SUPERSYMMETRIC DARK MATTER

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A review of supersymmetric dark matter in minimal supergravity unification with R-parity invariance and with radiative breaking of the electro-weak symmetry is given. The analysis shows the lightest neutralino is the LSP over most of the parameter space of the supergravity model. The event rates in neutralino-nucleus scattering in dark matter detectors are also discussed. It is found that the event rates are sensitive to the constraint from the $b \rightarrow s\gamma$ experiment. It is also found that the event rates are sensitive to the constraints of relic density and in our analysis we have used the accurate method for the computation of the neutralino relic density. Finally, the effect of the new results on quark polarizabilities, from the data of the Spin Muon Collaboration, on event rates is also discussed. The analysis shows that the event rates for the Ge detectors and for other detectors which use heavy targets are only negligibly affected.

1. Introduction

Considerable evidence for the presence of dark matter in the universe exists: in our galaxy, in other galaxies and in galactic clusters. The rotation curves of luminous matter in spiral galaxies, point to massive amounts of non-luminous matter in the halo of galaxies and provide perhaps the strongest evidence for the existence of dark matter\(^1\). There are many possible candidates for dark matter both in particle physics and in astronomy. Thus in astronomy possible candidates for dark matter are Jupiters, brown dwarfs, neutron stars, black holes etc, while in particle physics one has the possibility of axions, neutrinos, sneutrinos, neutralinos etc, An important constraint arises from the fact that not all of the dark matter in the Universe can

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be baryonic in nature. First, the baryonic dark matter is constrained severely from analysis of nucleosynthesis which show $\Omega_B \leq 0.1$. Second, the recent results of the MACHO Collaboration\(^2\), and from EROS\(^3\) indicate that at best only $20 - 30\%$ of the halo of galaxies is composed of MACHO’S ( Massive Compact Halo Objects), and thus the remainder must be non-baryonic dark matter (NBDM). The non-baryonic dark matter could be either hot (HDM) or cold (CDM). The HDM could be one of the neutrino species (most likely possibility is the tau neutrino $\nu_\tau$), while the CDM could be either an axion, or a SUSY particle\(^4\) (a sneutrino $\tilde{\nu}$ or a neutralino $\chi$). In the following we shall pursue the SUSY possibility and assume that CDM is either a $\tilde{\nu}$ or a $\chi$. Actually it turns out that in supergravity unification\(^5\), with radiative breaking of the electro-weak symmetry, the lightest neutralino turns out to be the LSP over most of the parameter space of the model\(^6\). Thus the model actually predicts the neutralino to be the CDM\(^6\). Further, we assume a mix of cold and hot dark matter in the ratio of 2:1 as indicated by the COBE data. The quantity that appears in theoretical analyses is $\Omega_h^2$, where $h$ is the Hubble parameter in units of $100Km/sMpc$. Currently the experimental uncertainty in $h$ is given by\(^7\) $0.82 \pm 0.17$ (Freedman et. al) ; $0.53 \pm 0.05$ (Sandage et. al). For the purpose of our analyses here we shall assume an $h$ in the range $0.4 \leq h \leq 0.8$ consistent with the above data. Then assuming that $\Omega_B = 0.1$, and $\Omega_{CDM} : \Omega_{HDM}=2:1$ one finds that

$$0.1 \leq \Omega_h^2 \leq 0.4$$ \hspace{1cm} (1)

As mentioned above there are two neutral states, the lightest neutralino $\chi$ and the sneutrino $\tilde{\nu}$, which are possible candidates for CDM in supersymmetric models, for example, the MSSM. However, in the MSSM there are many arbitrary parameters and the model does not predict what the LSP is. The situation is radically different in supergravity unified models\(^5\). Here the parameter space is five dimensional and reduces to a four dimensional space after fixing the Z-mass using radiative breaking of the electro-weak symmetry. The parameter space of the model is then characterized by $m_0, m_{1/2}, A_0, \tan\beta$ where $m_0$ is the universal scalar mass, $m_{1/2}$ is the universal gaugino mass, $A_0$ is the universal trilinear coupling at the GUT Scale and $\tan\beta = v_2/v_1$ where $v_2$ gives mass to the top quark and $v_1$ gives mass to the bottom quark. We limit the parameter space by the fine tuning criterion $m_0, m_{gluino} \leq 1TeV$ where $m_{gluino} = (\alpha_3/\alpha_G)m_{1/2}$ and also limit $\tan\beta$ (since $\tan\beta$ is the ratio of two VeVs, a large $\tan\beta$ also implies a finetuning) so that $\tan\beta \leq 20$. In this domain, we find that the lightest neutralino is also the LSP over most of the allowed region.

