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ON DARK MATTER DETECTION AND SOLAR NEUTRINOS

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Contents

Introduction

Particle Physics has recently provided possible explanations of several astrophysical problems such as the dark matter of the universe and the longstanding solar neutrino puzzle.

The dark matter problem is the well established evidence that indicates that at least 90 % of the mass in the universe is not luminous, that is, that it neither emits nor absorbs electromagnetic radiation of any frequency and is observed only by its gravitational effects. The determination of the nature of this "dark matter" (hereafter DM) is one of the important open questions confronting modern physics nowadays.

Many DM candidates have been proposed. It is possible that ordinary baryonic matter is present in the form of "Jupiters" (planet-size gas balls too small to shine) or black holes. However, the agreement between standard primordial nucleosynthesis and the observed abundances of the light nuclides imposes an upper limit on the total baryonic density of the universe that does not allow baryons alone to account for all the non luminous matter. So, non baryonic DM is necessary, and the most accepted candidates for it are weakly coupled elementary particles. In fact, massive neutrinos or particles as the axion (proposed to solve the strong CP problem), the monopoles or the lightest supersymmetric particle can contribute significantly to the overall energy density of the universe. Other DM candidates have recently been proposed in relation to the solar neutrino problem.

The solar neutrino problem corresponds to the fact that only one third of the predicted neutrinos produced in the ^{8}B reactions in the solar core have been observed. Two types of solutions were proposed to explain this. The first involves changes in the properties of the neutrinos: a small neutrino mass ($\sim 10^{-6}$ eV) may produce oscilations between the three neutrino flavours so that the detected ν_e are just 1/3 of the initial ones, a large magnetic moment ($\sim 10^{-10}\mu_B$) may allow an helicity flip in the high magnetic fields in the sun, etc.. The second kind of solution require changing the solar properties in order that the central temperatue of the sun, to which the production of ^{8}B neutrinos is very sensitive, results slightly lower than the standard solar model estimate. Faulkner, Gilliland, Press and Spergel realized that this could be achieved if a thermalization of the central core with the neighbouring (colder) regions arises due to the presence of a large population of weakly interacting particles, since this would efficiently increase the solar heat transport, without affecting the other observed solar properties. These hipothetic particles have been named "cosmions". They should be trapped by the sun from the halo of our galaxy (a spheroidal region composed of non luminous matter extending up to several optical radii). To get the large concentration required, cosmions should constitute an important fraction of the halo density, so that they also naturally become a DM candidate.

To be trapped and transport heat efficiently, the cosmion properties (mass and cross sections) should lie in a narrow range. None of the conventional dark matter candidates (heavy neutrinos, sneutrinos, photinos, higgsinos,...) satisfy these constraints, but some proposals for cosmions were done, involving new interactions and extensions of the standard model. Clearly, it is important to see how these conjectured particles can be tested or, moreover, if they are not just ruled out by existing experiments. This is one of the subjects addressed to in this thesis, where it is shown that hadron colliders working in the TeV region will be able to probe some cosmions models, since a cosmion produced in a $p\bar{p}$ collision escapes detection and can give rise to signals with large missing transverse momentum. It is also shown that the indirect detection through the observation at underground proton-decay experiments of the neutrino flux arising from the annihilation of the cosmions trapped in the sun, places stringent bounds on several candidates.

The other subject faced in this thesis is related to the recent recognition by Griest and Barbieri, Frigeni and Giudice that in the context of the minimal supergravity theory, the general neutralino (a mixture of photino, zino and higgsinos) can have coherent interactions with nuclei, contrary to the previous belief that only spin dependent interactions affect Majorana particles. If the lightest neutralino is an important component of the DM in the galaxy, this enhanced cross section with heavy nuclei increases significantly the capture of neutralinos by the sun or the earth, specially when the Higgs mediating the coherent interaction is light ($m_H < 30$ GeV). In this condition, the large population of captured particles can, as the cosmions, be constrained by the neutrino flux they would produce when they annihilate. Finally, it is shown that in some regions of the supersymmetric parameter space, and if the scalar Higgs is very light ($m_H \simeq$ 2 GeV) the neutralino can not only provide a good DM candidate, but also satisfy the constraints necessary to solve the solar neutrino problem, this is, it becomes a cosmion.

Chapter I gives a brief introduction to the topics that were mentioned here and that are employed in the succesive chapters, where the contributions of this thesis are presented.

CHAPTER 1

Dark matter and solar neutrinos

1.1-The Dark Matter

The matter content of the universe is usually measured in terms of $\Omega = \rho/\rho_c$, which is the cosmological density ρ in units of the critical density $\rho_c = \frac{3H_o^2}{8\pi G}$ necessary to close the universe. $H_o \equiv h \cdot 100\ Km/s\ Mpc$ is the Hubble constant, with $h \simeq 0.5 - 1$ being the Hubble parameter (1 parsec=3.26 light years).

Although the visible mass accounts for $\Omega_{vis} \simeq 0.01$, several observations indicate that $\Omega_{tot} \geq 0.1 - 0.2$ [1.1]. One of the most striking examples is given by the measurements of the orbital velocities of stars or gas clouds orbiting around spiral galaxies. In fact, it is observed that these velocities remain approximately constant outside the regions where the light falls exponentially off and at least up to 2-3 optical radii [1.2]. Since $v(r)^2 = \frac{GM(r)}{r}$, the flatness of the rotational curves means that the mass contained within a radius r grows linearly with r. This is an indication of the existence of an halo of non luminous matter and it requires the presence of $M_{tot} > 3 - 10\ M_{vis}$ in order to provide the gravitational attraction necessary to bind objects of such high rotational velocities. The existence of DM is also indicated by observations at larger scales, for example $\Omega \simeq 0.2$ is associated with halos of groups and clusters of galaxies. The amount of DM on even larger scales is uncertain, in part because it is not known how this DM is distributed. If its distribution is the same as that of the galaxies, then $\Omega_{DM} \sim 0.2$[1.3]. But if a smooth component is assumed, it is possible that $\Omega_{DM} \simeq 1$. Tentative observations from the Infrared Astronomical Satellite, which is the largest scale redshift survey that has been used to estimate Ω, suggest that $\Omega \geq 0.7$ [1.4]. Finally, there is also a theoretical prejudice (inflation theory)[1.5] that favours $\Omega = 1$.

Which is the nature of this DM? Standard primordial nucleosynthesis calculations fit well with the observed abundances of the light nuclides (H, D, ^{3}He, ^{4}He, ^{6}Li, ^{7}Li) only if the density Ω_b of ordinary ("baryonic") matter lies in the range $0.014\ h^{-2} \leq \Omega_b \leq 0.025\ h^{-2}$, with a firm upper limit $\Omega_b \leq 0.14$ [1.6]. If $\Omega \simeq 1$, the majority of the DM must be non baryonic, but if the cosmological density is actually at the lower end of the observationally allowed range, $\Omega \simeq 0.2$, it could be that the DM is mainly baryonic. However, in this case many constraints restrict the possible forms it may take to "Jupiters" or black holes, and it is difficult to invent schemes in which 90 % of the baryonic matter in galaxies is converted to these unusual forms rather than to stars. Though, it is quite plausible that the majority of the DM is nonbaryonic.

Many nonbaryonic elementary particles can provide the DM [1.7]. Among the most accepted candidates are massive neutrinos, axions and the lightest supersymmetric particle. In an early stage of the evolution of the universe, these particles are in thermal and chemical equilibrium with the remaining ones, but as the universe expands and cools, the "freeze-out" temperature is reached when the annihilation rate becomes too low to allow the equilibrium condition to be preserved. If the particles are stable, after this decoupling their density diminishes only by their posterior (not very efficient) annihilation and because of the expansion of the universe. This leads, in general, to a significant contribution to Ω. In this thesis, two types of DM are considered: the neutralinos, one of the most likely candidates for the lightest supersymmetric particle, and the cosmions, proposed recently in relation with the solar neutrino problem.

1.2-The neutralino

Many theoretical problems of the standard Glashow Weimberg Salam model, as the description of fundamental scalars free from the naturalness problem [1.8] or the inclusion of gravity, can be solved in the context of a locally supersymmetric theory. The minimal supergravity model is the simplest locally supersymmetric extension of the standard model. At low energies[1.9], it preserves the gauge group

$SU(3) \times SU(2) \times U(1)$, but each known fermion (boson) has an associated new bosonic (fermionic) superpartner. This doubles the number of degrees of freedom, with the addition that two Higgs doublets are necessary to give mass to all the fermions. In the same way as other symmetries are also present in the standard model (lepton number for instance), in supersymmetric theories there is generally an additional "R-symmetry", that leaves, after the electroweak breaking, a multiplicatively conserved quantum number called R-parity. In terms of the baryon number B, the lepton number L and the particle spin S, the R-parity is

$$R = (-1)^{3(B-L)+2S}. \tag{1.2.1}$$

Since the standard particles are R-even while their superpartners are R-odd, the conservation of this symmetry (that holds when B-L is conserved) has the important consequence of leaving the lightest supersymmetric particle (LSP) stable and making it a possibly important cosmological constituent.

The LSP cannot be a charged or coloured particle, or else it would have been observed in searches of annomalous heavy isotopes [1.10]. The remaining possibilities are the scalar neutrino and the neutralino. The first one has the problem that in general a lighter scalar charged lepton is also predicted, although this pattern might be rearranged by radiative corrections in some models. There are no arguments instead against the neutralino [1.11] and though, it will be considered here as the most likely LSP.

