Astroparticle Physics and Colliders

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Abstract

In this talk I discuss the interplay between collider physics and four topics of astroparticle physics: neutrino oscillations, electroweak baryogenesis, LSP Dark Matter, and ultra–high energy cosmic rays (UHECR). Some astrophysical scenarios can (only) be tested decisively at colliders. In other cases input from collider experiments is required to sharpen predictions for future astro–particle physics experiments, e.g. for the LSP detection rate or the UHECR spectrum in top–down models.
1) Introduction

The organizers of this conference asked me to give a review talk on what is commonly known as astro–particle physics. This describes a vast field of research. I obviously have neither the time nor the expertise to cover most of it adequately. I therefore decided to narrow the focus on some topics that have, or at least may have, fairly direct connections to experiments (to be) performed at high energy particle colliders. This not only matches the main theme of this conference, it also suits my personal interests. In the following four sections I will therefore describe lepton flavor violation (LFV) as predicted in certain supersymmetric models of neutrino masses that can explain the observed pattern of neutrino flavor oscillations; electroweak baryogenesis, which does not work in the Standard Model (SM) but may work in its supersymmetric extension, the MSSM; the LSP as cold Dark Matter candidate; and the “top–down” interpretation of ultra–high energy cosmic rays, which posits that these events originate from the decay or annihilation of ultra–massive particles.

2) Neutrino Oscillations

It may be appropriate to begin this survey by reminding ourselves that astro–particle physics experiments, most notably SuperK [1], succeeded in committing the murder which a great many experimenters working at colliders had spent twenty–odd years attempting: by unambiguously showing that muon (anti)neutrinos oscillate into other flavors, most likely into $\bar{\nu}_\tau$, the minimal SM was killed, since it predicts vanishing neutrino masses and hence separately conserved $e$, $\mu$, $\tau$ lepton numbers. This is a constructive proof that astro–particle physics experiments can produce results which impact “mainstream” particle physics. Similarly, recent SNO data [2] have shown that $\bar{\nu}_e$ change into something else somewhere between the Sun and solar neutrino detectors on Earth, which implies that electron number is not conserved. This is again most naturally explained by neutrino flavor oscillations [3], but in this case other explanations are still possible [4] (since no “$L/E$” plot showing the characteristic energy dependence of neutrino oscillations has yet been established). We now know that some neutrinos are massive, and that none of the three lepton flavors is conserved.

These discoveries are undoubtedly of great importance both conceptually and for the planning of future Earth–based neutrino oscillation experiments. However, its relevance to collider physics is a priori much less clear. In the framework of the SM, small neutrino masses can most naturally be accommodated by introducing a dimension–5 term in the Lagrangian, of the form

$$L_\nu = \frac{\lambda_{ij}}{M} L^c_i H L_j H + h.c.$$ (1)

Here $H$ is the Higgs doublet, the $L_i$ are left–handed lepton doublets ($i$, $j$ being flavor labels), $M$ is a large mass scale, and the $\lambda_{ij}$ are numerical coefficients. The data can be accommodated by allowing $\lambda$ to have non–vanishing off–diagonal entries, which clearly violates lepton flavor. This LFV also manifests itself in processes of interest to collider physicists. For example, the large $\nu_\mu - \nu_\tau$ mixing implied by atmospheric neutrino oscillations gives rise to $\tau \to \mu \gamma$ and $Z \to \mu^+\mu^-$ decays through the one–loop diagrams shown in Fig. 1. However, explicit evaluation of these diagrams reveals that they are suppressed by powers of the neutrino mass, leading to predictions for the branching ratios of these LFV decays many orders of magnitude below the sensitivity of conceivable experiments. This result can be generalized: In the SM
augmented by the term (1), or by tiny Yukawa couplings giving rise to Dirac neutrino masses, LFV at colliders is unobservably small \[5\].

\[
\begin{align*}
\gamma, Z & \rightarrow \mu \gamma \\
\tau & \rightarrow \mu, \mu & \tau
\end{align*}
\]

**Fig.1:** Diagram giving rise to $\tau \rightarrow \mu \gamma$ and $Z \rightarrow \mu \tau$ decays in the SM augmented by the dimension–5 term (1). The external boson ($\gamma$ or $Z$) has to be appended in all possible ways, and the $\times$ denotes a $\nu_\tau \leftrightarrow \nu_\mu$ flavor transition. These diagrams are suppressed by powers of $m_\nu/M_W$ and are thus unobservably small.

The situation can be quite different in supersymmetric extensions of the SM, if the scale of LFV is below the energy scale where SUSY breaking is transmitted to the visible sector (i.e., where “ordinary” sparticles obtain their masses) \[7\]. Consider a supersymmetric seesaw model \[6\] with three $SU(2) \times U(1)_Y$ singlet superfields $N_i$ appearing in the superpotential via

\[
W_N = f_{ik} H_u L_i N_k + \frac{1}{2} M_{kl} N_k N_l,
\]

where $H_u$ is the Higgs doublet coupling to up–type quarks. At energies below the smallest eigenvalue $M_N$ of the matrix $M$, the three singlets can be integrated out, producing a term of the form (1) in the effective Lagrangian (with $H$ replaced by $H_u$). The coefficients $\lambda$ of eq.(1) are then bilinear functions of the new Yukawa couplings $f$ appearing in eq.(2), leading to neutrino masses of order

\[
m_\nu \sim |f|^2 \langle H_u^0 \rangle^2 / M_N^2.
\]

The crucial difference between the SM and the MSSM is that the latter contains additional parameters (besides the $\lambda_{ij}$) that can violate lepton flavor: the soft breaking masses and trilinear $A$ parameters in the slepton sector. High–scale LFV therefore does not necessarily decouple in the MSSM. In particular, even if we assume that slepton masses are universal at some high scale $M_U > M_N$, with value $m_{\tilde{L}}$, the LFV violation encoded in the $f_{ij}$ will radiatively induce LFV slepton masses of order \[8\]

\[
\delta \left( m_{\tilde{L}}^2 \right)_{ij} \sim \frac{1}{8 \pi^2} \sum_k f_{ki} f_{kj} \left( 2 m_{\tilde{L}}^2 + m_{H_u}^2 + |A_{\tilde{L}}|^2 \right) \log \frac{M_U}{M_N};
\]

the exact (1–loop) RGE can be found in \[8\]. Eqs.(3) and (4) together imply that this new source of LFV scales like $m_\nu M_N/M_W^2 \cdot \log(M_U/M_N)$. For given neutrino mass this becomes maximal for $M_N$ only slightly below $M_U$. In other words, this new physics effect increases with increasing energy scale where the effect originates!

