the universe likely has a non-thermal cosmological history. This has received increasing attention as a generic feature of a comprehensive underlying theory; we comment briefly on this and a more detailed discussion is given in\textsuperscript{12}. If however the LSP has a relatively large wino component but has a suitable admixture of other eigen-components, a thermal mechanism for explaining its relic abundance is quite possible. We review this idea as well. Our goal in this chapter is simply to emphasize that in the past decade a wino-like LSP has been recognized to be both theoretically and phenomenologically well motivated. We proceed by describing several examples for readers who wish to see more details.

1.2. Non-thermal winos from moduli stabilized on a $G_2$ manifold

1.2.1. Soft breaking from the $G_2$

A recent string model of interest that predicts a LSP that is dominantly wino is that of the $G_2$-MSSM. The $G_2$-MSSM was formally introduced in Ref.\textsuperscript{6} and was motivated by the work of Acharya and Witten,\textsuperscript{5} where it was demonstrated that $M$-theory compactifications on a manifold $X$ of $G_2$ holonomy can give rise to chiral fermions in four dimensions only if $X$ is not smooth; specifically the fermions need to be localized at conical singularities. Thus in Ref.\textsuperscript{6} it was shown that moduli stabilization is manifest in a large class of $M$-theory compactifications in the zero-flux sector, with minimally two non-abelian asymptotically free gauge groups which gives rise to a moduli potential. The simultaneous breaking of supersymmetry becomes possible when at least one of the hidden sectors contains charged matter, and the cosmological constant can be tuned towards zero.

The underlying framework is described by $\mathcal{N} = 1$ supergravity where a generalized sector of soft susy breaking is derived from both a tree level supergravity contribution\textsuperscript{17} and an anomalous supergravity contribution\textsuperscript{14,16}, as in\textsuperscript{15}. The soft parameters can be parametrized at the unification scale as $m_0 = s \cdot m_{3/2}$, $m_a = f_a \cdot m_{3/2}$, $A_3 = a_3 \cdot m_{3/2}$ where $m_{3/2} \sim \mathcal{O}(10^{-100})$ TeV is the gravitino mass, $m_0$ is a universal scalar mass, $m_a$ are the gaugino
masses, \( A_3 \) are the tri-linear couplings of the third generation, and \( \tan \beta \) is found to lie in the range \( \tan \beta \sim 1.5 - 2.0 \). Here the parameters \((s, f_\alpha, a_3)\) are functions of the microscopic theory which are determined by specification of the Kähler potential \( K \), superpotential \( W \) and gauge kinetic function \( f \) (see below). The soft parameters are well approximated by (for the complete analytical expressions see\(^6\)) \( s \sim 1 \), \( f_\alpha = f_\alpha' \alpha_G - \epsilon \eta \), where \( f_\alpha' = (0.35, 0.58, 0.64) \), and \( \eta = 1 - \alpha_G \delta \) parametrizes gauge coupling corrections in the tree level sector of the gaugino masses, and \( \epsilon \sim 0.024 \) depends on the parameters of the hidden sector potential which are responsible for tuning the cosmological constant to zero. The terms entering for the tri-linears of the third generation are well approximated by \( a_3 = 1.67 - 3 \epsilon \ln(4\pi^{1/3}/V_7) - 2 \epsilon \ln(|Y_{t,b,\tau}|) \), where \( Y \) are normalized Yukawas and \( V_7 \) is the normalized volume of the \( G_2 \) manifold which also enters in the determination of the gravitino mass. Here the largeness of the gravitino mass decouples the scalars while the gaugino masses are suppressed relative to gravitino mass, where the suppression enters via the volume of hidden sector three cycles. The physical values of the soft parameters are sensitive to the precise value of the unified gauge coupling and threshold corrections. The largeness of the gravitino mass drives the \( \mu \) term to be order \( m_3/2 \) for electroweak symmetry breaking, which in turn induces a relatively large self energy correction\(^{18}\) to the electroweak gaugino masses.\(^7\)

The above arises from coupling the moduli and hidden sector matter to the visible SUGRA sector with the following Kähler potential \( K \), superpotential \( W \) and gauge kinetic function \( f \) at the compactification scale \( \sim M_{\text{unif}} \).\(^{6,7}\)

\[
\frac{K}{M_{\text{pl}}^2} = -3 \ln(4\pi^{1/3}/V_7) + \phi \bar{\phi}, \quad V_7 = \prod_{i=1}^{N}s_i^2, \quad a_i \in Q^+
\]

\[
W = \frac{M_{\text{pl}}^3}{2} \left(C_1 P \phi^{-2/(P)} \bar{\epsilon}^{ib_1 f_1} + C_2 Q \bar{\epsilon}^{ib_2 f_2}\right); \quad b_1 = \frac{2\pi}{P}, \quad b_2 = \frac{2\pi}{Q}
\]

\[
f_1 = f_2 \equiv f_{\text{hid}} = \sum_{i=1}^{N} N_i z_i; \quad z_i = t_i + is_i.
\]