In the minimal standard model, the neutralino is a mass eigenstate of $\tilde{W}_3$, $\tilde{B}$, $\tilde{H}_1^0$, $\tilde{H}_2^0$, corresponding to the superpartners of the neutral SU(2) and U(1) gauge bosons and of the neutral components of the two Higgs. Their mass term has the form

$$-\frac{1}{2}\chi^T M_N \chi, \tag{1.2.2}$$

where $\chi^T = (\tilde{W}_3,\ \tilde{B},\ \tilde{H}_1^0,\ \tilde{H}_2^0)$ is a four component Majorana spinor and

$$M_N = \begin{pmatrix} M' & 0 & -M_Z cos\beta sin\theta_W & M_Z sin\beta sin\theta_W \\ 0 & M & M_Z cos\beta cos\theta_W & -M_Z sin\beta cos\theta_W \\ -M_Z cos\beta sin\theta_W & M_Z cos\beta cos\theta_W & 0 & -\mu \\ M_Z sin\beta sin\theta_W & -M_Z sin\beta cos\theta_W & -\mu & 0 \end{pmatrix} . \tag{1.2.3}$$

This matrix depends on the gaugino masses M' and M, on the vacuum expectation values of the Higgs fields $v_{1,2} = \langle H^0_{1,2} \rangle$ through $tg\beta = v_2/v_1$ and on the mass parameter μ (positive or negative) that couples the two Higgs in the superpotential

$$\mathcal{L} \ni \propto \ \mu H_1 H_2 \ . \tag{1.2.4}$$

The simplifying assumption M=M', corresponding to an underlying unifying group, is generally done and will be employed here. The neutralino states are obtained by a diagonalization of this matrix and some simple and important examples are the photino $\tilde{\gamma} = cos\theta_W \tilde{B} + sen\theta_W \tilde{W}_3$, with $m_{\tilde{\gamma}} = M$, which becomes the lightest eigenstate in the limit $M \to 0$, and the higgsino $\tilde{h} = sin\beta \tilde{H}^0_1 + cos\beta \tilde{H}^0_2$, that is the lightest eigenstate for $\mu = 0$ or $\beta = \pi/4$ and $\mu \ll M, M_Z$, with $m_{\tilde{h}} \propto \mu$.

The mass of the LSP remains an open question, although if supersymmetry is invoked to solve the naturalness problem one expects superpartner masses to be associated to the weak scale. So, it is not unreasonable to postulate that the LSP mass is in the range of a few GeV to tens of GeV. Experimental limits are rather indirect. If produced in accelerator experiments, charged or coloured superpartners could be readily detected, and so lower limits on their masses have been obtained [1.12]. However, although the neutral LSP may be produced, it will not be directly detectable. Its production must be inferred from missing energy experiments, such as "monojets" at $p\bar{p}$ colliders or anomalous single photons events at e^+e^- machines. The requirement that the universe not be overclosed, suggests that a pure photino must have mass $m_{\tilde{\gamma}} \geq 2\ GeV$ [1.11]. However, for the generic neutralino no absolute limits on its mass are known. Some constraints on the parameters entering in the model (μ, β, M) can be obtained from collider experiments [1.13] and, if the neutralino is an important constituent of the DM in the halo, from cosmological requirements and from its direct interaction in germanium detectors [1.14], or by indirect means in underground proton-decay experiments [1.15], as will be shown in chapter 4.

In forthcoming chapters, the capture of DM particles by the sun or the earth will be considered. This trapping results specially important when the particles have coherent (spin independent) interactions with nuclei, particularly in the case of the earth where most elements are heavy. Since massive neutrinos or Majorana particles (as the neutralinos under consideration) couple to the Z^0 gauge boson axially, a spin dependent interaction arises due to Z^0 exchange, yielding to a low energy cross section with nucleus proportional to the total spin of the nucleus. Since most nucleons in nuclei have their spins sum pairwise to zero, this coupling does not increase with the size of the nucleus. However, it has recently been pointed [1.16,1.14] that the exchange of scalar particles (Higgs and in the neutralino case also squarks) give rise to a coherent interaction, that can enhance significantly the scattering from heavy nuclei with respect to the pure Z^0 exchange case.

1.3-Coupling of scalars with nuclei

Many scalar particles (Higgs, axions, etc) have the common feature that they couple to quark masses according to

$$\mathcal{L}_{int} \propto S \sum_q m_q \bar{q} q \ , \qquad (1.3.1)$$

where S represents the bosonic field. Shifman, Vainshtein and Zakharov [1.17] considered the effect of this coupling with nucleons, and a nontrivial result was found in the computation of the matrix element $\langle p | \sum m_q \bar{q} q | p \rangle$ of the interaction (1.3.1) between the nucleonic state $|p\rangle$. In fact, being the valence quarks very light ($m_{u,d} \simeq 10$ MeV), the direct coupling of scalar particles to usual hadrons is manifestly small. However, the effective coupling is not small, since the heavy quark terms in (1.3.1) induce a scalar-gluon effective interaction which dominates at low energy (the energy scale is set by the heavy quark mass).

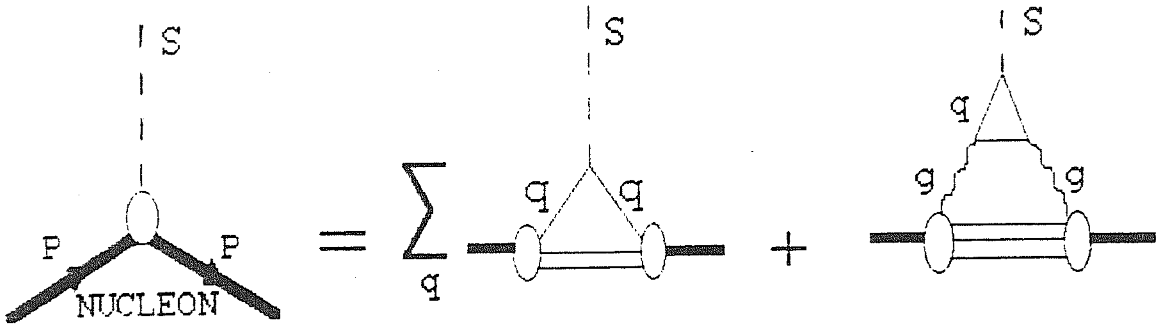

Figure 1

This is pictorially shown in fig.1. The first term in the sum, proportional to the light quark masses, is allways small, but the second one instead is not small when the virtual quarks in the loop are heavy. This piece can be computed in the so called heavy quark expansion [1.18], and the leading term (the only one surviving in the limit $m_q \to \infty$) is completely analogous to the contribution of the regulator fields to the triangle anomaly in the trace of the energy momentum tensor [1.19]. This is

$$\langle p| \sum m_q \bar{q} q |p\rangle \simeq -\frac{\alpha_s n_h}{12\pi} \langle p| G_{\mu\nu} G^{\mu\nu} |p\rangle \left(1 + \mathcal{O}\left(\alpha_s \frac{M_N^2}{m_h^2}\right)\right) , \tag{1.3.2}$$

where n_h is the number of heavy quarks h, $G_{\mu\nu}$ is the gluon field strength tensor and α_s is the quark-gluon coupling constant. Since the nucleon mass M_N can be expressed in terms of the matrix element of the trace of the energy momentum tensor (including the anomaly) at vanishing momentum transfer:

$$\langle p|\Theta_{\mu\mu}|p\rangle = M_N \overline{u}(p)\ u(p) \tag{1.3.3}$$

with

$$\Theta_{\mu\mu} = \sum_q m_q \bar{q} q - \frac{1}{8\pi}\left(11 - \frac{2}{3} n_q\right) \alpha_s G_{\mu\nu} G^{\mu\nu} \tag{1.3.4}$$

From (1.3.2) and (1.3.4) one readily obtains

$$M_N \bar{u}(p) u(p) = \left(\frac{\alpha_s (n_q - n_h)}{12\pi} - \frac{11}{8\pi}\alpha_s\right) \langle p | G_{\mu\nu} G^{\mu\nu} | p \rangle \ , \tag{1.3.5}$$

yielding a scalar-nucleon coupling

$$\langle p | \sum m_q \bar{q} q | p \rangle = \left(\frac{2 n_h}{33 - 2 n_l}\right) M_N \bar{u}(p) u(p) \ , \tag{1.3.6}$$

where $n_l = n_q - n_h$ is the number of light quarks, taken to be u, d and s.

In this way, each heavy quark gives a contribution proportional to $\frac{2}{27} M_N$ to (1.3.2) and though, in the low energy scattering from a heavy nucleus[1], the contribution of each nucleon adds coherently producing large enhancements. As an example, the Higgs mediated scattering of a neutralino from a nucleus i containing A nucleons of mass M_N is

$$\sigma_i = \frac{8 G_F^2}{9\pi} \frac{M_W^2}{m_H^4} \alpha_H \left(\frac{2}{27} n_h M_N A\right)^2 \frac{m_\chi^2 m_i^2}{(m_\chi + m_i)^2} \ , \tag{1.3.7}$$

where the coupling α_H depends on the parameters of the supergravity model and will be discussed in chapter 4. Also a contribution of the form (1.3.2) appears when one consider the squark mediated scattering of neutralinos from nuclei, giving rise to a coherent interaction at low energies [1.16].