These flavor off–diagonal slepton masses can contribute to $Z \rightarrow \mu^\pm \tau^\mp$ and $\tau \rightarrow \mu \gamma$ decays through the one–loop diagrams shown in Fig. 2 \[7\]. There is again some GIM–type cancellation; however, the relevant suppression factor is now $\left( m_{\tilde{L}}^2 \right)_{\mu \tau} / m_{\tilde{L}}^2$, which is usually much bigger than
This suppression is nevertheless crucial for bringing supersymmetric see–saw models with large neutrino mixing angles and high–scale SUSY breaking into agreement with stringent experimental limits on $\tau \rightarrow \mu \gamma$ and $\mu \rightarrow e \gamma$ decays.

\[
\begin{align*}
\gamma, Z \hskip 0.5cm \tau \hskip 0.5cm & \xrightarrow{\chi_0^0} \hskip 0.5cm \tau_L \hskip 0.5cm \mu_L \hskip 0.5cm + \hskip 0.5cm \gamma, Z \\
\hskip 0.5cm \tau \xrightarrow{\chi_i} \hskip 0.5cm \nu_\tau \hskip 0.5cm \mu \hskip 0.5cm & \xrightarrow{\chi_i} \hskip 0.5cm \nu_\mu
\end{align*}
\]

**Fig.2:** SUSY contributions to $\tau \rightarrow \mu \gamma$ and $Z \rightarrow \mu \tau$, where $\times$ now denotes a slepton flavor transition. These contributions are only suppressed by powers of $(\delta m_\mu^L) / m_\mu$. 

This suppression of loop–induced LFV decays may also offer the chance to observe tree–level LFV effects in sparticle production and decay at high energy colliders. For example, one can search for sneutrino pair production at lepton colliders, followed by $\nu_i \rightarrow \ell_k \tilde{\chi}_1^0 \rightarrow \ell_k jj \tilde{\chi}_1^0$ decays. This can give rise to final states containing $\mu^+ \tau^-$ pairs together with up to four jets and missing energy/momentum, which have very little background [8]. These LFV signals can be understood as being due to Majorana slepton oscillations [4]. They are therefore only suppressed by powers of $\delta m_L / \Gamma_L$, where $\Gamma_L \sim \alpha \overline{m}_L$ is the average decay width of the oscillating sleptons. Since the largest new Yukawa couplings $f$ are assumed to (again) occur in the third generation, and $\nu_\tau$ is known to strongly mix with $\nu_\mu$ but only weakly with $\nu_e$, for given luminosity $\mu^+ \mu^-$ colliders offer better sensitivity for these $\tilde{\nu}$ oscillation signals than $e^+ e^-$ colliders do [8]; however, it now seems likely that a muon collider, if it can be built at all, will have far smaller luminosity than that currently foreseen for linear $e^+ e^-$ colliders.

Slepton oscillations can also be tested in slepton production at hadron colliders [10, 11], as well as through the production of heavier neutralinos and charginos followed by their leptonic decay [10, 11, 12]. These signals can probe regions of parameter space that are compatible with bounds on LFV $Z$, $\tau$ and $\mu$ decays. Moreover, analysis of such signals allows (at least in principle) to determine the quantities $f_{ij}$ and $M_{ij}$ separately, rather than only the combinations of these parameters that produces the effective neutrino mass term (1) [13], if the slepton masses at scale $M_U$ are known.

However, all these SUSY–mediated LFV effect vanish if the scale of LFV is higher than the scale where SUSY breaking is mediated to the visible sector, as is the case e.g. in most see–saw models with gauge–mediated supersymmetry breaking. On the other hand, supersymmetry also permits an entirely novel way to generate neutrino masses, if $R$–parity is broken. In general one neutrino will become massive through mixing with neutralinos, while the other neutrino masses are generated through loop diagrams. Many models of this kind predict relations between neutrino oscillations and decays of the lightest superparticle (LSP), where the latter may include LFV modes; this again offers the possibility to detect LFV at colliders. I will not discuss this approach here, since it has been covered at this conference by E.J. Chun [14].
3) Electroweak Baryogenesis

The Universe around us contains baryons and leptons, but very few antibaryons and antileptons. Analyses of Big Bang nucleosynthesis [15] show that the baryon to photon ratio,

\[ \eta \equiv \frac{n_b - n_\bar{b}}{n_\gamma}, \] (5)

must be a few times \(10^{-10}\). In other words, there is a very small, but non–vanishing, excess of particles over antiparticles in the Universe. This is not an existence proof of “new physics”, since such a small asymmetry could simply be imposed as an initial condition on the Universe. However, a dynamical explanation of such a small excess, starting from a Universe with vanishing total baryon and lepton number, would certainly be more attractive.

In his seminal 1967 paper [16], Sakharov listed three necessary conditions any dynamical mechanism must satisfy: it must (obviously) violate baryon number; it must violate both C and CP, since the baryon number generator \(B\) is odd under both C and CP, so that \(\langle 0|B|0 \rangle\) would have to vanish if \(C\) or \(CP\) were good quantum numbers; and it must be active out of thermal equilibrium, since otherwise the rate for any reaction \(A \rightarrow B\) equals that of the inverse reaction \(B \rightarrow A\), making it impossible to create a net baryon number.

Somewhat later it was realized [17] that a fourth condition has to be satisfied. The reason is that the SM predicts electroweak baryon number violating “sphaleron” transitions. Roughly speaking, these correspond to transitions between degenerate (ground) states with different “topology” [in the space of \(SU(2)\) gauge configurations]. These states are separated by “potential barriers” whose height is given by the vev that breaks the electroweak gauge symmetry. At zero temperature these transitions are suppressed by a tunneling factor \(\exp(-16\pi^2/g_2^2)\), where \(g_2\) is the \(SU(2)\) gauge coupling, and can thus safely be neglected. However, these transitions are in thermal equilibrium at temperatures \(T \geq 100\) GeV. Note that these transitions violate both baryon and lepton number, but conserve \(B-L\). The fourth condition that any dynamical explanation of the baryon asymmetry of the Universe (BAU) must satisfy in order to be “immune” to sphaleron transitions is therefore the following: it must either produce a non–vanishing \(B-L\), or else produce non–vanishing \(B\) at temperature (well) below 100 GeV.

Of course, the SM also violates C and CP. Together with the sphaleron–induced violation of baryon number, this satisfies two of the three Sakharov conditions. Moreover, the third condition, deviation from thermal equilibrium, is satisfied during phase transitions.† Such a phase transition might be associated with electroweak symmetry breaking. To see this, consider the temperature dependent effective potential of the Higgs field\‡\(\phi\):

\[ V_{\text{eff}}(\phi, T) \simeq \frac{\gamma T^2 - m^2}{2} \phi^2 - \frac{\lambda}{2} \phi^4. \] (6)

Here \(-m^2\) (the infamous negative mass–squared) and \(\lambda\) describe the potential at \(T = 0\), while \(\gamma\) and \(E\) describe the leading effects of the thermal plasma. The qualitative behavior of this potential is illustrated in fig. 3. Since \(\gamma\) and \(E\) are both positive, at sufficiently high

\*The small \(e^+\) and \(\bar{p}\) components in the primary cosmic ray flux have most likely been produced in the collisions of relativistic \(e^-\) and \(p\) with ordinary matter.

†For example, boiling water is a non–equilibrium process. This can be seen from the fact that boiling water behaves quite differently from condensing steam, even if temperature and pressure are the same.