### 2. Neutralino Relic Density

The neutralino relic density at current temperatures is given by\(^8\)

$$\Omega h_0^2 = 4.75 \times 10^{-40}(T_{Z1}/T_\gamma)^3(T_\gamma/2.75)^3N_{F}^{1/2}(GeV^{-2}/J(x_f))g/cm^3$$ \hspace{1cm} (2)
Here $T_\gamma$ is the current photon temperature, $n_F$ is the effective degrees of freedom computed at the freeze-out temperature, $(T_\chi/T_\gamma)^3$ is a reheating factor, and $J(x_f)$ is given by

$$J(x_f) = \int_0^{x_f} <\sigma v> \, dx$$  \hspace{1cm} (3)

$<\sigma v>$ is the thermal average of the neutralino annihilation cross section and $x \equiv kT/m_{\tilde{Z}_1}$. Now $J$ receives contributions from Z-exchange, Higgs exchange and from the s-fermion exchange in the t-channel. For the s-fermion exchange the conventional approximation of expanding $<\sigma v> = a + bv^2$ in the integrand in eq. (3) is valid and we use this approximation. However, for the Z and Higgs pole terms such an expansion is a poor approximation in the region below the poles. Thus for $J_{\text{Higgs}}$ and $J_{Z\text{-pole}}$ we use a rigorous thermal averaging over the poles. For example, consider the annihilation via the lightest Higgs pole. Here

$$\sigma v = \frac{A_{\text{Higgs}}}{m_{\tilde{Z}_1}^4} \frac{v^2}{((v^2 - \epsilon_h)^2 + \gamma_h^2))}$$ (4)

$$\epsilon_h = \frac{m_h^2 - 4m_{\tilde{Z}_1}^2}{m_{\tilde{Z}_1}^2}$$ (5)

$$\gamma_h = \frac{m_h \Gamma_h}{m_{\tilde{Z}_1}^2}$$ (6)

where $m_h$ is the Higgs mass, and $\Gamma_h$ is the Higgs width.

Computationally eq. (3) implies a double integration over a pole. Since the pole is associated with a very small width numerical integrations are tricky as a sharp resonance can be easily missed. A more reliable procedure is to reduce eq. (3) to a single integral so that

$$J_{\text{Higgs}}(x_f) = \frac{A_{\text{Higgs}}}{2\sqrt{2}m_{\tilde{Z}_1}^4} [I_{1h} + \frac{\epsilon_h}{\gamma_h} J_{2h}]$$ (7)

$$I_{1h} = \frac{1}{2} \int_0^\infty dy y^{-\frac{1}{2}} e^{-y} \log \left( \frac{(4yx_f - \epsilon_h)^2 + \gamma_h^2}{\epsilon_h^2 + \gamma_h^2} \right)$$ (8)

$$I_{2h} = \frac{1}{2} \int_0^\infty dy y^{-\frac{1}{2}} e^{-y} \left[ \tan^{-1} \left( \frac{(4yx_f - \epsilon_h)^2 + \gamma_h^2}{\gamma_h} \right) + \tan^{-1} \left( \frac{\epsilon_h}{\gamma_h} \right) \right]$$ (9)

A similar analysis can be carried out for $J_Z$. What one finds then is that eq. (7) and the similar expression for $J_Z$ give a smooth result on integration over the poles.

3. Analytic Analysis of the Neutralino Composition

The lightest neutralino $\tilde{Z}_1$ is a linear combination of the four neutral states $(\tilde{W}, \tilde{B}, \tilde{H}_1, \tilde{H}_2)$ so that

$$\tilde{Z}_1 = n_1 \tilde{W}_3 + n_2 \tilde{B} + n_3 \tilde{H}_1^0 + n_4 \tilde{H}_2^0$$ (10)
where $n_i (i = 1 - 4)$ are the co-efficient of the components of the $\tilde{Z}_1$ eigenvector and are discussed below. The co-efficients $n_i$ play an important role both in the relic density analyses as well as in the analyses of event rates in neutralino-nucleus scattering. It is useful then to gain an analytic understanding of the parametric dependence of $n_i$ on the basic parameters of the model. The $n_i$ are determined by the neutralino mass matrix

$$M_{\tilde{Z}} = \begin{pmatrix} \tilde{m}_2 & 0 & a & b \\ 0 & \tilde{m}_1 & c & d \\ a & c & -\mu & 0 \\ b & d & -\mu & 0 \end{pmatrix}$$