(1.1)

Here, \( V_7 \) is the volume of the \( G_2 \) manifold \( X \) (in units of the eleven-dimensional Planck length \( l_{11} \)), the composite scalar \( \phi \) is the effective meson field; there is one for each pair of massless quarks which forms the composite. \( f_{1,2} \) are the gauge kinetic functions of two hidden sectors at tree-level, \( s_i \) are the \( N \) geometric moduli of the \( G_2 \) manifold and \( t_i \) are axionic scalars. The parameters \( P \) and \( Q \) are proportional to the beta
function coefficients of the gauge groups. The difference \( Q - P \) is fixed by the solution corresponding to a metastable minimum with spontaneously broken supersymmetry and one finds \( Q - P \sim 3 \). The effective parameter 
\[ P_{\text{eff}} = \text{ln}(C_1/C_2) \]
 takes a value of 84 for \( Q - P = 3 \) as governed by the tuning of the cosmological to zero. It is this relative smallness of the factor 
\[ 1/P_{\text{eff}} \]
 which enters for example in the tree-level gaugino sector, that allows the anomaly terms and tree terms to contribute to the soft masses on the same footing.

1.2.2. **Moduli Masses**

The system of moduli fields which participate in reheating the universe at various epochs are the \( N \) geometric moduli \( s_i \), and the hidden sector meson field \( \phi \). Since these moduli and meson will mix in general, the physical moduli correspond to mass eigenstates. There is one heavy eigenstate with mass 
\[ m_X = (7K_1/3 + K_2)^{1/2}m_{3/2} \]
 along with \((N-1)\) degenerate light eigenstates with mass 
\[ m_{X_j} = (K_2)^{1/2}m_{3/2} \]
 and an eigenstate with mass 
\[ m_\phi = (K_4 - K_2^2)^{1/2}m_{3/2} \]. The \( K \) factors are uniquely determined in terms of the VEV of the meson field and the parameters discussed above. A remarkable result is that 
\[ m_{X_j} \approx 1.96 \times 10^{-7} \text{GeV} \]
(1.2)
That is, the light moduli are phase space restricted and the decay into gravitinos are suppressed which essentially eliminates the moduli induced gravitino problem.\(^6,7\)

1.2.3. **The right halo cross section and just about the right abundance from non-thermal winos**

For a pure wino in the large \( \mu \) limit such as that which arises in the \( G_2 \) models, one has 
\[ \tilde{N}_1^0 = \tilde{W}, \quad \tilde{N}_2^0 = \tilde{B}, \quad \tilde{N}_3^{0,4} = \frac{1}{\sqrt{2}}(\tilde{H}_1 \mp \tilde{H}_2), \]
where 
\[ \tilde{N}_n^0 = N_{nm}\tilde{Z}_m, \quad \tilde{Z}^T = (\tilde{W}, \tilde{B}, \tilde{H}_1, \tilde{H}_2) \] and the annihilation cross section of the winos into \( W \) Bosons proceeds through an s-wave via chargino exchange
\[ \langle \sigma v \rangle = \frac{g_4^2}{2\pi m_W^2} \frac{(1 - m_W^2/m_W^2)^{3/2}}{(2 - m_W^2/m_Z^2)^2} \] (1.3)
with \( \langle \sigma v \rangle \sim (1.6 - 2.2) \times 10^{-7}\text{GeV}^{-2} \) for \( m_W \in (170, 200) \text{ GeV} \). The lack of velocity dependence and any non-pertubative effects in this limit fixes the halo cross section with the wino mass determined. As far as the
relic abundance, quite generally, in the absence of co-annihilations for s-wave annihilations the relic abundance follows from solving the Boltzmann equations. This is equivalent to

$$\Omega h^2 = \frac{h^2 m_W^3 n_W/s}{\rho_c/s_0} = \frac{h^2}{4(\rho_c/s_0)} \frac{90}{\pi^2 g_*} \frac{1}{M_{pl}} \frac{1}{\langle \sigma v \rangle} \frac{m_W}{T}$$

where $n_W = \frac{H}{\langle \sigma v \rangle}$, $H^2 = T^4 g_* T^2/(90 M_{pl}^2)$, $s = (2\pi^2/45) g_* T^3$ and $\rho_c/s_0 = 3.6 \times 10^{-9}$ GeV $h^2$, with $g_* = g_*(T)$. Note that the above holds for both the thermal case and the non-thermal case; the difference for the non-thermal case arises in that the neutralino density can become repopulated well after freeze out (where $m/T \sim 20$ at $T = T_{\text{freeze}}$). Instead, for the non-thermal case, the reheat temperature $T_R$ occurs at temperatures on the scale (5-100) MeV, depending on the particular model.