‡\(\phi\) only describes the field acquiring a vev, e.g. the real part of the neutral component of the Higgs doublet in the SM, or a linear combination of the two analogous real parts in the MSSM.
temperature the origin ($\phi = 0$) will be a minimum, not a maximum. Indeed, for $T > T_c$ it will be the absolute minimum, i.e. the $SU(2) \times U(1)_Y$ symmetry will not be broken. However, a second minimum may exist at $\phi = \phi_0 \neq 0$, which becomes more prominent as $T$ decreases. At some critical temperature $T_c$ this second minimum satisfies $V(\phi_0, T) = 0$, i.e. it is degenerate with the minimum at the origin. It is easy to show that this requires

$$\phi_c \equiv \phi_0(T = T_c) = \frac{ET_c}{\lambda}.$$  \hspace{1cm} (7)

At slightly lower temperature the Universe begins to tunnel from $\phi = 0$ to $\phi = \phi_0$. The rate of (thermal) sphaleron transitions at this tunneling temperature depends very sensitively on $\phi_0$ as well as $T$. It can be shown that these transitions are sufficiently suppressed if

$$\phi_c \gtrsim T_c,$$  \hspace{1cm} (8)

which requires $E \gtrsim \lambda$ from eq.(7). Recall that $\lambda$ is proportional to the square of the mass of the (lightest CP–even) Higgs boson. The experimental lower bound on the mass of (SM–like) Higgs bosons therefore translates into a lower bound on $\lambda$. Moreover, $E$ receives significant contributions only from bosons with mass not greatly in excess of $T$. In the SM, which contains few elementary bosons, condition (8) is therefore badly violated for all experimentally allowed values of $\lambda$; indeed, here $E/\lambda$ is so small that no proper phase transition occurs \[18\].

Fig.3: The temperature–dependent effective potential at temperatures above, at, and below the critical temperature $T_c$.

The situation is somewhat more favorable in the MSSM. The experimental lower bound on $\lambda$ is (slightly) weaker than in the SM. More importantly, there are many more elementary scalars in the theory. The potentially most important ones are the $\tilde{t}$ squarks, due to their large (Yukawa) couplings to $\phi$. Recall, however, that $\tilde{t}$ can only contribute significantly to $E$
if $m_{\tilde{t}}$ does not greatly exceed $T_c$. At the same time, some $\tilde{t}$ squarks must be quite heavy to produce sufficiently large loop corrections to the mass of the lightest CP-even Higgs boson, raising it from the tree-level upper bound of $M_Z$ to at least the experimental lower bound. Moreover, electroweak precision data imply lower bounds of several hundred GeV on the masses of $SU(2)$ doublet $\tilde{t}$ and $\tilde{b}$ squarks. This leads one to consider scenarios with heavy $SU(2)$ doublet squarks, but light $SU(2)$ singlet $\tilde{t}_R$. Indeed, it can be shown \cite{19} that some region of parameter space satisfying the constraint (8) still exists, with $m_{\tilde{t}_L} \geq 1 \text{ TeV}$, $m_{\tilde{t}_R} \leq m_t$ and $m_h \leq 117 \text{ GeV}$; the latter two predictions can easily be tested at future collider experiments.

The existence of a light $\tilde{t}_R$ and a fairly light Higgs boson ensures that condition (8) is satisfied, which in turn implies that the third Sakharov condition is met. As well known, the second Sakharov condition is satisfied already in the SM, where CP violation originates in complex Yukawa couplings, and manifests itself in the complex phase of the KM matrix. However, it turns out that this source of CP violation is too weak to produce a sufficiently large baryon asymmetry in the present framework. One needs large \cite{20}, perhaps even maximal \cite{21}, CP violation in the chargino mass matrix, together with relatively light charginos. This prediction can also be tested through future precision measurements, e.g. at $e^+e^-$ linear colliders \cite{22}.

Electroweak baryogenesis can thus easily be falsified by future collider experiments. This remains true in extensions of the MSSM which contain $SU(2) \times U(1)_Y$ singlet Higgs superfields \cite{23}; although this makes it somewhat easier to satisfy the constraint (8), the required CP violation in the chargino system is essentially the same as in the MSSM. Unfortunately, other proposed explanations of the BAU cannot be tested in the laboratory. For example, in “leptogenesis” \cite{24} one creates a lepton asymmetry at high $T$, which is partially converted to a baryon asymmetry by electroweak sphaleron transitions. The lepton asymmetry is generated in the production and decay of “right–handed” [$SU(2) \times U(1)_Y$ singlet] (s)neutrinos with mass typically exceeding $10^8$ GeV, well beyond the reach of conceivable colliders. The same heavy neutrinos influence the mass matrix of the light neutrinos through the see–saw mechanism \cite{8}; however, the CP violating phase required for leptogenesis is not necessarily related to the phase that can (in principle) be measured in neutrino oscillation experiments \cite{25} (but in supersymmetric models with universal soft breaking terms at some high scale, it can be measured in low–scale LFV processes \cite{26}). Similarly, the Affleck–Dine mechanism \cite{27} produces a baryon or lepton asymmetry from a sfermion condensate just after the end of inflation; its dynamics is governed by parameters that cannot be measured at colliders. However, all proposed dynamical explanations of the BAU require the existence of new CP violating phases. We just saw that in many cases these phases cannot be measured at collider energies; nevertheless the existence of new sources of CP violation can serve as motivation to search for CP violation wherever experimentally feasible, including channels where no CP violation is expected in the SM.

4) Neutralino Dark Matter

The existence of Dark Matter (DM) in the Universe is quite well established. Historically the first evidence came from the analysis of the motion of gravitationally bound systems. The simplest example of this kind is given by the well–known galactic rotation curves. An object on a circular orbit (with radius $R$) around a spherical galaxy must have orbital velocity $v(R) = \sqrt{G_N M(R)/R}$, where $G_N$ is Newton’s constant and $M(R)$ is the mass inside the
orbit. If a galaxy’s mass was confined to its luminous parts, one should thus find \(v(r) \propto 1/\sqrt{R}\) for sufficiently large \(R\); however, observation yields essentially flat rotation curves, \(v(R) \sim \text{constant}\), which implies that the luminous part of the galaxy is embedded in a much larger dark halo. Analogous arguments can be made for the motion of galaxies inside clusters. More recently, measurements of the anisotropy of the cosmic microwave background on degree scales as well as observations of type–Ia supernovae at high redshift have determined the total matter density of the Universe to fall in the range \([28]\)

\[
0.1 \lesssim \Omega_{\text{matter}} h^2 \lesssim 0.3.
\]

Here, \(\Omega_X\) denotes the mass (or energy) density of component \(X\) in units of the critical density (i.e. \(\sum X \Omega_X = 1\) corresponds to a flat Universe), and \(h\) is the Hubble constant in units of 100 km/(sec\cdot Mpc). \(\Omega_{\text{matter}}\) includes all forms of non–relativistic matter, but analyses of Big Bang nucleosynthesis have shown that \(\Omega_{\text{baryon}} h^2 < 0.05\) \([28]\), so that the range \((9)\) essentially also holds for the DM component.

The lightest neutralino \(\tilde{\chi}_1^0\) is probably the best motivated, and certainly the most widely studied \([29]\), particle physics DM candidate. In many SUSY models \(\tilde{\chi}_1^0\) is the lightest superparticle (LSP) over much of parameter space; it is then absolutely stable, if \(R\)–parity is conserved. It is also neutral, which is required experimentally for any massive particle that contributes significantly to DM; charged particles would bind to nuclei, thereby forming isotopes with exotic mass–to–charge ratios, which have not been found \([30]\). Finally, the present thermal LSP relic density is at least within a few orders of magnitude of the desired range \((9)\).