(11)

where $\tilde{m}_i = (\alpha_i/\alpha_3)m_{3/2}$, $a = M_Z\cos\theta_W\cos\beta$, $b = -M_Z\cos\theta_W\sin\beta$, $c = -M_Z\sin\theta_W\cos\beta$ and $d = M_Z\sin\theta_W\sin\beta$. Here $\tilde{m}_2(\tilde{m}_1)$ are the SU(2)(U(1)) gaugino masses determined by the relation $\tilde{m}_i = (\alpha_i/\alpha_G)m_{1/2}$. The parameter $\mu$ is determined by fixing the Z-mass using radiative breaking of the electro-weak symmetry. It is found that over most of the parameter space $\mu$ is determined to be large i.e. $|\mu^2/M_Z^2| >> 1$. In this domain one can carry out a perturbative expansion of $n_i$ in $M_Z/\mu$. One finds that to second order in $(M_Z/\mu)$ one has

$$n_1 \approx -\frac{1}{2} \frac{M_Z}{\mu} \frac{1}{(1 - \tilde{m}_1^2/\mu^2)} \frac{M_Z}{m_2 - \tilde{m}_1} \sin2\theta_W \left[ \sin2\beta + \frac{\tilde{m}_1}{\mu} \right]$$

(12)

$$n_2 = 1 - \frac{1}{2} \frac{M_Z^2}{\mu^2} \frac{1}{(1 - \tilde{m}_1^2/\mu^2)} \sin^2\theta_W \left[ 1 + \frac{\tilde{m}_1}{\mu} \sin 2\beta + \frac{\tilde{m}_1^2}{\mu^2} \right]$$

(13)

$$n_3 = \frac{M_Z}{\mu} \frac{1}{(1 - \tilde{m}_1^2/\mu^2)} \sin\theta_W \sin\beta \left[ 1 + \frac{\tilde{m}_1}{\mu} \cot\beta \right]$$

(14)

$$n_4 = -\frac{M_Z}{\mu} \frac{1}{(1 - \tilde{m}_1^2/\mu^2)} \sin\theta_W \cos\beta \left[ 1 + \frac{\tilde{m}_1}{\mu} \tan\beta \right]$$

(15)

The expansion of eqs.(12-15) is found to be accurate to $(3 - 5)\%$ over a significant region of the parameter space. From the above one can easily see that $n_2 > n_1, n_3, n_4, |n_3| > |n_4|$ where in getting this result we have used the radiative electro-weak symmetry breaking relation $\tan\beta > 1$. We note that eq. (13) implies that the neutralino is mostly a Bino in the scaling limit. However, a note of caution is needed in that one should not take the $|M_Z/\mu| \to \infty$ limit. This limit is dangerous since the gaugino-higgsino interference terms which are proportional to $(n_1, n_2),(n_3, n_4)$ vanish in this limit. In realistic analyses, as in the computation of the coherent part of neutralino-nucleus scattering, such terms make significant contributions and cannot be set to zero.

4. Neutralino Detection via Neutralino - Nucleus Scattering

Various possibilities for the detection of neutralino dark matter have been dis-
cussed in the literature. For example, annihilation of neutralinos in the galactic halos can produce an observable signal, i.e., \( \chi + \chi \rightarrow A + X \), where \( A \) can be an energetic gamma ray, positron or an anti-proton\(^4\). However, the backgrounds in these processes are rather significant so this process does not seem very encouraging for the detection of the neutralino. A second possibility is that the halo neutralinos are captured in the center of earth and sun, annihilate and produce upward moving neutrinos (and muons) in detectors on earth\(^{14,15,16}\). The background in these processes are significantly reduced due to the angular windows around earth and sun. Current estimates, however, show that one needs around \( 10^3 - 10^4 m^2 \) neutrino telescopes to see any significant effect. The telescopes currently being planned aim to approach \( O(10^3 m^2) \) area. So once again this mode of detection also does not appear very optimistic for the neutralino dark matter. The most optimistic mode for detection of neutralinos appears to be scattering of neutralinos off nuclei in cryogenic detectors\(^1,4,19,13,14,17,18\). Several detectors using this mode are in various stages of development. We shall focus on this mode of detection for the rest of the talk. More recently there has also been a discussion of detection of neutralinos via atomic excitations\(^20\), but at the moment this possibility requires more investigation. The prime detector in neutralino-nucleus scattering is the quark and the effective interaction that governs the neutralino-quark scattering consists of a spin-dependent (incoherent) part and a spin-independent (coherent) part, and is given by\(^4\)

\[
\mathcal{L}_{\text{eff}} = \tilde{Z}_1 \gamma_\mu \gamma_5 \tilde{Z}_1 \bar{q} \gamma^\mu (A_q P_L + B_q P_R) q + \tilde{Z}_1 \tilde{Z}_1 m_q \bar{q} C_q q
\]