Thus the presence of moduli (denoted here now more generally by $M$) are ubiquitous in string theories whose low energy limit are supergravity models (as we have just discussed in the previous section). If such a scalar decays after freeze out it can decay into susy particles repopulating the neutralino abundance. Here we discuss the phenomena in some generality.

Assuming the energy density during modulus decay is transferred completely to radiation one has $H = \Gamma_M$, and thus $T_R = (90 \pi^2/g_*)^{1/4} \sqrt{\Gamma_M M_{pl}}$, where the modulus decay is parametrized as $\Gamma_M = c_M M^3 / \Lambda^2$. The scale $\Lambda$, the modulus mass $M$ and its coupling are determined by the underlying model. In order to avoid the gravitino problem (for recent work see\cite{19}), the modulus is taken to satisfy $M < 2 m_{3/2}$ (as mentioned previously this is remarkably manifest in the $G_2$ models\cite{6}), where $m_{3/2}$ is the gravitino mass. The relic density calculation of the abundance of winos is then entirely straightforward with the temperature $T \to T_R$ and one obtains

$$\Omega_{\tilde{W}} h^2 \simeq 0.32 \frac{1}{\sqrt{c_M}} \left( \frac{3 \times 10^{-7} \text{GeV}^{-2}}{\langle \sigma v \rangle} \right) \left( \frac{m_{\tilde{W}}}{200 \text{GeV}} \right) \left( \frac{m_{3/2}}{100 \text{TeV}} \right)^{-3/2}$$

where $\langle \sigma v \rangle$ is to be velocity averaged with a Boltzmann distribution and we have used $g_* = 10.75$ (the appropriate degrees of freedom for MeV scale temperatures), $M \sim 2 m_{3/2}$ and $\Lambda = M_{pl}/\alpha$, where $\alpha$ parametrizes deviations from the reduced Planck scale. For example, moduli couplings at a string scale would correspond to $\alpha \sim \sqrt{\Lambda}$, where $\Lambda$ is the (dimensionless) volume of compactification and thus for a pure wino with mass of $\sim 180$ GeV, the WMAP relic density $(\Omega_{\tilde{W}} h^2 \sim 0.1)$ is then achieved for $c_M \sim 1$, \...
\[ \Lambda = M_{\text{String}} \sim 4 \times 10^{17} \text{GeV}, \] (which is similar to the heterotic string scale)

for gravitino mass of \( \sim 40 \) TeV.

For the case of the \( G_2 \)-MSSM \( \alpha = 1 \) and \( c_M \in (0.5, 4) \) over the part of the model parameter space investigated. Thus a gravitino mass on order of 100 TeV is needed for a pure wino with mass of 180 GeV to give the observed relic density within the WMAP error band. A more detailed understanding of the underlying string theory may allow for a lighter modulus mass.

1.3. Wino-Like Dark Matter in the Stueckelberg Extensions and Kinetic Mixings

1.3.1. The right relic abundance and just about the right halo cross section from extra \( U(1)_X \) Factors

The presence of matter in the hidden sector can have observable effects on the visible (MSSM) sector. One such effect is that extra Majorana matter very weakly coupled to the MSSM neutralinos can coannihilate with the LSP which has the effect of enhancing the relic density for the LSP by as much as an order of magnitude or more.\(^8,40\) This enhancement of the relic abundance can occur through the presence of \( n \ U(1)_X \) gauge symmetries in the hidden sector and correspondingly \( n \) new scalars (axions) each being absorbed through a Stueckelberg mechanism\(^{40,43}\) generating masses for the \( n \ U(1)_X \) gauge Bosons.\(^8\)