1.4-The cosmions and the solar neutrino problem

The only direct signals that we can receive from the nuclear reactions taking place in the interior of the sun are the neutrinos they produce, since the solar medium is transparent to these particles. Some reactions producing neutrinos are

[1] where the momentum transfer is less than the inverse of the nuclear radius

$$p + p \rightarrow D + e^+ + \nu_e \qquad E_\nu^{max} = 0.4\ MeV$$

$$^8\mathrm{B} \rightarrow {}^8\mathrm{Be}^* + e^+ + \nu_e \qquad E_\nu^{max} = 14\ MeV$$

The total energy production of the sun is dominated by the first one, but the ^{8}B reaction, although it provides only a tiny fraction of the total energy, is the only that produces neutrinos of sufficiently large energy as to allow their detection in the ^{37}Cl Davis experiment [1.20]. This experiment has been running for more than 15 years and obtained a count rate of 2.0±0.3 Solar Neutrino Units (1 SNU=$1 \cdot 10^{-36}$ events/(target atom×sec). The very high temperatures required to trigger the ^{8}B reaction are only achieved in the region of radius $r \leq 0.05 R_\odot$ in the central core of the sun. However, since the only measured solar temperature is that of the surface, our knowledge of the core temperature T_c comes from computations that assume a particular model for the solar interior. Bahcall [1.21] predicted, using the standard solar model, a rate for the ^{8}B reaction that would yield to the observation of 7.9(1±0.33) SNU in the Davis experiment, clearly not in agreement with the experiment. Recently, a confirmation of this neutrino deficit has also been obtained at the Kamioka underground detector [1.22].

To solve this puzzle, many attempts have been done. One of them, the cosmion solution, is based on the fact that a reduction in T_c by $\simeq$ 7% with respect to the standard solar model estimate would diminish the rate of the ^{8}B reaction just by the factor 1/3 necessary to account for the ~2 SNU observed. Faulkner and Gilliland [1.23] and Spergel and Press [1.24] proposed that one way to reduce T_c is by improving the energy transport from the neutrino producing regions to the outer regions of the sun with a population of weakly interacting massive particles. A significant concentration (more than 10^{-12} per solar nucleus) should be trapped in the sun. Press and Spergel [1.25] proposed that these particles have been captured by the sun from the halo of our galaxy.

In order to be captured and conduce heat efficiently in the sun, the effective scattering cross section per nucleus should be about $\sigma_c = 4 \cdot 10^{-36} cm^2$. With this cross section, the mean cosmion free path in the sun is of the order of the solar radius, giving a high probability for an halo cosmion to be trapped in a traversal through the sun, and this also optimize the heat conduction. The cosmion mass should be between 4 and 10 GeV. If lighter, it evaporates from the sun and if heavier, the orbits are too small to produce enough heat transport away from the core. To accumulate the required number of cosmions, these particles must be stable and not annihilate fast in the sun. This last requirement is severe, since requiring that the number of annihilations at present be less than the cosmions captured, per unit time, leads to [1.26]

$$\langle \sigma v \rangle_A \leq 10^{-4} \sigma_{scatt} \tag{1.4.1}$$

where $\langle \sigma v \rangle_A$ is the annihilation cross section times the relative cosmion velocity averaged in phase space. Since $\langle \sigma v \rangle_A$ is generally obtained by crossing the scattering cross section, they are usually of the same order and (1.4.1) is difficult to satisfy. There are mainly three ways to suppress annihilations in the sun: *i*) The annihilation cross section is suppressed, as could be the case of Majorana particles for instance; *ii*) the trapping rate of cosmions and anticosmions in the sun are different, yielding an asymmetry in their numbers only in the sun; *ii*) a cosmic asymmetry is assumed, i. e., the number of cosmions and anticosmions in the universe are different. In the last two cases, the annihilation rate in the sun is determined by the minority component.

A very remarkable property of the above mentioned bounds is that they can be made compatible with a cosmion relic abundance suitable to account for the dark matter.

Krauss et al.[1.26] showed that none of the "conventional" dark matter candidates have the expected properties. Namely, the standard weak cross section is several orders of magnitude smaller than required. Exotic cosmion candidates were proposed. A spin $\frac{1}{2}$ neutral particle interacting with nucleons via the exchange of a colour triplet scalar of mass arround 100 GeV [1.27]. The magnino[1.28], a neutral fermion which interacts through an anomalous magnetic moment . The $EXon$[1.29], a fourth generation neutrino interacting through a very light (500-1000 MeV) standard Higgs boson (just with the coupling discussed in the previous section). Finally, a spin $\frac{1}{2}$ neutral particle with interaction mediated by an extra Z' neutral gauge boson, arising from an E_6 breaking with an extra $U(1)$ factor [1.30].

In most of these models, annihilations are supressed by means of a cosmion asymmetry, and it is interesting to note that if this asymmetry could be related to the baryonic asymmetry in the context of a Grand Unified Theory, since the cosmion mass is about ten times the nucleon mass, its contribution to the density of the universe could be enough to account alone for the missing mass.

A new proposal for cosmion, the old neutralino, will be presented in chapter 5.

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CHAPTER 2

Testing cosmions at Tevatron

Since many cosmion models have been proposed recently, an interesting problem to address to is the detection of these conjectured particles in order to see the constraints imposed by existing and future experiments. Cosmions from the halo of our galaxy should produce detectable signals in the next generation of Germanium and Silicon detectors [2.1], by depositing energy in collisions with nuclei in the crystal. Another possibility is to search for signals arising from the cosmions trapped in the sun. In fact, the annihilation of cosmions and anticosmions can give rise to energetic light neutrinos that could be seen at underground proton decay detectors, as will be discussed in the next chapter. Cosmion production in particle accelerators is also possible. Magninos could be tested in e^+e^- colliders [2.2], since they should couple strongly to photons. The E_6 solution instead, requires a light Z' boson (of mass around 60 GeV) that couples only to quarks, being though testable at hadron colliders [2.3].

This chapter considers the search at the Tevatron collider of the cosmions corresponding to a class of models proposed by Gelmini, Hall and Lin [2.4]. These models predict the existence of a scalar coloured triplet with a mass around 100 GeV, giving rise to jet events with large missing transverse momentum at hadron colliders working in the TeV range. It is found [2.5] that the expected signal is well above the Standard Model background, so that Fermilab experiments will be able to support or neatly rule out these models.

In the proposals of ref.[2.4], the cosmion X is assumed to be a spin 1/2 neutral fermion, which interacts with the up quark via exchange of a scalar coloured triplet ϕ with the following term in the lagrangian

$$g\overline{X}(1-\gamma_5)u\phi + h.c.. \tag{2.1}$$

Note that the interaction involves only the up quark of the first generation, in order to avoid large flavour changing neutral current contributions.

X satisfies all the requirements to become a cosmion [2.6-10]. It is stable and its mass can be chosen in the correct range of 4-10 GeV. The annihilation rate can be suppressed in different ways [2.4]: *i)* X is a Majorana fermion and therefore its annihilation rate is proportional to $(\frac{m_u}{m_X})^2$, or else p-wave suppressed. *ii)* X is a Dirac fermion and the halo is composed of much more cosmions than anticosmions. The annihilation rate is basically controlled by the minority component. *iii)* X is again Dirac, but the halo is cosmion symmetric. However, the solar capture rate for cosmions and anticosmions is different. A sizable concentration in the sun is preserved, after partial annihilation. In this case, besides the interaction (2.1), an analogous term involving the opposite chirality projector and a different coupling constant is needed. Although in this analysis only the interaction (2.1) is considered, this last modification would not change significantly the results.

The interaction term (2.1) gives rise to a non relativistic scattering cross section of a cosmion from a nucleus N

$$\sigma = \frac{\lambda}{\pi}\left(\frac{g}{m_\phi}\right)^4 \frac{m_N^2 m_X^2}{(m_N+m_X)^2}. \tag{2.2}$$

In the case of a Majorana cosmion, $\lambda = 3\xi^2$, where $\xi = .85$ is the axial charge associated with the up quark in the proton and the total cross section in the sun is dominated by the scattering from hydrogen. If X is a Dirac particle, the coherent interaction with the nucleons gives $\lambda = (A+Z)^2$, where A and Z are the atomic number and charge of the nucleus respectively. The helium contribution dominates now the total cross section. In both cases, one can easily achieve correct values for the cosmion cross section $\sim 4 \cdot 10^{-36} cm^2$ when the mass of the scalar triplet ϕ is around 100 GeV.

As already noticed in ref.[2.4], the presence of a scalar triplet in this mass range is well testable at a hadron collider operating in the TeV region. The aim of this chapter is to give precise predictions of ϕ production at Tevatron. The relevant subprocesses for ϕ production in $p\bar{p}$ collisions are:

$$gg \to \phi\phi^* \tag{2.3a}$$

$$q\bar{q} \to \phi\phi^* \tag{2.3b}$$

$$gu \to \phi X. \tag{2.3c}$$

The Feynman diagrams contributing to (2.3) are depicted in fig.1. Besides the pure QCD contributions, fig.1 shows also the effect of the interaction (2.1), represented as a dot.

The scalar triplet subsequently decays

$$\phi \to uX, \tag{2.4}$$

giving rise to a hadronic jet plus missing transverse momentum $\not{p}_T$ carried out by the undetected weakly interacting neutral X. The final signal will be an excess over the standard model background of unbalanced monojet and dijet events.

The cross section for $\not{p}_T$ events arising from ϕ production and subsequent decay, can be obtained through a Monte Carlo computation, that makes a convolution of the parton cross section for the processes (2.3) with their structure functions. The coupling constant g of the interaction (2.1) can be fixed by the condition of having an average cosmion scattering cross section from nuclei, eq.(2.2), in the correct range. This gives

$$\frac{g^2}{4\pi} \simeq .04 \left(\frac{m_\phi}{100GeV}\right)^2 \tag{2.5}$$

for a Dirac cosmion. In the case of a Majorana X, g^2 is about one order of magnitude larger, so that assuming the more conservative choice (2.5) one can obtain a lower bound on the expected signal.

Fig.2 shows the total cross section (or the number of events for an integrated luminosity of $1pb^{-1}$) at $\sqrt{s} = 2$ TeV, as a function of the ϕ mass [1]. The results are shown for three differents cuts on $\not{p}_T$, respectively 40, 60 and 80 GeV.