Here \( A_q, B_q \) are the spin-dependent amplitudes which arise from the s-channel Z-exchange and the t-channel squark exchange, and \( C_q \) is the spin-independent amplitude arising from the s-channel Higgs exchange and t-channel squark-exchange. Realistically, of course the quarks are bound inside nuclei so a reasonable amount of nuclear physics enters in the analysis. The total event rate is given by\(^4\)

\[
R = [R_{\text{coh}} + R_{\text{inc}}] \left[ \frac{\rho_{\tilde{Z}_1}}{0.3 \text{ GeV cm}^{-3}} \right] \left[ \frac{\langle v_{\tilde{Z}_1} \rangle}{320 \text{ km/s}} \right] \text{ events/kg da (17)}
\]

\[
R_{\text{coh}} = \frac{16 m_{\tilde{Z}_1} M_N^3 M_Z^4}{[M_N + m_{\tilde{Z}_1}]^2} |A_{\text{coh}}|^2
\]

\[
R_{\text{inc}} = \frac{16 m_{\tilde{Z}_1} M_N}{[M_N + m_{\tilde{Z}_1}]^2} \lambda^2 J(J + 1) |A_{\text{inc}}|^2 , \quad (19)
\]

In the above \( \rho_{\tilde{Z}_1} \) is the local density of dark matter \( v_{\tilde{Z}_1} \) is the relative velocity, \( M_N \) is the nucleus mass, \( J \) is the nucleus spin and \( \lambda \) is defined by \( \lambda < N |\vec{J}|N > \) where \( \vec{S}_i \) is the nucleon spin. From eqs. (18) and (19) one finds that for large \( M_N \) one has \( R_{\text{coh}} \sim O(M_N) \), \( R_{\text{inc}} \sim O(\lambda^2 J(J + 1)/M_N) \). Thus in principle one has two qualitatively different types of detectors, i.e., those with large \( M_N \) and
those with large values of $\lambda^2 J(J + 1)$. In practice there is seldom a case where $R_{coh}$ is negligible and realistic analyses even for light target material (e.g., $^3$He, CaF$_2$) require inclusion of both $R_{inc}$ and $R_{coh}$. For heavy targets (e.g., Ge, NaI, Pb) $R_{inc}$ is typically small, i.e. only a few percent of the total $R$.

We note that there are uncertainties in the evaluation of both $R_{inc}$ and $R_{coh}$. Uncertainties in $R_{inc}$ arise due to experimental uncertainties in the determination of $\Delta q$ on which $R_{inc}$ depends, where $\Delta q$ are the quark polarizabilities defined by

$$< p(n) | \bar{q} \gamma_\mu \gamma_5 q | p(n) > = S_\mu^{p(n)} \Delta q, \text{ where } S_\mu^{p(n)} = (0, \vec{S}_\mu^{p(n)}) \text{ is the nucleon spin.}$$

We shall discuss later the sensitivity of the results to the determination of $\Delta q$. There is also an uncertainty in the determination of $R_{coh}$. This arises due to the uncertainty in the determination of $s$-quark matrix elements

$$< n | m_s \bar{s} s | n > = f_s M_n \text{ that enter in the computation of } C_s. \text{ Currently } f_s \text{ has a significant uncertainty, about 50%, which can lead to an uncertainty of } 0(30 - 50)% \text{ in the determination of } R_{coh}.$$