Thus, in top-down approaches to building realistic models based on D-Branes the SM gauge group can be produced but one also encounters residual Abelian group factors where the extra \( U(1) \)s usually correspond to massive vector fields. In particular frameworks they lead to terms in the action of the form \( B \wedge F \), i.e. terms of the form \( \frac{1}{2} \epsilon_{\mu\nu\alpha\beta} B_{\mu\nu} F_{\alpha\beta} \) which are needed for anomaly cancellation via a [four dimensional] Green-Schwarz (GS) mechanism.\(^{35}\) In other frameworks the \( B \wedge F \) couplings can arise for the non-anomalous cases as well,\(^{34,38-40}\) for the anomalous case see i.e.\(^{41,42}\). These types of couplings can give rise to Stueckelberg mass terms. For example, under the duality transformation \( \partial_\mu \sigma \sim \epsilon_{\mu\nu\rho\sigma} \partial^\nu B^{\rho\sigma} \), as illustrated in,\(^34\) vector fields gain mass via the GS mechanism leading to Stueckelberg mass terms in the Lagrangian of the form

\[
L_{\text{St}} = -\frac{1}{2} \sum_I (\partial^\mu \sigma^I - k^I a C_a^\mu)^2, \tag{1.6}
\]

with the index \( I \) running over all the 4D fields and thus runs over each of pseudo-scalars \( \sigma^I \) (Ramond axions) and over the killing vector coefficients.
Here \( a \) indexes a Brane stack with a supporting gauge group (for example: \( U(N_a) = SU(N_a) \times U(1)_a \)) with each Abelian vector field denoted as \( C^a_\mu \). It is then clear that the quadratic term for \( C^a_\mu \) in Eq.(1.6) gives rise to mass terms for the Abelian vector fields. The orthogonality of the killing vectors determines which Bosonic states will be rendered massive or if they remain massless. Upon addition of gauge fixing terms, the cross terms in Eq.(1.6) cancel and thus the psuedoscalars and vector fields decouple. The mechanism outlined here is distinct from a Higgs mechanism for mass generation as there is no extra dynamical scalar field.

In the supersymmetric case, a Stueckelberg extension of the MSSM\(^{40}\) arises from an extended Lagrangian with \( N_V \) abelian gauge fields and \( N_S \) axions

\[
\mathcal{L}_{St} = \int d^2 \theta d^2 \bar{\theta} \sum_{j=1}^{N_S} \left( S_j + \bar{S}_j + \sum_{i=1}^{N_V} M_{ij} C_i \right)^2,
\]

where \((S_j, C_i)\) are (chiral,vector) supermultiplets.

The case of interest to us here is when \( N_S = N_V - 1 \) so generically all \( N_V - 1 \) number of axions are absorbed by \( N_V - 1 \) number of vector bosons making them massive and leaving one vector boson massless. To generate communication between the hidden and the visible sectors we thus identify one of the fields in the \( N_V \) set to be the hypercharge, e.g., \( B_Y = C_1 \). When one includes the electroweak sector of MSSM, there would then be an automatic coupling between the hidden and the visible sectors through \( B_Y \) which enters both in the MSSM sector and in the Stueckelberg sector. Consequently a spontaneous breaking in the MSSM sector with inclusion of the above interactions will lead to a massless photon, a weakly coupled \( Z' \) boson and a system of \( N_V - 1 \) extra \( Z' \) bosons. Additionally there will be \( N_V - 1 \) number of CP even Higgs fields.

In the fermionic sector we gain a Stueckelberg chiral fermion for each \( S_j \) and a Stueckelberg gaugino for each \( C_i \). These mix with the neutral fermions of the MSSM sector producing a neutralino mass matrix which is \( 4 + N_S + N_V \) dimensional. Thus for \( n \) number of \( U(1)_\chi \), there are \( 2n + 4 \) Majorana states:

\[
(N^0_1, n^0_1, n^0_2, \ldots n^0_{2n}), N_2^0, N_3^0, N_4^0)
\]

where \( N_i^0 (i = 1, 2, 3, 4) \) are essentially the four neutralino states of the MSSM and \( n^0_\alpha, (\alpha = 1, \ldots, 2n) \) are the additional states.\(^5\) With the Majorana fields of the hidden sector interacting extra weakly, their interaction cross sections are suppressed, and one finds that there is an enhancement
of the relic density by a factor $B_{\text{Co}}$ through coannihilation effects entering through the number of degrees of freedom supplied by the hidden sector Majoranas. This enhancement to the relic density (to be distinguished physically from a boost in the halo!) is given by

$$B_{\text{Co}} \simeq \frac{\sum_{a,b} \int_{x_1}^{\infty} (\sigma_{ab} v) \gamma_a \gamma_b \frac{dx}{x^2}}{\sum_{A,B} \int_{x_1}^{\infty} (\sigma_{AB} v) \Gamma_A \Gamma_B \frac{dx}{x^2}}.$$ 

$$\gamma_a = \frac{g_a (1 + \Delta_a)^{3/2} e^{-\Delta_a x}}{\sum_b g_b (1 + \Delta_b)^{3/2} e^{-\Delta_b x}} \text{MSSM}$$

$$\Gamma_A = \frac{g_A (1 + \Delta_A)^{3/2} e^{-\Delta_A x}}{\sum_A g_A (1 + \Delta_A)^{3/2} e^{-\Delta_A x}} \text{MSSM} \otimes \text{Hid.}$$