Let us investigate this last argument in some more detail, since it is often used to single out regions of SUSY parameter space as being “cosmologically favored”. The basic assumption is that \(\tilde{\chi}_1^0\) was in thermal equilibrium with the hot plasma of SM particles after the last period of entropy production, which in simple cosmological models occurred at the end of inflation. This requirement is satisfied if the “re–heat” temperature at the end of this epoch of entropy production exceeded \(0.1 m_{\tilde{\chi}_1^0}\); note that the scale of inflation is supposed to be around \(10^{13}\) GeV, so re–heat temperatures in excess of \(0.1 m_{\tilde{\chi}_1^0}\) are easy to achieve. Neutralinos will then stay in equilibrium as long as the rate of reactions that change the number of LSPs is larger than the Hubble expansion rate. \(R\)–parity conservation implies that LSPs can only be created or annihilated in pairs, so that the equilibrium condition reads

\[
n_{\tilde{\chi}_1^0}(T) \langle v \sigma_{\tilde{\chi}_1^0} \rangle > H(T).
\]

Here, \(n_{\tilde{\chi}_1^0}\) is the LSP number density, \(\sigma_{\tilde{\chi}_1^0}\) the total \(\tilde{\chi}_1^0\) pair annihilation cross section, \(v\) the relative velocity between the two LSPs in the cm frame, and \(\langle \ldots \rangle\) denotes thermal averaging. At temperatures \(< m_{\tilde{\chi}_1^0}\), \(n_{\tilde{\chi}_1^0} \sim (m_{\tilde{\chi}_1^0} T)^{3/2} \exp \left(-m_{\tilde{\chi}_1^0}/T\right)\) is exponentially suppressed, whereas the Hubble parameter in the radiation–dominated era, \(H(T) \sim T^2/M_{\text{Pl}}\), \(M_{\text{Pl}} \simeq 2.4 \cdot 10^{18}\) GeV being the Planck mass, only decreases like a power of \(T\). Clearly the condition \((10)\) can therefore not be satisfied at arbitrarily low temperatures. Due to the \(1/M_{\text{Pl}}\) factor in \(H(T)\), LSPs stay in chemical equilibrium with the SM plasma until they are quite non–relativistic; freeze–out (decoupling) typically occurs at \(T_F \simeq m_{\tilde{\chi}_1^0}/20\). Clearly decoupling will occur later, at a smaller value of \(n_{\tilde{\chi}_1^0}\) for larger values of \(\langle v \sigma_{\tilde{\chi}_1^0} \rangle\). Indeed, in this scenario one finds that the present \(\tilde{\chi}_1^0\) relic mass density is essentially inversely proportional to \(\langle v \sigma_{\tilde{\chi}_1^0} \rangle\); this is quite

*Note that LSPs will stay in kinetic equilibrium with the SM plasma, i.e. will have a thermal velocity distribution, long after they drop out of chemical equilibrium.
into a massless fermion–antifermion pair only from a \( P \)
wave of the sparticle mass scale. Moreover, two identical Majorana fermions can annihilate
at the very least establishes an intriguing coincidence between particle physics and cosmology.

Dimensional analysis shows that \( \sigma \chi^0_1 \) should decrease, i.e. \( \Omega \chi^0_1 h^2 \) should increase, like the
square of the sparticle mass scale. Moreover, two identical Majorana fermions can annihilate
into a massless fermion–antifermion pair only from a \( P \)-wave (or higher) initial state.\(^1\) Since \( \chi^0_1 \)'s are already quite non–relativistic at freeze–out, this reduces the thermal average of \( v \sigma \chi^0_1 \)
by about an order of magnitude. If the LSP is mostly a bino, which is true over most of
parameter space in many SUSY models, the upper bound in (4) can only be satisfied in three
different regions of parameter space [31, 32]:

- **The bulk region:** If \( \chi^0_1 \approx \tilde{b} \), the biggest contribution to \( \sigma \chi^0_1 \) comes from \( \ell^+ \ell^- \) final states,
which in turn receive the most important contributions from \( \tilde{\ell}_R \) (or \( \tilde{\tau}_1 \)) exchange in the \( t– \) or \( u– \)channel \( (\ell = e, \mu, \tau) \). One reason for this is that \( \tilde{\ell}_R \) has the largest
hypercharge, and hence the strongest coupling to \( \tilde{b} \), of all sfermions. Moreover, since
\( \tilde{\ell}_R \) is an \( SU(3) \times SU(2) \) singlet, its mass only receives rather small contributions from
gaugino loop diagrams; such diagrams always increase the sfermion mass, and are e.g.
responsible for the fact that squarks without sizable Yukawa couplings cannot be much
lighter than gluinos. Here \( \Omega \chi^0_1 h^2 \leq 0.5 \) requires [33] \( m_{\chi^0_1}, m_{\ell_R} \leq 250 \text{ GeV} \), which can
readily be tested at a first–generation linear collider.

- **The co–annihilation region:** The \( \tilde{\ell}_R \) (or \( \tilde{\tau}_1 \)) exchange contribution to \( \sigma \chi^0_1 \) is maximal if
\( m_{\tilde{\ell}_R} \) (or \( m_{\tilde{\tau}_1} \)) is essentially equal to \( m_{\chi^0_1} \). In such a situation \( \Omega \chi^0_1 h^2 \) can no longer be
estimated from \( \chi^0_1 \) pair annihilation alone [34]. If another superparticle \( \tilde{P} \) is close in
mass to \( \chi^0_1 \), reactions of the kind \( \chi^0_1 + f \leftrightarrow P + f' \), with rate \( \propto \exp (–m_{\tilde{P}}/T) \), will stay
in equilibrium much longer than reactions of the kind \( \chi^0_1 + \chi^0_1 \leftrightarrow f + \bar{f} \), which have
rate \( \propto \exp (–2m_{\chi^0_1}/T) \); here \( f, f' \) are some SM particles with mass \( \lesssim T \). In this case
the total number of superparticles, which determines the final LSP relic density, can
also be reduced by \( \chi^0_1 \tilde{P} \) co–annihilation or \( \tilde{P} \bar{\tilde{P}} \) annihilation. Note that the \( \tilde{P} \) number
density is suppressed by the Boltzmann factor \( \exp [–(m_{\tilde{P}} – m_{\chi^0_1})/T] \) relative to the \( \chi^0_1 \)
number density. Since the relevant temperature is \( T \approx T_F \lesssim m_{\chi^0_1}/20 \), this Boltzmann
suppression implies that co–annihilation can only be important for scaled mass difference
\( (m_{\tilde{P}} – m_{\chi^0_1})/m_{\chi^0_1} \lesssim 0.2 \) or so. On the other hand, the total co–annihilation cross section
can be significantly larger than the \( \chi^0_1 \) pair annihilation cross section. In case of co–
annihilation with \( \tilde{\ell}_R \) (or \( \tilde{\tau}_1 \)), the enhancement is primarily due to the fact that co–
annihilation, as well as annihilation of slepton pairs, can occur from \( S \)-wave initial states;
it can thus reduce \( \Omega \chi^0_1 h^2 \) by about one order of magnitude, for vanishing mass
difference [33]. On the other hand, co–annihilation with a light \( \tilde{t}_1 \) squark, which has
strong QCD and Yukawa interactions, can reduce the relic density by more than three
orders of magnitude [36]. Co–annihilation with sleptons can therefore “only” increase
the upper bound on slepton masses by about a factor 3.5; this range may still be covered
by future \( e^+e^- \) colliders. On the other hand, co–annihilation with \( \tilde{t}_1 \) can bring a bino–
like LSP as heavy as 3 TeV into agreement with the constraint (4), i.e. it allows to
construct cosmologically allowed SUSY spectra well beyond the reach of any currently

\(^1\)This statement is strictly true only for a CP–invariant theory. However, bounds on the electric dipole
moments of the neutron and electron tightly constrain the CP–violating phases relevant to \( \chi^0_1 \) annihilation.
planned collider; recall that in models with gaugino mass unification a bino mass of 3 TeV implies a gluino mass near 20 TeV!