A reasonable range between 50 to 150 GeV is taken for m_ϕ. As m_ϕ increases, the interaction (2.1) must become stronger, according to eq.(2.5). Therefore, the enhancement of the cross section due to the term (2.1) (with respect to the pure QCD contribution) becomes larger as m_ϕ grows. This enhancement is more evident for higher cuts in $\not{p}_T$ because the largest contribution to high $\not{p}_T$ events comes from monojets, that are due to the interaction (2.1) through processes (2.3c).

The standard model background comes from Z, W and heavy flavour production with subsequent semileptonic decays, where the neutrinos carry the $\not{p}_T$ and the leptons either escape detection or are lost inside the jets. It is expected that the standard model signal should be smaller than 100 pb for $\not{p}_T > 40$ GeV or 10 pb for $\not{p}_T > 80$ GeV [2.12]. Therefore, the cosmion generated $\not{p}_T$ signal is well detectable in the full range of interest for m_ϕ.

A non conventional source of background is supersymmetry. The production of squarks and their decay into a quark and a stable neutral weakly interacting supersymmetric particle gives a similar experimental signal. In the limit of very massive gluinos and turning off the interaction (2.1), the squark cross section differs from the one relative to ϕ only by an overall factor of ten[2]. At large m_ϕ, the term (2.1) largely enhances the QCD contribution and neatly modifies the signal with respect to the supersymmetric case. Therefore, it may be possible to distinguish the ϕ from a squark since, due to the interaction (2.1), the ratio of monojets over dijets is expected to be larger.

In conclusion, this chapter presents the experimental signals at Tevatron collider predicted by the cosmion model proposed in ref.[2.4]. The cosmion interaction with matter is mediated by a scalar coloured triplet ϕ with mass around 100 GeV. ϕ can

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[1] The structure functions of ref. [2.11] (set 1) were employed

[2] This factor comes from the usual assumption that ten kinds of squarks are taken to be degenerate (the partners of two chiral states of five flavours).

be copiously produced at Tevatron and its decay into the cosmion gives rise to large $\not{p}_T$ events, well above the standard model background. If no excess of $\not{p}_T$ events is observed, experiments can rule out completely the model.

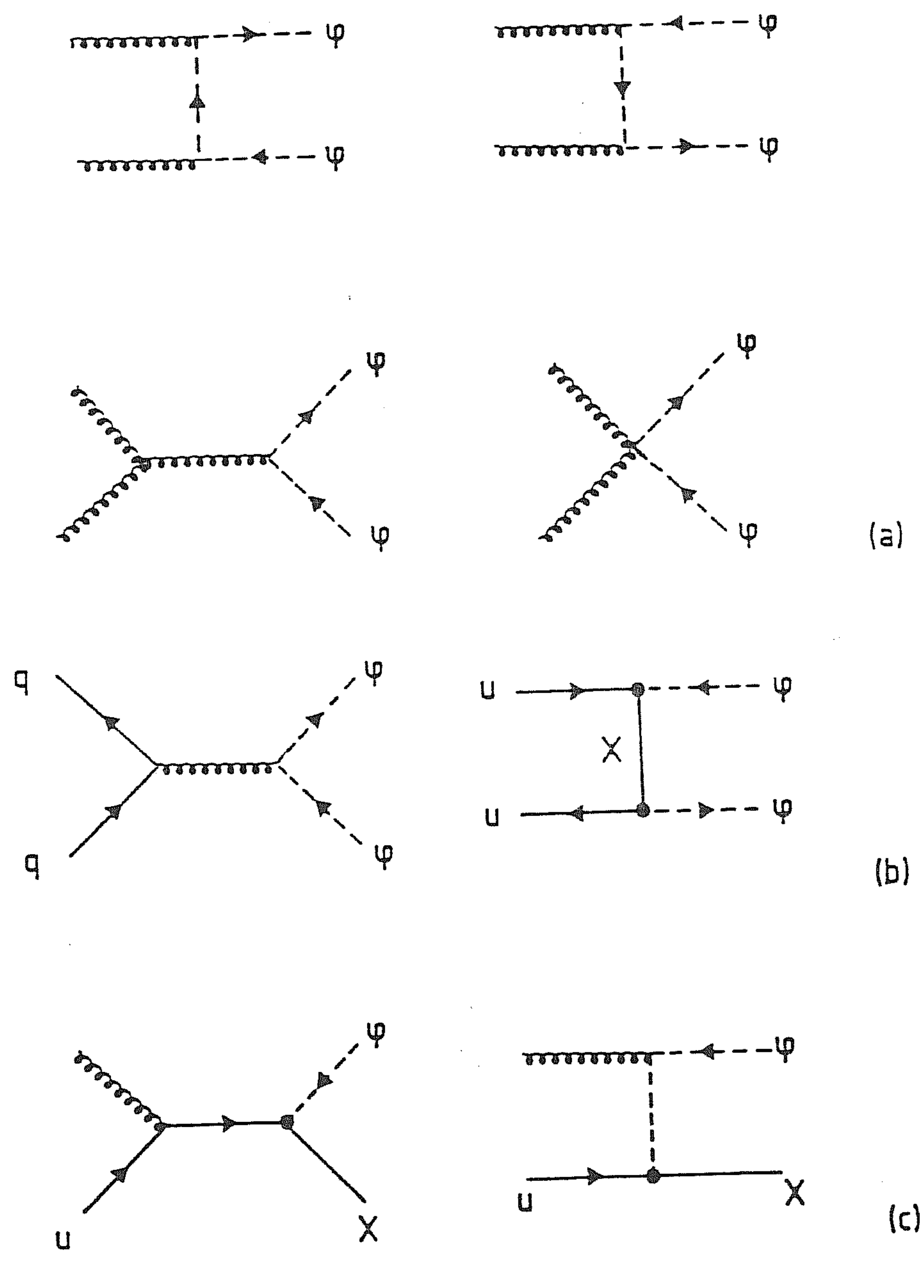

Fig.1. Feynman diagrams for the parton processes contributing to ϕ production. The dot represents the cosmion interaction (2.1).

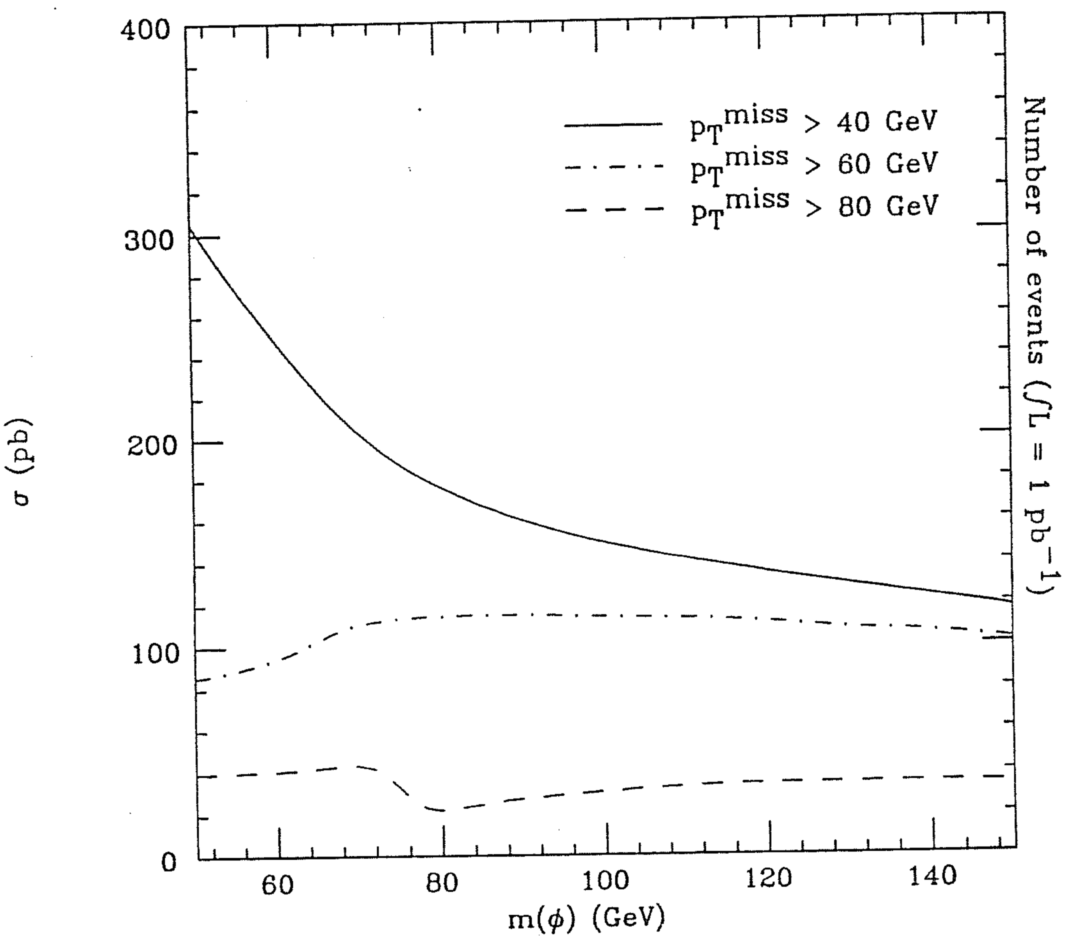

Fig.2. The cross section for $p\bar{p} \to$ jets + $\not{p}_T$ at $\sqrt{s} = 2$ TeV, due to ϕ production. The signal for three different cuts on $\not{p}_T$ is shown, as a function of the ϕ mass.