Here $a$ runs over the channels which coannihilate in the MSSM sector, while $A$ runs over channels both in the MSSM sector and in the hidden sector (i.e., $A=1, \ldots, n_v+n_h$). For the case when these fields that coannihilate are completely degenerate with the LSP, and the hidden sector interactions are suppressed, one has $B_{\text{Co}} = (1 + d_h/d_v)^2$, where $d_s = \sum_s g_s$, for $s = (v, h)$, i.e.

$$\left(\Omega h^2\right)_{N^0} \simeq (1 + \frac{d_h}{d_v})^2 \left(\Omega h^2\right)_{\text{MSSM}},$$

and one has for the case of $n$ hidden sector $U(1)$s the result

$$B_{\text{Co}} = (1 + 2n)^2 \quad \text{(no MSSM particle coannihilations).}$$

As studied in,\textsuperscript{8,43} such an extra weakly interacting set of hidden sector neutralinos can be sourced not only by Stueckelberg mass mixings but also through kinetic mixings. This has recently been observed in other stringy environments where the extra weak Majoranas, such as Stinos\textsuperscript{43} have been dubbed string photini.\textsuperscript{44} In fact such a situation can arise from both Stueckelberg mass and Kinetic mixings operating at comparable scales.\textsuperscript{43,45}

Quite generally, in order for the above models to produce a significant flux in the halo they must lead to the LSP having a significant wino component. Such eigen content arises in non-universal SUGRA models and string models with non-universalities in the gaugino sector (for early work see\textsuperscript{30}). A pure wino can be obtained by driving down $M_2(GUT)$ and halo cross sections of the size $2 \times 10^{-7}\text{GeV}^{-2}$ can manifest.\textsuperscript{8} However in this case, the degenerate set of $U(1)$ factors must be large, since MSSM coannihilations reduce the maximal $B_{\text{Co}}$ to $B_{\text{Co}} = (1 + 2n/3)^2$. Yet when MSSM co-annihilations are reduced, the wino component can also remain large $N_{1,2} \sim 0.70$, and there can be non-negligible mixtures of bino and
Higgsino components. In this wino mixed case, the annihilation cross section in the halo can be about a factor of 5 lower than the pure wino case, while the $B_{C_0}$ enhancement of the relic density, is large enough (maximally about 45 in practice with a $U(1)_X^3$ gauge symmetry) to bring the mixed wino neutralino abundance up to the lower limit on the WMAP constraint. Fits to the PAMELA data require a small boost in the halo on the order of (3-5). Thus such a mixed wino model is less susceptible to potential overproduction of protons and photons but can still produce a good fit to the PAMELA data and the WMAP data. Nevertheless, a large wino component is needed.

### 1.4. Directly Detecting Wino-Like Dark Matter

Experiments are actively attempting to detect dark matter via their spin dependant and spin independent scattering with nuclei. The LSPs have a velocity distribution near the earth and in the local galactic halo, and they are travelling with non relativistic speed order $0.001c$. As such their momentum transfer is rather small (order 100 MeV for LSP masses of order 100 GeV). Therefore, the relevant interactions for the direct detection of dark matter (LSP collisions with nuclei) may be calculated in the limit of zero momentum transfer. For the case of the MSSM, over most of the viable parameter space, the relevant piece of the interaction Lagrangian is

$$
\mathcal{L} = \bar{\chi} \gamma^\mu \chi \gamma^5 \chi \gamma^5 q_i \rightarrow \left(\alpha_1 + \alpha_2 \gamma^5 \right)q_i + \alpha_3 \bar{\chi} q_i q_i + \alpha_4 \bar{\chi} \gamma^5 \chi \gamma^5 q_i q_i + \alpha_5 \bar{\chi} \gamma^5 \chi q_i q_i + \alpha_6 \bar{\chi} \chi \gamma^5 \chi q_i q_i .
$$