Next we discuss briefly the relative contribution of the heavy neutral Higgs to the coherent part of the scattering. Naively one might expect the heavy Higgs contribution to be negligible since it would be suppressed by the heavy Higgs ($mass)^2$ while the light Higgs contribution is suppressed only by the light Higgs ($mass)^2$. However, this assessment is correct for the up quark but not for the down as can be seen by the expression for $C_q$ below:

$$C_q^{Higgs} = \frac{g_2^2}{4M_W} \left[ \left\{ \begin{array}{c}
\cos \alpha \\
\sin \beta
\end{array} \right\} \frac{F_h}{m_H^2} - \left\{ \begin{array}{c}
\sin \alpha \\
\cos \beta
\end{array} \right\} \frac{F_{H}}{m_H^2} \right] \text{ u-quark}
$$

$$+ \left\{ \begin{array}{c}
\cos \alpha \\
\sin \beta
\end{array} \right\} \frac{F_h}{m_H^2} - \left\{ \begin{array}{c}
\sin \alpha \\
\cos \beta
\end{array} \right\} \frac{F_{H}}{m_H^2} \right] \text{ d-quark} \quad (20)$$

where

$$F_h = (n_1 - n_2 tan \theta_W) (sin \alpha n_3 + cos \alpha n_4) \quad (21)$$

$$F_H = (n_1 - n_2 tan \theta_W) (-cos \alpha n_3 + sin \alpha n_4) \quad (22)$$

Here $\alpha$ is the Higgs mixing angle. Now the Higgs mixing angle $\alpha$ is typically small so from eq. (20) one finds that there is a suppression of the d-quark contribution in the light Higgs sector which often can be more than the ($mass)^2$ suppression in the heavy Higgs sector. Thus the heavy Higgs contribution cannot be neglected as it can make a substantial contribution to $C_q$.

5. $b \rightarrow s \gamma$ Branching Ratio Constraint on Event Rates

Recently the CLEO Collaboration obtained the first experimental determination of the photonic penguin process $b \rightarrow s \gamma$. For the inclusive $b \rightarrow s \gamma$ decay they find the result

$$BR(b \rightarrow s \gamma) = (2.32 \pm 0.51 \pm 0.29 \pm 0.32) \times 10^{-4} \quad (23)$$
Now in the SM $b \to s\gamma$ decay receives contributions from the W-exchange. For the supersymmetric case there are additional contributions arising from the exchange of $H^+, \tilde{W}, \tilde{Z}, \text{and } \tilde{g}$. There are many uncertainties in the theoretical evaluation of $b \to s\gamma$. These uncertainties arise from uncertainties in the determination of $\alpha_s$, quark masses K.M matrix elements, and uncertainties arising from the next to leading order (NLO) QCD corrections which can be as large as $0(30)\%$. There is a significant debate in the literature currently regarding what exactly the $BR(b \to s\gamma)$ value is in the SM$^{23}$. An accurate answer to this question can only result after the next-to-leading order(NLO) QCD corrections have been computed reliably. Similar uncertainties are present in the evaluation of $b \to s\gamma$ in SUSY theory. The discussion of $b \to s\gamma$ constraint on neutralino relic density$^{24,25}$ can be facilitated by use of the parameter $r_{\text{SUSY}}^{18}$ which we define by the ratio $r_{\text{SUSY}} = BR(b \to s\gamma)_{\text{SUSY}} / BR(b \to s\gamma)_{\text{SM}}$. Many of the uncertainties discussed above cancel out in the ratio $r_{\text{SUSY}}$. However, we must keep in mind that the NLO corrections would in general be different for the SUSY case than for the SM case. In this analysis, however, we limit ourselves to the leading order evaluation. Analogous to $r_{\text{SUSY}}$ we can also define$^{18}$ $r_{\text{exp}} = BR(b \to s\gamma)_{\text{exp}} / BR(b \to s\gamma)_{\text{SM}}$. The CLEO Collaboration use an SM value of $(2.75 \pm 0.8) \times 10^{-4}$. Using this value and the result of eq. (23) one finds that $r_{\text{exp}}$ lies in the range $r_{\text{exp}} = 0.46 - 2.2$. This range turns out to be a rather strong constraint on SUSY theory if we assume that $r_{\text{SUSY}} = r_{\text{exp}}^{18}$. This is so because in the SUSY case one normally gets a much larger range for $r_{\text{SUSY}}$, i.e., a range of $\approx (0, 10)$.