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CHAPTER 3

Testing cosmions at underground detectors

In this chapter, the indirect detection in underground proton decay detectors of the existing cosmion models is considered [3.1]. Severe constraints on their masses and on the cosmic asymmetry between cosmions and anticosmions result from these experiments for many candidates. Also bounds comming from cosmological and astrophysical requirements (that they do not overclose the universe and that they solve the solar neutrino puzzle) have been considered.

The large number of cosmions required to be trapped in the sun to solve the solar neutrino problem, can give rise to a large flux at the earth of neutrinos produced by cosmion-anticosmion annihilations. The neutrino detection in underground detectors (as Frejus, IMB or Kamioka) set limits on the flux comming from the direction of the sun [3.2] that, in turn, constrain the possible cosmion models. The best chances to observe the signal from the sun is when these neutrinos are very energetic ($E_\nu \sim$ GeV) since their cross section with the nuclei in the detector is proportional to E_ν, and also because the atmospheric background diminishes significantly for increasing energies [3.3]. This background of neutrinos at the earth is mainly due to the decay of muons or pions and kaons produced in the atmosphere by cosmic rays.

Energetic neutrino production is particularly important when the Z^o annihilation channel is allowed, since this can give rise (if the cosmion is a Dirac fermion) to direct prompt neutrinos of energy equal to the cosmion mass. This is the case for the $EXon$ [3.4], the magnino [3.5] and in a particular case also for the E_6 candidate [3.6][1]. For the other proposals [3.6-7], neutrinos are only produced by secondary decays, their mean energy being diminished and also their branching ratios being suppressed.

The flux of high energy neutrinos coming from the direction of the sun, detected at present in proton decay experiments, is of the order[3.8] of the expected atmospheric background. This background flux will be compared with the flux yield by cosmions, since a flux larger than this background by a factor of order one is already rejected. No attempt of a detailed comparison with experimental data will be done.

Attention will be focused on the $EXon$ [3.4], a fourth generation neutrino with only one Higgs doublet, since it is the simplest extension of the minimal Standard Model. The generalization to the other models will be done at the end of the chapter.

Griest and Seckel [3.9] have considered in detail the effect of a cosmological asymmetry for massive "standard" neutrinos (with scattering due only to Z^0 exchange), analysing relic abundances, the capture of neutrinos by the sun and the possibility of observing their annihilation products. The main difference in the cosmion case under study is the small value assumed for the Higgs boson mass. In fact, the non relativistic $EXon$ scattering cross section from a nucleus N_i is [3.4]

$$\sigma_i = \frac{8G_F^2}{27^2\pi}\frac{{n_h}^2}{{m_H}^4}\frac{\left(m_{N_i}m\right)^4}{(m_{N_i}+m)^2}, \tag{3.1}$$

where m and m_H are the cosmion and Higgs masses respectively. n_h is the number of heavy quarks that couple to the Higgs, corresponding in this model to c, b, t and the two new quarks of the fourth generation, *i.e.*, $n_h = 5$ (this is exactly the coupling discussed in **1.3**). The effective scattering cross section in the sun is obtained as a weighted sum

[1] Two of the three E_6 cosmion candidates are $SU(2)\times U(1)$ singlets and do not annihilate directly into light neutrinos. They correspond to the $SO(10)$ singlet and the $SU(5)$ singlet of the $SO(10)$ **16**, both in the **27** of E_6.

$\sigma_{eff} = \sum x_i\sigma_i$, where the solar number abundances x_i are tabulated in ref.[3.10]. A convenient way to express σ_{eff} is

$$\sigma_{eff} = \left(\frac{m_H^*}{m_H}\right)^4 4 \cdot 10^{-36} cm^2, \tag{3.2}$$

where m_H^* is the Higgs mass that gives the expected value for σ_{eff}. From eq.(3.1), given the values of x_i, it results that the largest contribution to σ_{eff} comes from scattering on He. The values of m_H^* range from $\sim 530 MeV$ for m=4 GeV to $\sim 1GeV$ for m= 10 GeV. The thermally averaged annihilation cross section is

$$\langle\sigma v\rangle_A = \frac{G_F^2 m^2}{2\pi}\left(N_A + O\langle v^2\rangle\right). \tag{3.3}$$

$N_A \simeq 7.4$ is the effective number of channels for the Z^0 s-wave annihilation[3.9] and the O$\langle v^2\rangle$ term corresponds to the p-wave annihilation (both through Z^0 and through Higgs exchange), giving a contribution essentially negligible since $\langle v^2/c^2\rangle \simeq 10^{-6}$ for cosmion thermalized in the sun.

The number of cosmions trapped in the sun (N_x) depends both on the capture of cosmions from the halo and on their annihilation in the sun. The range of masses considered ensures that there is no evaporation of cosmions from the solar surface. This gives

$$\dot{N}_{x(\bar{x})} = C_{x(\bar{x})} - C_A N_x N_{\bar{x}}, \tag{3.4}$$

where $C_x(C_{\bar{x}})$ give the capture rate of cosmions (anticosmions) and C_A gives the annihilation rate per cosmion pair. In this model, the annihilations are suppressed by means of a cosmological asymmetry between cosmions and anticosmions, so that the annihilation rate $\Gamma_A = C_A N_x N_{\bar{x}}$ is small due to the smallness of $N_{\bar{x}}$.

Using the expressions obtained by Press and Spergel [3.11], we get

$$C_{x(\bar{x})} \simeq f_G\ 1.2 \cdot 10^{30} n^h_{x(\bar{x})} \frac{f_\sigma}{v_{300}} \frac{cm^3}{sec} \tag{3.5a}$$

and

$$C_A \simeq 1.0 \cdot 10^{-54} \left(\frac{m}{10GeV}\right)^{7/2} \left[\frac{\langle\sigma v\rangle_A}{\frac{7.4}{2\pi} G_F^2 m^2}\right] sec^{-1} \tag{3.5b}$$

$n^h_{x(\bar{x})}$ is the local number density of cosmions (anticosmions) in the halo, v_{300} is the mean cosmion velocity in the halo in units of 300 Km/sec and f_σ takes into account the dependence of the capture on σ_{eff}. For $\sigma_{eff} < \sigma_c$, $f_\sigma \simeq .89\frac{\sigma_{eff}}{\sigma_c}$, and since for $\sigma > \sigma_{eff}$ all cosmions traversing the sun have a probability near one of being captured, f_σ saturates to the value one. The correction factor $f_G \simeq 1.5$ was obtained by Gould [3.12]. It accounts for the relative motion of the sun with respect to the isotropic cosmion distribution. The term in square brackets ($\simeq 1$ for EXons), is relevant for the other cosmion models. If cosmions provide the dark matter of the universe, their local density should provide the halo density $\rho^h \simeq .3$ GeV/cm^3. To be more general, we shall consider the situation in which part of the halo density is due to another type of cold dark matter (axions, neutralinos, etc.) that, however, do not become trapped efficiently by the sun or to dark matter in the form of jupiters or black holes, so that the only effect is to reduce the cosmion density in the halo to a fraction $\kappa = \frac{\Omega_x+\Omega_{\bar{x}}}{\Omega_{DM}}$ of ρ^h. An estimate of the cosmion mass density in the halo is then $\rho^h_{x(\bar{x})} = \kappa\frac{\Omega_{x(\bar{x})}}{\Omega_x+\Omega_{\bar{x}}}\rho^h$, where $\Omega_{x(\bar{x})} = \frac{\rho_{x(\bar{x})}}{\rho_{crit}}$ give the cosmological abundance of cosmions (anticosmions).

The relic cosmic densities $\rho_{x(\bar{x})}$ depend on the cosmion evolution in the early stages of the universe. After cosmions become non relativistic ($T < m$), the equilibrium number density is given by

$$n^o_x = 2\left(\frac{mT}{2\pi}\right)^{3/2} e^{-\frac{1}{T}(m-\mu)}, \tag{3.6}$$

where the chemical potential μ is related to the asymmetry $A = \frac{1}{2}(n_x - n_{\bar{x}})$ through $sh(\mu/T) = \frac{A}{n^o_{(A=0)}}$ (for anticosmions, the sign of the chemical potential reverses). As the universe cools further, the decoupling temperature T_d is reached when the interaction rate of cosmions becomes smaller than the expansion rate of the universe, and for $T < T_d$ cosmions "freeze out". The equation describing this evolution is

$$\frac{dn_{x(\bar{x})}}{dt} = -3\frac{\dot{R}}{R}n_{x(\bar{x})} - \langle\sigma v\rangle_A\left(n_x n_{\bar{x}} - n^o_x n^o_{\bar{x}}\right), \tag{3.7}$$

where n is the actual number density at time t and R is the cosmic scale factor. It is convenient to recast eq.(3.7) in the form

$$\frac{dY}{dz} = \frac{c}{z^2}\left(Y^o(Y^o + 2\alpha) - Y(Y + 2\alpha)\right). \tag{3.8}$$

In this last expression $Y = \frac{n_{\bar{x}}}{s}$ is the abundance of the minority component (assumed to be the anticosmions) scaled to the entropy density, $\alpha = \frac{A}{s}$ is the corresponding asymmetry, $z = m/s^{1/3}$ and

$$c = \left(\frac{b}{2\pi}\right)^{\frac{1}{2}} \langle \sigma v \rangle_A m \mathrm{M_p}, \tag{3.9}$$

where $b = \frac{2}{45}\pi^2 g_*$ and $\mathrm{M_p}$ is the Planck mass. The factor g_* counts the effective number of degrees of freedom of the particles that still interact at temperature T [3.9]. It was determined using the freeze out temperatures given in the ref. [3.13] and interpolating between 150 MeV and 250 MeV for the temperature of the QCD confinement phase transition .