Indeed, the spin independent (SI) cross section is currently being probed in experimental searches. For the cross section of neutralinos scattering elastically off target nuclei, in terms of the reduced mass of the neutralino and the target system ($\mu_{\chi T}$), one has

$$
\sigma_{\chi(T)} = \frac{4 \mu_{\chi T}^2}{\pi} \left( Z f_p + (A - Z) f_n \right)^2 ,
$$

where $(Z, A)$ are the atomic (number,mass) of the nucleus, and the interactions with the quarks in the target nuclei are dominated by $t$-channel CP-even Higgs exchange, and $s$-channel squark exchange and are housed in

$$
f_{p/n} = \sum_{q=u,d,s} f_{Tq}^{(p/n)} a_q \frac{m_{p/n}}{m_q} + \frac{2}{27} f_{Tq}^{(p/n)} \sum_{q=c,b,t} a_q \frac{m_{p/n}}{m_q} .
$$
Here \( f_{TG}^{(p/n)} \) is given by \( 1 - f_{T_u}^{(p/n)} - f_{T_d}^{(p/n)} - f_{T_s}^{(p/n)} \) and arises via gluon exchange with the nucleon and the \( f_{T_u}^{(p/n)} \) are determined from light quark masses obtained from baryon masses via matrix elements and from the value of the pion-nucleon sigma-term. Numerical values and further details are given in, for example, in. Under constraints from collider data on the sparticle masses and mixings, the dominant couplings that enter in the spin independent cross section are given in, for example, in.

The first term arises from squark exchange and is typically much suppressed, while the dominant effects enter through the exchange of Higgs Bosons. The parameters \( \delta_{1,2} \) depend on eigen components of the LSP wave function and \( B, C, D \) depend on VEVs of the Higgs fields and the Higgs mixing parameter \( \alpha \) and are given by

for u quarks: \( \delta_1 = n_{13} \) \( \delta_2 = n_{14} \) \( B = \sin \beta \) \( C = \sin \alpha \) \( D = \cos \alpha \)

for d quarks: \( \delta_1 = n_{14} \) \( \delta_2 = -n_{13} \) \( B = \cos \beta \) \( C = \cos \alpha \) \( D = -\sin \alpha \).

In order for wino-like dark matter to give rise to a detectable spin independent cross section, at current sensitivities, a mixture of Higgsino content is needed, where the spin independent nucleon cross section of \( \sim 5 \times 10^{-44} \text{ cm}^2 \) at dark matter mass of \( \sim 60 \text{ GeV} \) is currently the minimum of the confidence limits. The LSP can have a large wino component and produce a large spin independent cross section as well as a large flux in the halo. Such can be achieved, for example in a model with split gaugino masses at the GUT scale such as \( M_a = m_{1/2}(1 + \Delta_a) \), \( a = 1, 2, 3 \) with \( (m_0, m_{1/2}, A_0, \tan \beta, (\Delta_1, \Delta_2, \Delta_3)) = ((1000, 800, 0)\text{GeV}, 10, (-47.7), (-4)) \), with \( \text{sign}(\mu) > 0 \), where \( (m_0, A_0) \) are the universal scalar mass and trilinear coupling, respectively, and \( \tan \beta \) is the ratio of the Higgs VEVs and \( \mu \) enters as the bilinear term in the superpotential (see i.e.29). After radiative

\(^{a}\)a natural class of models where this happens are grand unified models such as \( SU(5) \), \( SO(10) \), and \( E_6 \) where the GUT symmetry is broken by a non-singlet \( F \) term (for recent work in this direction see31-33).
electroweak symmetry breaking the produced spectrum and mixing leads to a wino-like eigenstate with both a strong halo cross section and scattering cross sections:

\[(N_\tilde{B}, N_\tilde{W}, N_\tilde{H}_1, N_\tilde{H}_2) = (0.234, -0.957, 0.161, -0.064),\]

\[\sigma_{SI} = 1 \times 10^{-8} \text{ pb}, \quad \sigma_{SD} = 2 \times 10^{-5} \text{ pb},\]

\[\langle \sigma v \rangle_{W^+W^-} = 2 \times 10^{-24} \text{ cm}^3/\text{s}, \quad m_{\tilde{W}} = 185 \text{ GeV}.\]

Thus, this class of model produces positrons in the halo which describe the PAMELA data (see next section), and produces a spin independent scattering cross section within reach of the CDMS and Xenon experiments.