6. Analysis and Results

We first discuss the event rates without inclusion of the $b \to s\gamma$ constraint. Results on maximum and minimum of event rates are exhibited in Fig. 1 as a function of $m_{\tilde{g}}$ for the target materials: Pb, Ge and CaF$_2$. We note that the maximum curves of the event rates all exhibit a dip when $m_{\tilde{g}}$ is in the mass range $\approx 300 - 450GeV$ independent of the target material. This dip arises$^{18}$ due to the Z-pole and the Higgs pole effects and the relic density constraint. Effectively, rapid annihilation near the Z and Higgs poles leads to values of $\Omega h^2$ which fall below the CHDM limit and have to be eliminated. The eliminated part of the parameter space contains the light SUSY spectrum and the large event rates. Thus one sees a dip in the event rates in the vicinity of $m_{\tilde{Z}_1} \sim M_Z/2$ and $m_{\tilde{Z}_1} \sim m_h/2$. The total event rate $R$ is generally dominated by $R_{\text{coh}}$ for all except the lightest target materials. Now $R_{\text{coh}}$ depends on the gaugino-higgsino interference term which is proportional to $(n_1, n_2) \times (n_3, n_4)$. Using the behavior of $n_i$ with large $\mu$ and the dependence of $R_{\text{coh}}$ on the gaugino-higgsino interference, we can understand the behavior of $R$ for large $m_{\tilde{g}}$. Typically $\mu$ is an increasing function of $m_{\tilde{g}}$, and the $n_i (i=1,3,4)$ are decreasing functions as $\mu$ increases. Thus $R_{\text{coh}}$ falls monotonically$^{18}$ with $m_{\tilde{g}}$ beyond the dip as can be seen in Figs. 1. (A similar analysis holds as a function of $m_0$). Now it is easily seen
that for $m_{\text{gluino}} \geq 650\text{GeV}$ for $\mu < 0$ and $m_{\text{gluino}} \geq 700\text{GeV}$ for $\mu > 0$ the event rates even for heavy targets (e.g. Pb) fall below the level of 0.01 event/kg.d. This is the level of sensitivity that detectors in the next 5-10 yrs hope to achieve.

Next we discuss the effect of the $b \to s\gamma$ constraint on dark matter. The quantitative effects of $b \to s\gamma$ depend on the value of $r_{\text{SUSY}}$. Generally, the $b \to s\gamma$ constraints is more severe for $\mu > 0$ than for $\mu < 0$. Similarly, one finds a significant effect on the event rates for $\mu > 0$, while the effect for $\mu < 0$ is relatively smaller. Results for the case $r_{\text{SUSY}} \leq 1.33$ is shown in Fig. 2 for the cases $\mu > 0$ (Fig. 2a) and $\mu < 0$(Fig. 2b). As indicated above the allowed region of the parameter space shrinks significantly for $\mu > 0$ (Fig. 2a) and the maximum allowed event rates also fall. The corresponding effect on $\mu \leq 0$ (Fig. 2b) is significantly smaller.

Finally we discuss the effect of the variations in quark polarizabilities $\Delta q$ on the event rates. The part sensitive to $\Delta q$ is $R_{\text{inc}}$. The previous determinations of $\Delta q$ using the EMC data gave, $(\Delta u, \Delta d, \Delta s) = (0.77 \pm 0.08, -0.49 \pm 0.08, -0.15 \pm 0.08)$. Recently there has been a reanalysis of $\Delta q$ using new data from SMC which gives, $(\Delta u, \Delta d, \Delta s) = (0.83 \pm 0.03, -0.43 \pm 0.03, -0.10 \pm 0.03)$. These determinations are consistent with each other within 1$\sigma$, but the variations of $\Delta q$, specifically the variation of $\Delta s$, can generate significant changes in $R_{\text{inc}}$. However, $R_{\text{inc}}$ is generally
a small component of the total $R^{18,13}$. We have analysed several target materials, He, CaF$_2$, Ge, Pb etc., and find that except for the lightest target materials, i.e. He and CaF$_2$, $R_{inc}$ is only a few percent of the total. Thus the event rates for targets such as Ge are not significantly affected by the new determination of $\Delta q^{18,13}$.

7. Concluding Remarks

We have given here an analysis of neutralino dark matter within minimal supergravity unification. Remarkably the model predicts that the lightest neutralino is also the LSP over most of the parameter space of the model, and thus the model gives a candidate for cold dark matter. We have analysed the relic density of neutralinos, using the accurate method which integrates over the Z and Higgs poles in thermal averaging of $\sigma v$. These effects are found to be significant for values of gluino masses $\leq 400\text{GeV}$. We have analysed the event rates in neutralino-nucleus scattering and find that there is a significant region of the parameter space where event rates $\geq 0.01$ event/kg.d are predicted for targets such as Ge and Pb. This region of the parameter space will be accessible to current and future technologies over the next 5-10 years. However, more sensitive detectors, two to three orders of magnitude more sensitive, are needed to probe most of the parameter space of the minimal supergravity model.

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