The relic abundances were computed by a numerical integration of eq.(3.8) subject to the initial condition that long before the decoupling they coincide with the equilibrium values.

Returning to eq.(3.4), it is possible to estimate now the number of cosmions trapped in the sun. In the zero asymmetry case, it results

$$N = N_o th\left(\frac{t}{\tau}\right). \tag{3.10}$$

$\tau = (C_A C_x)^{-\frac{1}{2}}$ is the time at which the annihilation rate becomes comparable to the capture rate. For times larger than τ, N reaches asymptotically the value $N_o = (C_x/C_A)^{\frac{1}{2}} \simeq 5 \cdot 10^{38}$. This is less than 10^{-18} times the number of nuclei in the sun, certainly not enough to produce a sufficient heat transport away from the solar core. When asymmetry is turned on, since the resulting characteristic time is $\tau \ll \tau_\odot$ ($\tau_\odot \simeq 1.7 \cdot 10^{17} sec$ is the age of the sun), the annihilation rate rapidly equals the capture rate of anticosmions, resulting

$$N \simeq (C_x - C_{\bar{x}}) \cdot \tau_\odot \simeq \kappa f_\sigma \frac{\rho_{.3}}{v_{300}} \cdot \frac{7.2 \cdot 10^{-11}}{m} \frac{\alpha}{\alpha + Y_{now}} N_\odot \tag{3.11}$$

where Y_{now} is the present anticosmion abundance corresponding to an asymmetry α, $N_\odot$ is the number of nuclei in the sun, $\rho_{.3} = \frac{\rho^h}{.3 GeV/cm^3}$ and m is expressed in GeV. To solve the solar neutrino problem, *i.e.* to predict a count rate of 2.1 Solar Neutrino Units in the ^{37}Cl experiment, a particular value of N is needed for each specified values of σ_{eff} and of the cosmion mass. Some examples were evaluated in ref.[3.14], but in general cosmion concentrations larger than 10^{-12} per nucleus are needed [3.15]. From eq.(3.11) we see that such concentrations can only be achieved if the asymmetry α is larger than a minimum satisfying

$$\alpha_{min} = Y_{now}(\alpha_{min}) \cdot (f_\sigma \kappa \frac{\rho_{.3}}{v_{300}} \frac{72}{m} - 1)^{-1}, \tag{3.12}$$

Cosmion-anticosmion pairs annihilate into pairs of light neutrinos of type i with a branching ratio $f_i \simeq .07$, giving rise to a flux on the earth surface

$$\Phi_i = \frac{2 f_i}{4\pi R_{es}^2} \Gamma_A. \tag{3.13}$$

The factor of two accounts for the two neutrinos produced in the annihilation and the rate $\Gamma_A = C_A N_x N_{\bar{x}}$ can be equated to the anticosmion solar capture rate $C_{\bar{x}}$. R_{es} is the earth-sun average distance. There is another contribution to Φ_i coming from secondary ν's produced in the decay of other annihilation products. Although they provide approximately half of the flux, they are harder to detect since their mean energy is roughly $\frac{1}{3}m$. Thus, they will not be considered here.

There are several ways in which they can produce visible signals in a deep underground detector. $V_e(V_\mu)$ events are those in which the e (μ) is produced inside the detector volume [2]. When an electron is produced by an interaction in the surrounding rocks, it is rapidly stopped. The muons can instead traverse the rock, loosing only a fraction of their energy before reaching the detector, where they can be stopped (S_μ

[2] ν_τ signals will not be considered since τ production is kinematically suppressed for this range of masses.

events) or traverse it (T_μ events). To avoid confusion with high energetic atmospheric muons, only upward going μ should be considered. Thus, the acquisition of data for S_μ and T_μ events is restricted to night hours alone, giving a time efficiency of about 0.3 . The computation of the total event rate $R_T = V_e + V_\mu + 0.3(S_\mu + T_\mu)$ follows closely that of "standard" massive neutrinos done by Griest and Seckel[3.9]. Rates in a $(20m)^3$ water detector running for one year will be considered, for which the background due to neutrinos produced by cosmic ray air showers was computed to be $R_{back} \simeq 22$ in ref.[3.9].

The event rates due to solar neutrinos is shown in $fig.1-a$, as a function of the asymmetry and the $EXon$ mass, assuming $\sigma_{eff} = \sigma_c$ and $\kappa = 1$. Although the relic abundances remain approximately the same as for "standard" neutrinos, the larger capture rate of cosmions by the sun increases proportionally the annihilations in it. Thus, event rates of several thousands per year result. However, for large asymmetries $(10^{-10} - 10^{-9})$, the relic abundance of anticosmions starts decreasing exponentially with α [3.9], reducing the signal well below the background. In this case, the abundance of the majority component is fixed roughly to 2α. An upper bound on α results by requesting that the cosmion abundance does not overclose the universe. This upperbound on the asymmetry is shown in the $\alpha - m$ plane in $fig.1-a$ for $\Omega h^2 \equiv (\Omega_x + \Omega_{\bar{x}})h^2 = .25$ and 1. Assuming a value for $\Omega_{tot}h^2$, the zone above the curve corresponding to $\Omega h^2 = \Omega_{tot}h^2$ is forbidden.

The range of masses and asymmetries excluded by experiments are those for which the solar signal is larger than the atmospheric background, 22 events in this case, by a factor of order one. We can take the isoevent curve corresponding to one hundred events as the lower limit and $\Omega h^2 = 1$ as the upper limit. Only the wedge between these two curves is not excluded. A decrease in the assumed maximal abundance to $\Omega h^2 = .25$, diminishes the allowed region drastically. Instead, this region is not very sensitive to the number of events chosen to get the lower limit. We see in fig. 1 that only a small range of asymmetries around the value $\alpha \sim 10^{-10}$ is not ruled out. This range diminishes as the cosmion mass decreases. For comparison, the baryon asymmetry is $\alpha_B = 2 - 7 \cdot 10^{-11}$ [3.16]. As a side remark, this is just the right ballpark to allow to relate both, the baryon and the cosmion asymmetry, in the context of a Grand Unified Theory.

In the case cosmions only contribute a fraction of the halo density ($\kappa < 1$), the event rate decreases just by a factor κ (since it is proportional to the cosmion density), while α_{min} increases, instead, approximately as κ^{-1}. Both these effects change the lower bound on α, however, the most important change is in the upperbound. For a given Ω_{tot}, the upperbound on α decreases, since the maximum of Ωh^2 is now $\kappa\Omega_{tot}h^2$. Thus, for values of $\kappa \leq 0.2$ there will be no allowed region.

The reliability of the above results depends on the values assumed for the magnitudes involved in the computation of the capture rates. A mean cosmion velocity in the halo $\bar{v} = 300Km/sec$, a local halo density of .3 GeV/cm^3 and $\sigma_{eff} = \sigma_c$ were assumed. The error in $\bar{v}$ is about 15% and model studies [3.17] indicate that probably $.2GeV/cm^3 < \rho^h < .43GeV/cm^3$. Both efects could amount to a 50 % change in the event rate. However, a change of this order produces only a slight modification in the lower bound of the allowed zone of the $\alpha - m$ plane. If $\sigma_{eff} \neq \sigma_c$, the event rates shown in fig.1 should be scaled with f_σ. This factor can not be too small, since for instance $\sigma_{eff} \simeq 0.1\sigma_c$ would require a cosmion concentration in the sun $\geq 10^{-11}$ in order to provide the heat conduction mechanism necessary to yield the required solar neutrino flux [3.15]. This is larger than the maximum achievable by capture from an halo cosmion distribution (see eq.(3.11), with $f_\sigma \simeq 0.1$).

As a sumary, the EXons are allowed only if their relic density $\Omega h^2 \geq .2$. For $\Omega h^2 \simeq .2$ only $m \simeq 10$ GeV is allowed (with $\alpha \simeq 4 \cdot 10^{-11}$). To get $m \simeq 5$ GeV allowed, then $\Omega h^2 \simeq 1$ is required (a value that may be too large to be compatible with the lifetime of the universe).

What about the others cosmion candidates? There are two main changes from one model to the other.

i) The annihilation cross section.

ii) The branching to energetic light neutrinos.

Since the effective scattering cross section must always lie near σ_c, the total capture rate $C_x + C_{\bar{x}}$ could not be altered significantly. However, an increase in the cosmion annihilation cross section would reduce largely the anticosmion relic abundance (for a given α), and hence, the annihilations in the sun, that are proportional to $C_{\overline{x}}$ (eq.3.13). Point ii) is relevant mainly when cosmions do not annihilate through a Z^o into a pair of primary neutrinos. This is the case for some E_6 candidates (see footnote 1). It is also the case of the cosmion models of ref.[3.7], where annihilation proceeds via the t-channel exchange of a scalar colour triplet, producing two u quarks. These light quarks subsequently hadronize and the products loose their energy in the solar medium before undergoing a weak decay [3.18]. In all these cases, even if a large quantity of cosmions is trapped in the sun, they do not annihilate enough into energetic light neutrinos as to produce a detectable flux at the earth.