1.5. Positron Flux from Wino-Like Dark Matter

A wino-like LSP can annihilate in the galaxy producing \(W^\pm\) bosons. The resulting positrons diffuse through the galaxy to satellite detectors. Several analyses have shown that the positron flux ratio from the annihilations of wino-like dark matter provides a relatively good description of the PAMELA data.\(^3\)\(^,\)\(^8\)\(^,\)\(^22\)

The positron flux can be described semi-analytically. The flux enters as a solution to the diffusion loss equation, which is solved in a region with a cylindrical boundary, and is well approximated under steady state conditions by

\[\Phi(E) = \frac{B_{\nu\nu}}{8\pi b(E)} \frac{\rho_\odot^2}{m_{\tilde{\chi}_0}^2} F(E), \quad [\text{GeV} \cdot \text{cm}^2 \cdot \text{s} \cdot \text{sr}]^{-1}\]

\[F(E) = \int_{E}^{M_{\tilde{\chi}_0}} dE' \sum_k \langle \sigma v \rangle^k_{\text{halo}} \frac{dN^k}{dE'} \cdot \mathcal{I}(E, E').\]

The particle physics depends on \(\langle \sigma v \rangle_{\text{halo}}\), the velocity averaged cross section in the halo of the galaxy, and \(dN/dE\), the positron fragmentation functions. The astrophysics depends on \(b(E) = E^2/(\text{GeV} \cdot \tau_E)\) with \(\tau_E = (1 - 3) \times 10^{16} \text{s}\) which parametrizes the energy loss in the flux from the presence of magnetic fields and from scattering off galactic photons. \(\mathcal{I}(E, E')\) is the a-dimensional halo function and the minimal parameters needed to describe it are the diffusion/propagation parameters \(\delta, K_0\), and \(L\), with diffusion coefficient \(K(E) = K_0(E/\text{GeV})^4\) and \(L\) being the half height of the cylinder, (for some fits see with various halo profiles see i.e.\(^21\)). The local halo density, \(\rho_\odot\), lies
in the range $\sim (0.3 - 0.6) \text{GeV/cm}^3$ (with recent result pointing towards a value closer to $\sim 0.4 \text{ GeV/cm}^3$) in the vicinity of galactic plane and at a distance of $r_\odot$ = 8.5 kpc. $B_\xi$ is a so-called boost factor which parametrizes the possible local inhomogeneities of the dark matter distribution. Large boost factors have been used in the literature, even as large as 10,000, to explain the PAMELA data, however analyses of simulations increasingly suggest boost factors near unity. A pure wino needs no multiplicative boost factor, while a wino with some mixture requires small halo boosts.

Thus, for wino dominated dark matter, $\frac{dN_k}{dE} \rightarrow \frac{dN_{WW}}{dE}$ with $\langle \sigma v \rangle_{WW} \rightarrow W^+ W^-$ the overwhelmingly dominant source of positrons. The relative strength of the halo cross section is governed then by the dark matter mass, and the strength wino component (see Eq.(1.3) for the dominant contribution of a pure wino) however the shape of the flux distribution is controlled by the fragmentation functions and the halo/profile model. The fragmentation functions for the $W$ boson fall off at the mass of the dark matter, and therefore one expects to see a significant dip in the positron fraction at threshold. The fall off may however be compensated by an astrophysical flux from shock waves/pulsars or other astrophysical remnants.

Of course, not only are positrons produced from the decays of the $W$ bosons, but hadrons and photons as well. Separate analyses have shown, with both semi-analytical fits and fits using Galprop, that the hadronic fluxes are in close accord with the PAMELA data. The precise nature of the fits depend sensitively on the halo/diffusion models and the proper handling of cosmic backgrounds. Recent preliminary analyses on the photon spectrum indicates that a pure wino may be constrained by the production of photons. If this is the case, this may represent a test of pure wino dark matter, however, as discussed in the previous sections non-negligible non-wino components can play an important role in the size of the $\langle \sigma v \rangle$ which is a sensitive factor in the flux predictions.

### 1.6. Dark Matter and the LHC

Our understanding of what we expect to observe at the LHC is tied closely to the relations between the LHC space of signatures and their possible degeneracies. In addition, there has been recent effort to connect the LHC space of signals with dark matter signatures in broad classes of models. The above connection, dark matter and the LHC, becomes rather relevant in the context of models which predict light gauginos and in particular light gauginos.
gluinos (see i.e. \(^{31,50}\)). In such cases, the production of colored sparticles via gluino production\(^{49}\) becomes the central source of high PT jets and missing energy and can lead to early discovery prospects at low luminosity.