- **The Higgs pole region:** The $\tilde{\chi}_1^0\tilde{\chi}_1^0$ annihilation cross section can be greatly enhanced if some $s$–channel diagram becomes resonant. The potentially most important such diagrams involve the exchange of the neutral CP–odd Higgs boson $A$, since it can couple to an $S$–wave initial state. If $2m_{\tilde{\chi}_1^0} \sim m_A$ the annihilation cross section scales like $1/\Gamma_A^2$ rather than $1/m_{\tilde{\chi}_1^0}^2$, which introduces an enhancement factor $1/\alpha^2 \sim 10^4$. This over–compensates the small coupling (of order $g_Y \sin \theta_W M_Z/|\mu|$) of bino–like neutralinos to Higgs bosons, and thus again allows to construct cosmologically acceptable scenarios with $m_{\tilde{\chi}_1^0} > 1$ TeV.

This discussion shows that $\Omega_{\tilde{\chi}_1^0} h^2$ depends on many details of the SUSY spectrum, i.e. on many soft breaking parameters. It is thus usually analyzed in models that describe the entire spectrum in terms of a small number of free parameters, the most common one being the minimal SUGRA model (sometimes called cMSSM). Here one postulates a common scalar mass $m_0$, a common gaugino mass $m_{1/2}$, and a common trilinear scalar interaction parameter $A_0$, at the high “input” scale $M_{\text{GUT}} \simeq 2 \cdot 10^{16}$ GeV. The ratio of vevs $\tan \beta$ is also a free parameter, as is the sign of $\mu$, while $|\mu|$ is fixed by the requirement $M_Z = 91.2$ GeV. In this scenario the bulk region is limited by $m_0 \lesssim 250$ GeV, $m_{1/2} \lesssim 350$ GeV, with little dependence on $A_0$ and $\tan \beta$. $\tilde{\chi}_1^0 - \tilde{t}_1$ co–annihilation is difficult to achieve in this model \cite{32, 37}, but co–annihilation with sleptons does occur over a narrow strip of gaugino masses satisfying $m_{1/2} \simeq 2.5 m_0$, the exact location depending on $A_0$ and $\tan \beta$. In this model the Higgs pole region only exists \cite{38} at large $\tan \beta \gtrsim 50$. This model allows a fourth region where the constraint \cite{1} can be satisfied:

- **The focus point region:** If $m_0^2 \gg m_{1/2}^2$, the soft breaking contribution to the mass of the Higgs boson that couples to the top quark is very small at the weak scale, independent of its value at the input scale (this is the “focusing”) \cite{39}. In this case $|\mu|$, the supersymmetric contribution to the Higgs boson masses, also comes out to be small, so that $\tilde{\chi}_1^0$ has significant, perhaps even dominant, higgsino components. This implies sizable annihilation cross sections into final states containing electroweak gauge and/or Higgs bosons. In the extreme case of higgsino–like LSP this again allows $m_{\tilde{\chi}_1^0} > 1$ TeV \cite{10}; note that annihilation now occurs from an $S$–wave initial state, via $SU(2)$ [rather than $U(1)_Y$] couplings.

The situation is illustrated in Fig. 4, which shows (in green) the region in the $(m_0, m_{1/2})$ plane satisfying $\Omega_{\tilde{\chi}_1^0} h^2 \leq 0.5$ for $A_0 = 0$, $\mu > 0$ and $\tan \beta = 10$ (left) and 60 (right). In the left panel, the “bulk” region is near the origin (and thus severely constrained by collider experiments), while the “co–annihilation” and “focus point” regions are the nearly horizontal and nearly vertical strip, respectively. These regions also exist at $\tan \beta = 60$ (right), but are supplemented by the much larger “Higgs pole” region. This last region is rather wide, since the thermal motion allows resonant $A$–exchange even if $m_{\tilde{\chi}_1^0}$ is well below $m_A/2$. These plots, which are updates of ref.\cite{32}, vividly illustrate that even in this very constrained model the condition \cite{2} does not allow to derive useful upper bounds on sparticle masses, although it

\footnote{Unless $\tan \beta$ is very large, in which case $L - R$ mixing in the $\tilde{\tau}$ sector reduces $\Omega_{\tilde{\chi}_1^0} h^2$ significantly; this mixing depends on both $\tan \beta$ and $A_0$ (in the latter case, mostly via $|\mu|$).}
does exclude large regions of parameter space. These figures also show experimentally excluded regions (grey), the region where $113 \text{ GeV} \leq m_h \leq 117 \text{ GeV}$ in accordance with the $2\sigma$ LEP evidence for an SM–like neutral Higgs boson (red), and the region compatible with $g_\mu - 2$ at the $2\sigma$ level (blue). These laboratory “discoveries”, if taken seriously, do impose upper limits on sparticle masses; in particular, a TeV–scale $e^+e^-$ collider would have to discover $\tilde{\chi}_1^0\tilde{\chi}_2^0$ production in the overlap region (black) where all constraints are satisfied.

Fig. 4: Regions in mSUGRA parameter space allowed by the DM constraint (green), by the LEP evidence for an SM–like Higgs boson (red), and by the Brookhaven measurement of $g_\mu - 2$ (blue); in the black regions, all three constraints are satisfied simultaneously, while the grey regions are excluded experimentally. These figures are for $A_0 = 0$, $\mu > 0$ and $\tan \beta = 10$ (left) and 60 (right).

The Dark Matter argument by itself is therefore not sufficient to tell us in what region of mSUGRA parameter space we are likely to be. The situation is even worse in other widely discussed, constrained SUSY models. In models with gauge–mediated SUSY breaking (GMSB), the LSP is the gravitino $\tilde{G}$, whose mass is essentially independent of the other sparticle masses. In anomaly–mediated SUSY breaking (AMSB) the LSP is wino–like, and makes a good thermal DM candidate for $m_{\tilde{\chi}_1^0} \simeq 1.5 \text{ TeV}$. This prediction is falsifiable in principle, but not at the any of the colliders that are now under serious consideration.