The magnino can annihilate both through a Z^o and a photon

$$\langle\sigma v\rangle_A = \langle\sigma v\rangle_{Z^o} + \langle\sigma v\rangle_\gamma + \langle\sigma v\rangle_{int} \tag{3.14}$$

where $\langle\sigma v\rangle_{Z^o} = \frac{G_F^2}{2\pi} m^2 N_A \simeq 4\cdot 10^{-9} m_5^2 GeV^{-2}$ as before, and

$$\langle\sigma v\rangle_\gamma = \frac{\pi\alpha^2\mu^2}{m^2}\sum_j Q_j^2 (1-\frac{m_j^2}{m^2})^{\frac{1}{2}}(1+\frac{m_j^2}{m^2}) \simeq 5\cdot 10^{-9}/m_5^2 GeV^{-2} \tag{3.15}$$

is the photon contribution, and the interference term is

$$\langle\sigma v\rangle_{int} = \sqrt{2} G_F \alpha\mu \sum_j Q_j g_V^j (1+\frac{m_j^2}{2m^2}) \simeq 2\cdot 10^{-9} GeV^{-2}, \tag{3.16}$$

The sums above run over all the final particles of charge Q_j and vector coupling g_V^j. We have taken the magnetic moment $\mu \simeq \frac{1}{8\pi^2}$ needed to obtain the required value for the scattering cross section [3.5]. For a magnino of mass near 10 GeV, $\langle\sigma v\rangle_A$ is mainly $\langle\sigma v\rangle_{Z^o}$. It is increased by only a few percent with respect to the $EXon$ case. Thus, the results plotted in fig. 1a are essentially unaltered in this range of masses. For magnino masses near 5 GeV, the contribution of photons is large, $\langle\sigma v\rangle_\gamma \simeq \langle\sigma v\rangle_{Z^o}$, thus, the total annihilation cross section is doubled (even for these masses, the interference contribution is less than 20% of the total annihilation cross section). As a consequence, the freeze out temperature decreases and the resulting relic abundance is smaller. Also the branching ratio to primary neutrinos is reduced since photons only annihilate into charged fermions.

Finally, the remaining E_6 candidate is a neutral fermion with a similar coupling to Z^o as ordinary neutrinos. In addition, it couples to the Z' boson. The Z' mass is adjusted so that the factor $\left(\frac{g'^2}{M_Z'^2}\right)^2$ (appearing in the scattering cross section from nuclei) is approximately 20 times larger than the corresponding one for Z^o scattering. The annihilation cross section has also the same factor, but there are however less annihilation channels available since the Z' should not couple to ordinary leptons[10], otherwise it would be too light to have escaped detection in neutral current experiments. Both opposite factors will typically result in $\langle\sigma v\rangle_A \simeq 10\langle\sigma v\rangle_{Z^o}$. Hence, for this model the relic anticosmion abundance is largely reduced and also the branching to light neutrinos diminishes.

Figs. $1b - c$ show the event rates in underground detectors for the magnino and the E_6 cosmion respectively, assuming $\sigma_{eff} = \sigma_c$ and $\kappa = 1$. Large regions in the $\alpha - m$ plane are ruled out, but the allowed region is wider than in the $EXon$ case, specially for the E_6 model. For $\kappa < 1$, the same modifications as in the $EXon$ case result.

In conclusion, all the existing cosmion candidates have been tested through their indirect detection in underground proton decay experiments. The large trapping of cosmions by the sun needed to explain the solar neutrino puzzle could give rise to a flux of energetic light neutrinos, produced by cosmion annihilations, exceeding the atmospheric background. This is the case for the $EXon$, the magnino and one E_6 candidate, for which the allowed region in the $\alpha - m$ plane is severely restricted by present experimental results and cosmological and astrophysical constraints.

FIGURE 1–a

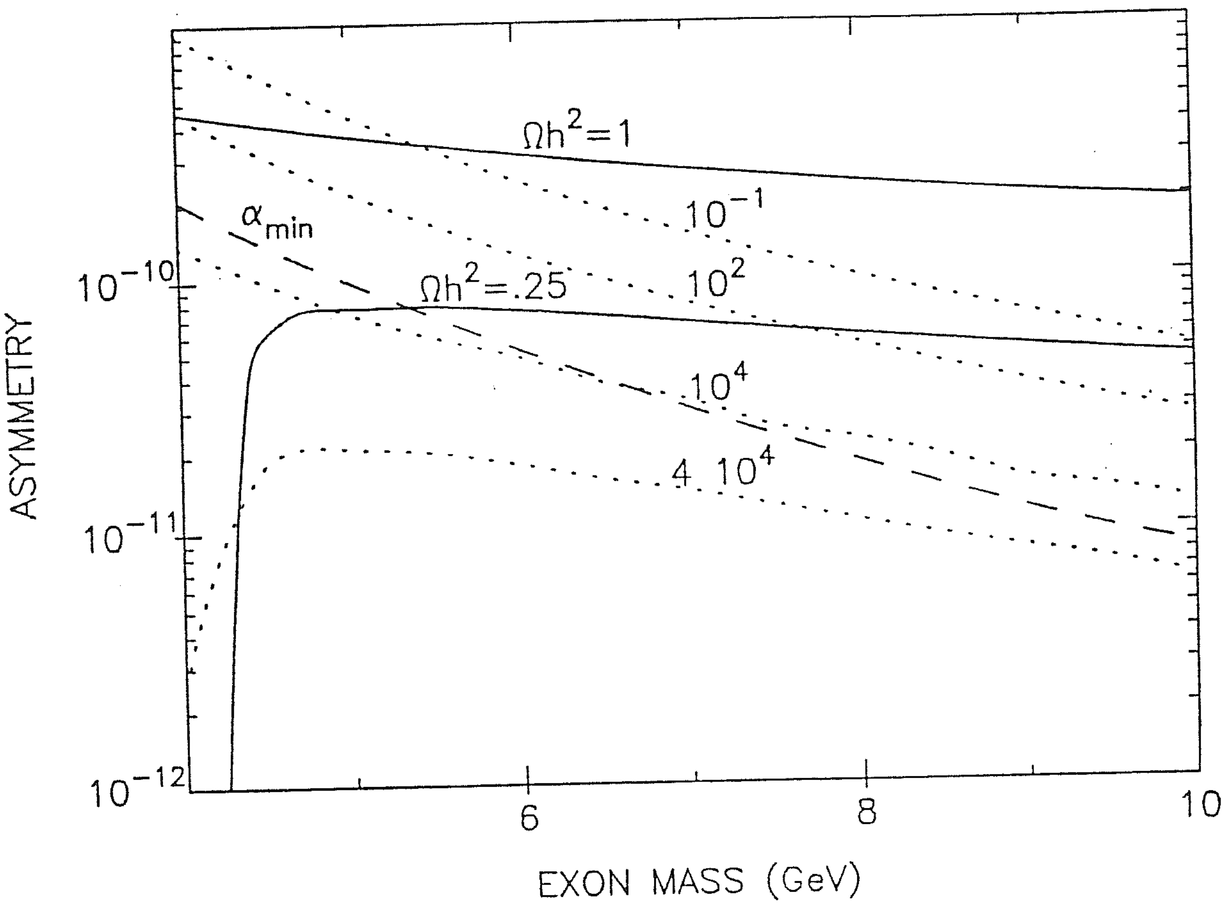


Fig.1. Number of events for underground detection and cosmological constraints as a function of the cosmion asymmetry α and the cosmion mass. Dotted lines correspond to isoevent curves for $10^{-1}, 10^2, 10^4$ and $4 \cdot 10^4$ events in a $(20m)^3$ water detector running for one year. Bold curves are isoabundances for $(\Omega_x + \Omega_{\bar{x}})h^2 = .25$ and 1. The minimum asymmetry (α_{min}) required to solve the solar neutrino problem is reppresented with a dashed line. Fig. 1-*a* corresponds to the fourth generation neutrino (*EXon*), fig. 1-*b* to the magnino and fig. 1-*c* to an E_6 cosmion candidate.