Specifically in the context of a wino LSP, collider implications have been discussed in some depth in the literature (see i.e. \(^{8,46,52-59}\)). The near degeneracy of a pure wino with the lightest chargino yields a mass splitting of order 160 MeV making its discovery prospects at hadron colliders challenging. The wino decay can lead a chargino and pion giving rise to a displaced vertex of a track length of a few centimetres. In association, dilepton signals with displaced vertices have been emphasized in.\(^{51}\) Typically however the role of colored sparticle production in models with a wino-like LSP has been put on a back burner. It has recently been observed however, that light gluinos can arise when the LSP has a significant wino component and when the scalars of the theory decouple.\(^{6}\) Here one finds dominant LHC production modes are

\[
pp \to [(\tilde{g}\tilde{g}), (\tilde{W}\tilde{C}_1), (\tilde{C}_1^\pm, \tilde{C}_1^\mp)]
\]  

and the 3 body decay modes lead to rich jet and missing \(E_T\) signatures translating into early discovery reach of the gluino as low as 500 GeV with \(\sqrt{s} = 7\) TeV for integrated luminosity in the neighbourhood of 800 pb\(^{-1}\) with LSP winos in the \(m_{\tilde{W}} \in [170, 200]\) GeV (as motivated by the reported PAMELA data).\(^{46}\) In recent works, the role of strong produced superpartners at the LHC in models with a wino-like LSP has been emphasized\(^{8}\) inspired by these recent observations that a wino can fit the PAMELA data\(^{3,8}\) and it has been found that a wino-like LSP can occur with a compressed spectrum, in particular for both the squarks and gluino.\(^{8}\) The lighted colored superpartners can produce stunning monojet, multijet and lepton signals with less than fb\(^{-1}\) of data when the mass splitting between the LSP and chargino opens up due in part to a non-negligible Higgsino component. The tagging of bjets in these models becomes very relevant as the three body decay of the \(\tilde{g} \to \tilde{W}b\bar{b}\) can dominate and the reconstruction of the invariant mass of the 2 bjet system has been shown to reveal a potential clue to the size of the gluino mass from its inferred kink in the spectrum. Indeed in greater generality, the gluino three body decay modes

\[
\tilde{g} \to b\bar{b}\tilde{W}, \quad \tilde{g} \to q\bar{q}\tilde{W} \quad \tilde{g} \to t\bar{t}\tilde{N}_2,
\]

\[
\tilde{g} \to tb\tilde{C}^- + h.c., \quad \tilde{g} \to q_u\bar{q}_d\tilde{C}^- + h.c,
\]

(1.22) (1.23)

can all become important depending on the part of the parameter space and loop decays of gauginos can also play a role.\(^{60}\)
1.7. Concluding Remarks

The implications of a wino-like LSP in connection with PAMELA\textsuperscript{3,8} and in connection with the LHC have recently been studied in some detail.\textsuperscript{8,46} In particular the rather exciting possibility of the observable effects from the production of light gluinos at the LHC and a possible interpretation of dark matter in the PAMELA Satellite data represent two prominent discovery channels for the identification of the existence of superpartners. A third related indication of wino-like dark matter could come from an enhanced spin independent cross section when the wino content is supplemented by non-negligible sources of Higgsino and bino content\textsuperscript{8} producing cross sections in the physically interesting region of $\sim \mathcal{O}(10^{-44})\text{cm}^2$.

The LHC could open a paradigm shift in our understanding of the nature of physics beyond the Standard Model of particle physics. Within the plethora of models that have been proposed, the clearest and most well motivated generalization of the Standard Model is clearly softly broken supersymmetry. The eigen content of the LSP plays a central role in determining observable manifestations of SUSY in forthcoming experiments. The LSP wavefunction indeed plays a central role in governing the spin independent cross section of LSPs scattering off of target nuclei, as well as the size of the annihilation cross sections of the LSPs in the early universe and in the galactic halo. The eigen content also implies certain relations on the amount of LSP missing energy produced in hadron collisions, as well as relations amongst the spectrum of the heavier unstable supersymmetric states under the constraints of radiative electroweak symmetry breaking and therefore the possible decays and productions modes of superpartners.

We have entered a data-rich period and some answers to the nature of the LSP are being directly tested. Shortly, it may indeed be possible to answer the seemingly simple question: what is the nature of the lightest supersymmetric particle? Here we have emphasized several theoretical and phenomenological motivations that indicate the distinct possibility that the LSP has a sizeable wino component; and should this be the case, we can expect to constrain or discover models which do produce a wino-like LSP in the very near future.
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References

12. For a detailed discussion see the Chapter by S. Watson.
20. For a detailed discussion see the Chapter by J. Wells.