On the other hand, collider physics experiments can make crucial contributions to DM physics if $\tilde{\chi}_1^0$ indeed constitutes most DM. Once a DM signal has been established, the first order of business would be to check whether the WIMP mass is compatible with $m_{\tilde{\chi}_1^0}$ as measured at colliders. This crucial check is probably much more demanding for DM search experiments than for collider experiments. Even at the LHC it is fairly easy to measure $m_{\tilde{\chi}_1^0}$ to $\sim 10\%$ [12]. In contrast, the DAMA “signal” [43] only determined the WIMP mass within a factor of two; the fact that DAMA almost certainly did not detect WIMPs – their result is now contradicted by two other experiments [44], one of which is background–free – does not change the conclusion that determining the WIMP mass from the observed nuclear recoil spectrum is quite difficult. Fortunately, “indirect” WIMP signals, e.g. high–energy neutrinos from WIMP annihilation in the Sun or Earth, or a $\gamma$–ray line from WIMP annihilation in the galactic halo [29] – should more easily allow to determine the WIMP mass.

§Thermal relic gravitinos can form the Dark Matter if their mass is about 0.5 keV [41], which implies that the NLSP is stable at the time scale of accelerator experiments. Observation of $\tilde{\chi}_1^0 \to \tilde{G}\gamma$ or $\tilde{\tau}_1 \to \tilde{G}\tau$ decays could thus exclude $\tilde{G}$ as DM candidate.
Input from collider experiments is also needed for the calculation of the thermal $\chi_1^0$ relic density. In the “bulk region” the required data – essentially $m_{\chi_1^0}$ and $m_{\tilde{t}_R}$ – can already be provided by the LHC. LHC data can even check whether $\chi_1^0$ has significant higgsino components, which would drastically change the di-lepton invariant mass distribution, e.g. from $\chi_1^0 \rightarrow \ell^+\ell^-\chi_0^0$ decays \[13\]. In the co–annihilation region one would need an accurate measurement of the $\tilde{\tau}_1 - \chi_1^0$ mass difference, which is impossible at the LHC and might even be difficult at future $e^+e^-$ colliders \[10\]. In the focus point region one would essentially only need to know the parameters of the neutralino–chargino system, which are easily measured at $e^+e^-$ colliders operating at $\sqrt{s} > 2m_{\chi_1^0}$ \[17\]; however, as mentioned above, this may require a multi–TeV collider if $\chi_1^0$ is higgsino–like. Finally, in the Higgs pole region one would need accurate measurements of the parameters of both the neutralino and Higgs sectors of the theory. The latter would be difficult at any planned collider if $m_A$ exceeds 1 TeV or so.

Suppose measurements at colliders indeed allow to predict the thermal $\chi_1^0$ relic density; what would we learn? If the value comes out higher than the upper bound in \[8\], we know that the assumptions going into this calculation are wrong, i.e. either $\chi_1^0$ is unstable, or it didn’t reach chemical equilibrium with the SM plasma after the period of last entropy production. Actually, the latter is not sufficient to avoid “overclosing the Universe”, since non–thermal mechanisms can easily over–produce LSPs. This is true in particular for direct LSP production in the decay of massive particles whose decays are responsible for the last entropy production, e.g. inflaton decays \[18\]. A too high thermal $\chi_1^0$ relic density would thus be evidence (though not proof) that $\chi_1^0$ is not stable at cosmological time scales. This could mean that $R$–parity is not conserved, or that $\chi_1^0$ is actually not the LSP, which instead resides in the “hidden sector” of the theory; this sector interacts with ordinary (s)particles only through gravitational–strength interactions. In order to clarify this issue one may then ultimately have to build an experiment searching for $\chi_1^0$ decay with lifetime up to one minute or so; longer $\chi_1^0$ lifetimes are essentially excluded by combinations of constraints from Big Bang nucleosynthesis and the spectrum of the cosmic microwave background.

Conversely, if the thermal $\chi_1^0$ relic density comes out too low, but WIMP search experiments indicate that $\chi_1^0$ contributes to the DM, one has proven the existence of non–thermal $\chi_1^0$ production \[19, 18\], typically from the decay of some more massive object. Finally, if the thermal $\chi_1^0$ relic density indeed comes out correct, and WIMP search experiments confirm that $\chi_1^0$ forms the DM, we’d have evidence that the Universe reached a re–heat temperature $T_R \gtrsim 0.1m_{\chi_1^0}$; note that currently we can only infer \[50\] $T_R \gtrsim 1 \text{ MeV}$, from the quantitative success of Big Bang nucleosynthesis.

In addition to the thermal relic density, one will also want to predict the cross section for $\chi_1^0$–nucleon scattering as well as $\sigma(\chi_1^0\chi_1^0 \rightarrow \gamma\gamma)$. The former determines the size of both the direct WIMP signal and the signal for WIMP annihilation in the Earth or Sun (for given WIMP flux), while the latter is required for the calculation of the size of the most obvious signal for WIMP annihilation in the halo of the galaxy. If a WIMP signal has been detected, with kinematics indicating that $\chi_1^0$ is indeed that WIMP\[1], knowledge of these cross section is required to glean information about the WIMP distribution in the galaxy from the observed signal(s), which would be of great interest to galactic modellers; indeed, it might open the door to some sort of “WIMP astronomy”. In the absence of a positive signal these cross sections would indicate the sensitivity an experiment would need to reach in order to exclude $\chi_1^0$ as

\footnote{WIMP discovery by itself should not be confused with discovery of SUSY. Other WIMP candidates already exist \[51\], and even more models are certain to be constructed as soon as a signal is found.}
DM. Unfortunately the $\tilde{\chi}_0^0$–nucleon scattering cross section is again sensitive \cite{29} to the Higgs sector, in particular to the mass and couplings of the heavy neutral CP–even Higgs boson, which are likely to be among the SUSY parameters that are most difficult to measure.

As mentioned earlier, most WIMP detection methods can determine the WIMP mass only with rather large uncertainty. Similarly, the extracted cross section will come with large uncertainty due to the rather poorly known $\tilde{\chi}_1^0$ flux. As discussed above, the flow of information regarding these quantities will ultimately be from collider physics to astro–particle physics, even if the first discovery comes from WIMP searches. Nevertheless, proof that $\tilde{\chi}_0^0$ contributes to the DM would provide invaluable information to “mainstream” particle physics. It would imply upper limits on the size of all $R$–parity violating couplings that are many orders of magnitude below the sensitivity of conceivable laboratory experiments (including proton decay). Moreover, it would imply that $\tilde{\chi}_0^0$ is indeed the LSP, i.e. would impose a strict lower limit on the mass of any superparticles in the hidden sector, in particular on $m_\tilde{G}$; this information would be of great interest for both model building and cosmology.

5) The Most Energetic Cosmic Rays

The most energetic particle collisions that are amenable to experimental analysis by humans do not occur at colliders; rather, their origin is the collision of cosmic ray (CR) “primaries” with nuclei in the atmosphere of Earth. The spectrum of these primaries has now been measured over more than ten orders of magnitude in energy, and about 30 orders of magnitude in flux. Here we are concerned with the high–energy end of this spectrum, events with energy around or above $10^{11}$ GeV. The very existence of such events is puzzling, for at least two reasons \cite{52}:

- It is believed that most CR primaries are accelerated in shock fronts surrounding celestial bodies. The basic requirement for such acceleration (or “bottom–up”) mechanisms to work is that there must be a sufficiently strong magnetic field $B$ extending over a sufficient length $L$, i.e. $B \cdot L$ must be sufficiently large. For example, in a field of $\sim 8$ T extending over $\sim 30$ km, one cannot accelerate protons to more than about 14 TeV; this describes the LHC, of course. Hillas has shown some time ago \cite{53} that few, if any, known objects in the Universe have sufficient $B \cdot L$ to accelerate protons to $10^{11}$ GeV even for perfect efficiency (i.e. if the $B$–field has collider quality, which does not strike me as a likely proposition). In such a “bottom–up” scenario neutral primaries would have to be produced when charged particles collide with ambient matter; these would require even higher energies for these charged particles than what is observed on Earth.

- At $E \gtrsim 5 \cdot 10^{10}$ GeV, protons can photo–produce pions on the cosmic microwave background, $p + \gamma_{\text{CMB}} \rightarrow \pi + N$ ($N = n, p$). A collision of this kind will reduce the proton’s energy by typically 20% or so. At yet higher energies, multi–pion production processes, which lead to even higher energy losses, also become important.\footnote{Protons can also lose energy through $p + \gamma_{\text{CMB}} \rightarrow p + e^+e^-$. This reaction has much lower energy threshold, but it reduces the proton’s energy only by one percent or so, and is thus of lesser importance.} The interaction length for such collisions is “only” a few tens of Mpc. As a result, protons with energy above this “GZK cut–off” \cite{54} must have been produced within 50 or 100 Mpc of Earth. The same is true for photons, which can be absorbed by $e^+e^-$ pair production on background photons in the radio band, as well as heavier nuclei, which can photo–dissociate on the microwave background, all with interaction lengths of a few to a few tens of Mpc.
the other hand, extra–galactic magnetic fields, as well as the fields in our own galaxy, are so weak that even protons at $10^{11}$ GeV should still point back to their source (within a few degrees), if it is within a couple of GZK interaction lengths. However, there are no nearby potential “bottom–up” sources of $10^{11}$ GeV particles near the measured arrival directions.

The solution to these puzzles most likely requires new astrophysics, new particle physics, or both. There have been several proposals of how to avoid the GZK cut–off. For example, the primaries might be $[53]$ “$R$–hadrons”, containing a light gluino (with mass of a few GeV). This would push the GZK cut–off to a few times $10^{11}$ GeV, just beyond the energy range covered by current observations. However, gluinos of the required mass have almost certainly been excluded by collider experiments $[30, 56]$. Moreover, this “solution” does not address the problem of how such energies can be reached in the first place. Another proposal postulates $[57]$ neutrinos as primaries, with mb cross sections on air. However, models of this kind are severely limited by unitarity arguments $[58]$. Besides, the neutrinos would presumably have to be produced in the collision of even more energetic protons, thereby aggravating the first of our two problems. The same is true for the so–called “$Z$–burst” scenario $[59]$, where one postulates a very large flux of neutrinos with energies extending to $10^{12}$ GeV or more impinging on neutrinos in the halo of our galaxy. These collisions produce on–shell $Z$–bosons, whose decay into (ultimately) protons and photons (among other things) produces the CR “primaries”. It has recently been argued $[60, 61]$ that the required neutrino flux violates current experimental upper limits (see however $[62]$). The perhaps most radical proposal $[63]$ is to simply remove the GZK cut–off by assuming that Lorentz invariance does not hold at sufficiently high energies. In this case the photoproduction cross section at $E_{\gamma} \sim 10^{-4}$ eV, $E_{p} \sim 10^{11}$ GeV may not be related to the measured cross section for $E_{\gamma} \sim$ GeV on protons at rest. However, this again does not address the first of our problems.

The to my mind best motivated “new physics” explanation is based on the hypothesis that the highest energy CR events come from the decay of some very massive $X$–particle. We clearly need $M_{X} \geq 10^{12}$ GeV, since the most energetic event ever observed $[64]$ has $E \approx 3 \cdot 10^{11}$ GeV, and only a (small) fraction of $M_{X}$ will go into the energy of any one $X$ decay product. On the other hand, values of $M_{X}$ as high as $10^{16}$ GeV are possible. In fact, the first models of this kind used GUT–scale particles bound in topological defects $[65]$; this could ensure their longevity. The $X$–particles might also roam freely, and be almost stable by virtue of having extremely suppressed couplings to normal matter $[66]$. Several mechanisms have been suggested that allow to produce sufficiently many such particles at the end of inflation $[67, 48]$. This solves the first of our two problems: basically $M_{X}$ stores energy from an early, extremely violent epoch in the history of the Universe. The second problem is solved since no association with known sources of high energy particles is needed here. Topological defects may not be associated with galaxies. In contrast, most “freely” moving $X$–particles should be enriched in galaxies $[68]$, just as other forms of matter are; in fact, they might even form the Dark Matter. One therefore expects most relevant decays of freely moving $X$–particles to occur in the halo of our own galaxy, which predicts an approximately (but not exactly $[69]$) isotropic distribution of arrival directions, in at least qualitative agreement with data.
An $N$–body primary $X$–decay would be like an $N$–jet event produced at an $e^+e^-$ collider operating at $\sqrt{s} = M_X$. As illustrated in Fig. 5, such jets would look quite different from the QCD jets we know. To begin with, the hierarchy between $M_X$ and the weak scale cries out for supersymmetry as a stabilization mechanism. Since $M_X \gg$ the sparticle mass scale $m_{\text{SUSY}}$, at the early stages of the jet evolution superparticles can be “showered off” just like any other partons. Moreover, we know that all gauge interactions will have similar strength at energy or momentum scales near $M_X$. Therefore electroweak interactions, and perhaps third generation Yukawa interactions, can play significant roles in the early stages of jet evolution. However, eventually (after $10^{-27}$ seconds or so) the shower (virtuality) scale will drop below $m_{\text{SUSY}}$. At this point superparticles will decouple from the shower and decay, possibly through lengthy decay cascades well known from LHC studies. Slightly later the same will happen to top quarks as well as electroweak gauge and Higgs bosons. From this point on electroweak interactions are indeed far weaker than QCD ones, so that the shower evolution is dominated by strong interactions. Later yet, all partons will hadronize. This is the end of jet evolution at collider experiments, but here we also have to treat the decays of unstable hadrons (mesons, neutrons etc.) as well as $\tau$–leptons and muons. The final output of such a calculation \cite{BroOK} is the spectra of stable particles “at source”, i.e. at the location of $X$–decay: electrons, protons, three kinds of neutrinos (including antiparticles in all cases), photons, and LSPs.
comparing with data, propagation effects may have to be included, especially when one is studying extra-galactic sources.

One finds that the photon flux at source exceeds the proton flux; this is in fact also true for jets produced at current colliders. The ratio is about 1.5 : 1 at large $x = E_{\text{particle}}/E_{\text{jet}}$ if the jet originates from a (s)quark, about 5 : 1 at small $x$ independent of the particle starting the jet, and exceeds 10 : 1 at large $x$ if the primary particle is a (s)lepton or Higgs boson. This is problematic, since data indicate that few, if any, of the highest energy events are due to photons. One thus has to get rid of most of the photons on the way from the source to Earth. This is actually expected to occur for sources at cosmological distances, but would require the galactic radio background to be much stronger than current best estimates if the sources reside in the halo of our own galaxy.

Another practical problem at present is the large discrepancy between the two measurements with the best statistics of the CR spectrum at highest energies, by the AGASA and HiRes collaborations. AGASA sees a spectrum with similar, or perhaps even larger, spectral index as at lower energy, where the spectrum can be described by a power law with power $\simeq -2.7$. In contrast, the HiRes spectrum shows a break at the GZK energy $\simeq 5 \cdot 10^{10}$ GeV, and drops quickly for $E > 10^{11}$ GeV.† The HiRes spectrum can be fitted assuming a homogeneous distribution of sources with $E^{-2}$ spectrum. However, the fact that such a fit is possible doesn’t tell us much about the actual nature of these sources; the GZK effect ensures that the information about the spectrum at source with $E > 10^{11}$ GeV is essentially lost if the sources are distributed homogeneously. In fact, the HiRes spectrum does not even prove that sources are distributed uniformly, since a top–down model with $M_X \simeq 10^{12}$ GeV will also give a rapidly falling spectrum above $10^{11}$ GeV. Both the HiRes and AGASA spectra can be fitted within such models. Performing such fits is necessary to fix the overall normalization of the fluxes of very energetic particles predicted by these models. However, this normalization is uncertain not only because of the discrepancy between the (supposedly) best data sets currently available, but also because the “background” is not well understood. Most people in the field assume that the spectrum below $10^{10}$ GeV is dominated by conventional (bottom–up) sources, while top–down models should at least explain the spectrum above $10^{11}$ GeV. However, since we do not know what these “conventional sources” actually are, let alone their spectral characteristics, it is not clear how much of the spectrum between $10^{10}$ and $10^{11}$ GeV should be explained through $X$ decays. This is of some concern, since in that energy range the spectrum is known far better experimentally than at even higher energy.

Nevertheless we can already make semi–quantitative predictions that should allow to test these models in the not too distant future. In particular, one expects the neutrino flux to be even higher than the photon flux, basically because jets contain more charged pions (each of which produces three neutrinos) than neutral ones (which decay into two photons each). Neutrino signals at very high energies have therefore early been recognized as potential test of this kind of model. More recent calculations indicate healthy event rates for next–generation experiments like Ice Cube, with detection possible in the near future in many scenarios. However, “bottom–up” models can also lead to significant fluxes of very energetic neutrinos, either through the interaction of accelerated protons with material near the source, or through the GZK effect. The observation of such a neutrino signal may therefore not be sufficient to discriminate between these classes of models.

†For some reason the most recent HiRes spectrum no longer contains the highest energy event ever observed by the Fly’s Eye collaboration, the direct predecessor of the HiRes experiment.
On the other hand, only “top–down” models predict a significant flux of LSPs with ultra–high energy. Even though the cm energy of the collision of a $10^{11}$ GeV proton with a nucleus (or proton) at rest is much larger than $m_{\text{SUSY}}$, only a tiny fraction of such collisions will produce superparticles, while each collision will produce (many) pions which decay into neutrinos. Existing bounds on the flux of very energetic neutrinos thus imply an undetectably small LSP flux in such models. Even in top–down models one needs an effective target mass of at least $10^{12}$ tons or so to detect these very energetic LSPs. The currently best hope for such experiments are space–based photo–detectors looking for fluorescence light (from the de–excitation of nitrogen molecules excited in the evolution of showers in air) in Earth’s atmosphere. In order to suppress the much larger background from neutrinos, one has to look for *upgoing* events. The trick is that (at least for bino–like LSP) there is an energy range (around $10^{9}$ GeV) where the Earth is opaque to neutrinos, but still allows a sizable fraction of the LSP flux to pass. This should (at least in principle) allow to detect the LSP flux predicted in many top–down models [76].

If it could be proven that the most energetic cosmic rays indeed originate from the decay of very energetic particles, we would have gained *experimental* access to energies far beyond the wildest dreams of accelerator physicists. This prospect is thrilling indeed, but we have to keep in mind that this “access” is bound to be rather indirect. Out of several tens of thousands of particles (with $x \geq 10^{-7}$) produced in any given $X$ decay, we at best observe a single one. In other words, the only measurables are the single–particle inclusive spectra of protons, neutrinos (alas, with all flavor information most likely washed out by oscillations), LSPs and, perhaps, photons (electrons certainly loose their energy through synchrotron radiation before reaching Earth). These spectra depend on various aspects of physics at energy scales all the way up to $M_{X}$. In particular, we’ll want to know the primary $X$ decay products (in order to “profile” the $X$–particle). Another important question is what kind of interactions are present at energies just below $M_{X}$, e.g. whether there are additional gauge interactions beyond those present in the SM. However, before we can address such questions we need to have a detailed understanding of physics at scale $m_{\text{SUSY}}$. In particular, sparticle cascade decays will affect the spectra of neutrinos, LSPs and (to a lesser extent) protons and photons. In this exciting scenario we would thus need input from experiments at TeV–scale colliders in order to learn about physics at energy all the way up to $M_{X} \gtrsim 10^{12}$ GeV.

6) Summary and Conclusions

In this talk I attempted to describe some of the connections between astro–particle and collider physics. The observation of neutrino oscillations gives additional impetus to searches for lepton flavor violation at accelerators, which in turn can help to pin down supersymmetric models of leptogenesis. Electroweak baryogenesis in the MSSM requires a light $\tilde{t}_{1}$ squark and sizable CP–violation in the chargino and neutralino sector, and can thus easily be tested at a high–energy $e^{+}e^{-}$ collider (if not sooner). Information from collider experiments will be essential to predict the thermal $\chi_{1}^{0}$ relic density, as well as the cross sections that are needed to predict the size of WIMP signals; conversely, proof that $\chi_{1}^{0}$ is sufficiently long–lived to form the Dark Matter in the Universe will imply constraints on SUSY parameters that cannot be obtained at colliders. Finally, if it can be established that the highest energy cosmic rays originate from the decay of very massive particles, detailed information about sparticle masses and decay patterns will be required to deduce the physics at very high energy scales from the spectra of very energetic
protons, neutrinos and LSPs that should eventually be measured by astro–particle physics experiments. These kinds of connections should prove fruitful for both fields.

Nevertheless a word of caution might be in order. At present collider physicists should not put too much stock in “cosmologically favored” scenarios or regions of parameter space. All the phenomena discussed here (neutrino masses, the baryon asymmetry of the Universe, Dark Matter, and cosmic rays with $E > 10^{11}$ GeV) currently have several competing explanations, with usually quite different repercussions for collider physics (or none at all). When planning and performing experiments at colliders, it is therefore probably better to simply ignore cosmological considerations, although it might prove very rewarding to keep these considerations in mind when analyzing the data. Similarly, astro–particle physicists should be aware that many limits from collider physics hold only under specific assumptions. To name an important example, there is no model–independent limit on WIMP masses from colliders. However, the interrelations between astro–particle and collider physics should become very important indeed once signals replace bounds, i.e. once an unambiguous discovery of “new physics” has been made in either area. Neutrino physics has shown the way; I fervently hope that additional signals of equal or even greater importance will be discovered sooner rather than later.

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