FIGURE 1–b

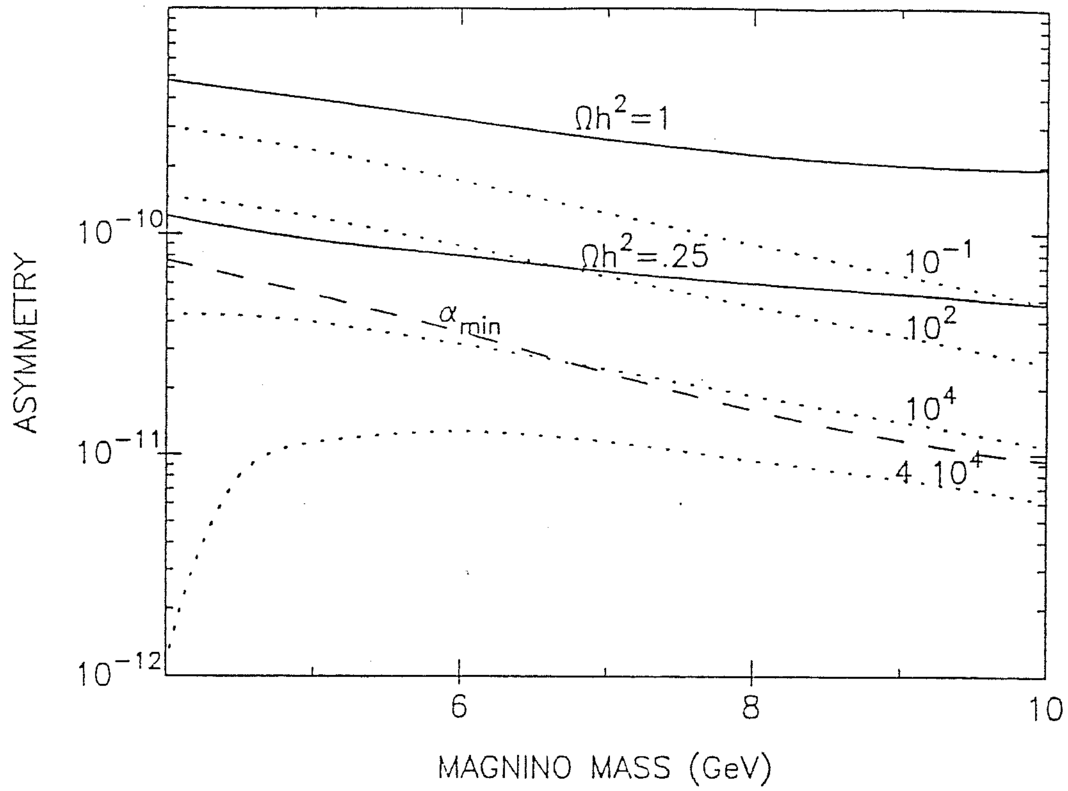
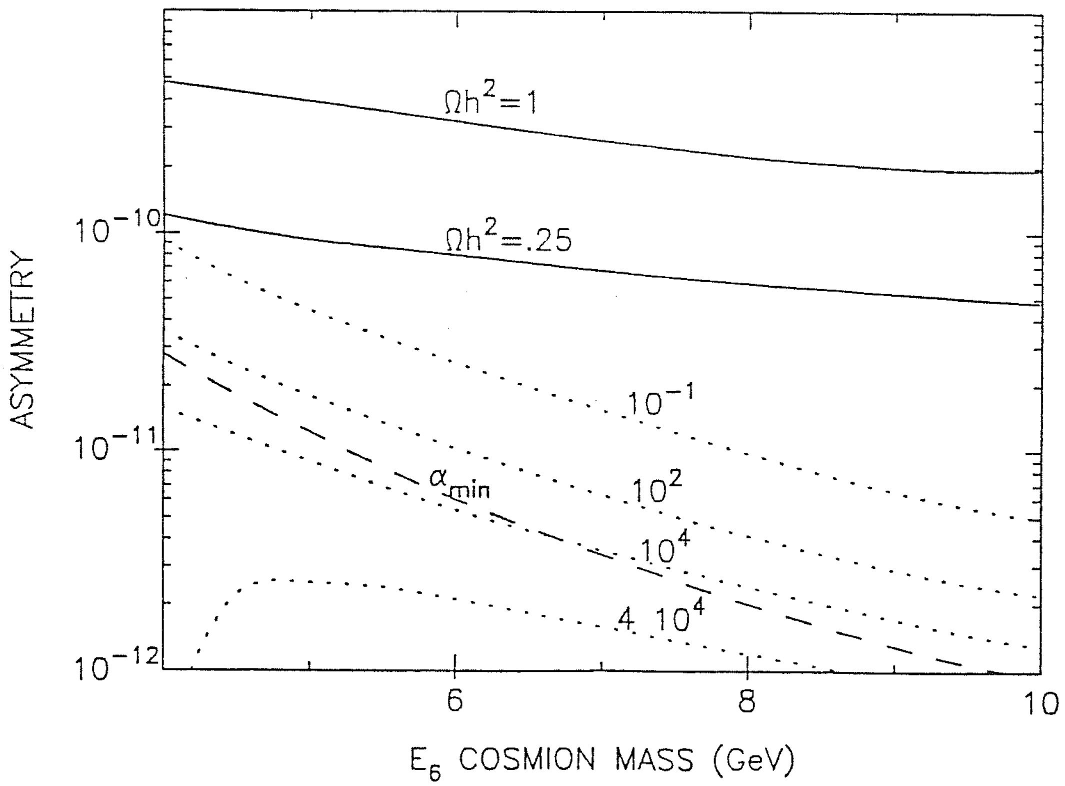


FIGURE 1–c



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CHAPTER 4

Testing neutralinos at underground detectors

4.1-Introduction

As discussed in **1.3**, the exchange of scalars can give rise to coherent interactions of neutralinos with matter. Griest [4.1] was the first to consider this effect for squark exchange, obtaining large enhancements in the coupling of neutralinos to heavy nuclei. However, because of the present experimental limits [4.2] and the theoretical predictions [4.3], squarks are expected to be rather heavy ($m_{\tilde{q}} \geq 95$ GeV) , so that the squark contribution could be very suppressed by the propagators. For this reason, Barbieri, Frigeni and Giudice [4.4] suggested that the main contribution to coherent neutralino-nucleus cross section comes from the exchange of a scalar Higgs boson. After all, a natural description of light scalars is the main motivation for introducing supersymmetry in the world of elementary particles at low energies [4.3]. In turn, the realistic supersymmetric models necessarily predict a light Higgs boson, often much lighter than the Z^0 [4.5].

In this chapter, the effect of Higgs boson exchange on the capture of dark matter neutralinos by the sun and the earth is considered [4.6]. If the Higgs boson is lighter than about 30 GeV, a large enhancement of the flux of energetic neutrinos coming from annihilations of the trapped neutralinos is predicted. Also the expected neutrino energy spectrum and the rates for neutrino generated events in underground detectors is studied.

4.2-Neutralino capture by the sun and the earth

Let us consider a neutralino χ with mass m_χ, mean velocity v_χ and mass density ρ_χ in the halo of our Galaxy. The capture rate for a body (sun or earth) of mass M_B is given by [4.7]

$$C = \left(\frac{6}{\pi}\right)^{\frac{1}{2}} \frac{\rho_\chi}{v_\chi} M_B \sum_i \langle v_{esc}^2 \rangle_i \frac{\sigma_i}{m_\chi m_i} f_i S_i. \tag{4.2.1}$$

The sum runs over all the elements contained in the body. f_i is the fraction of M_B due to the element of mass m_i, σ_i is the cross section of the neutralino with the nucleus of kind i. $\langle v_{esc}^2 \rangle_i$ is the square escape velocity mediated over the distribution of the element. Finally, S_i is a factor which takes into account the suppression due to the element mass mismatching with the neutralino and the possible lack of coherence of the interaction. The effect of the mismatching is particularly important for the earth since in this case the escape velocity is low, so that neutralinos are trapped only if they loose a sizable amount of their energy in the collisions. This is possible only if m_χ and m_i are sufficiently close. The effect of the lack of coherence becomes important when the transfer momentum is not negligible with respect to the inverse of the characteristic dimension of the nucleus. This has to be taken into account for heavy neutralinos. The analytical expression for S_i can be found in ref. [4.7][1].

As illustrated in the introduction, we will assume that squarks are heavy, focusing the attention to the exchange of a light Higgs boson. The relevant interactions are the Higgs-neutralino coupling:

$$\frac{g}{2} F H^0 \overline{\chi}\chi, \tag{4.2.2}$$

the Higgs-quark coupling:

[1] For the capture of light halo particles by the sun, eq.(4.2.1) reduces to the more simple expression used in (3.5a)

$$-\frac{g}{2}K_q\frac{m_q}{M_W}H^0\overline{q}q, \tag{4.2.3}$$

and the Z^0-neutralino coupling:

$$\frac{g}{4cos\theta_W}NZ_\mu\overline{\chi}\gamma^\mu\gamma_5\chi, \tag{4.2.4}$$

where g is the gauge coupling constant. The coefficients F, K_q and N are given once the supergravity model is specified. In section 4, we will show how they are connected to the fundamental parameters of the theory and give the numerical predictions of the minimal supergravity model.

Using the result of **1.3**, one finds that the Higgs exchange gives a coherent neutralino scattering cross section from a nuclei of mass m_i (containing A nucleons of mass m_N) [4.4]:

$$\sigma_i = \frac{\mathrm{G_F}^2}{2\pi}\alpha_H\left(\frac{8}{27}\frac{m_N M_W}{m_H^2}A\right)^2\frac{m_\chi^2 m_i^2}{(m_\chi + m_i)^2}, \tag{4.2.5}$$

$$\alpha_H \equiv F^2(2K_u + K_d)^2. \tag{4.2.6}$$

Assuming $\rho_\chi = 0.4$ GeV/cm^3 and $v_\chi = 300$ Km/sec, the Higgs exchange gives a capture rate:

$$C = c\alpha_H\left(\frac{10\ GeV}{m_H}\right)^4\sum_i A^2\phi_i f_i S_i\frac{4m_\chi m_i}{(m_\chi + m_i)^2}, \tag{4.2.7}$$

where $c = 8.4\cdot 10^{15}sec^{-1}$ for the earth and $8.7\cdot 10^{24}sec^{-1}$ for the sun. $\phi_i \equiv \langle v_{esc}^2\rangle_i/v_0^2$, where v_0 is the escape velocity at the surface of the body (11.2 Km/sec for the earth and 618 Km/sec for the sun).

For the sun, it is important to take into account also the axial interaction of the neutralino with hydrogen via Z^0 exchange. The corresponding cross section is

$$\sigma = \frac{\mathrm{G_F}^2}{2\pi}\alpha_Z 3g_A^2\frac{m_\chi^2 m_p^2}{(m_\chi + m_p)^2}, \tag{4.2.8}$$

$$\alpha_Z \equiv N^2, \tag{4.2.9}$$

where m_p is the proton mass and $g_A = 1.25$. This leads to a capture rate from hydrogen

$$C = c' \alpha_Z \phi_H f_H S_H \frac{4 m_\chi m_p}{(m_\chi + m_p)^2}, \tag{4.2.10}$$

where, for the sun, $c' = 7.8 \cdot 10^{26} sec^{-1}$.

Fig. 1a shows the predicted capture C of neutralinos by the earth, assuming $\alpha_H = 1$ (it is easy to rescale the results for different values of α_H, see eq.(2.7))[2]. The origin of the peaks in correspondence with different element masses is due to the mass mismatching effect [4.7], which has already been discussed. The largest capture is achieved near the "iron resonance". Unfortunately, this range of neutralino masses seems to be cosmologically disfavored since it is close to the condition $m_\chi \simeq M_Z/2$. In this case, the annihilation through Z^0 exchange is very effective and the relic density of neutralinos is expected to be rather small [4.1,4].

In fig. 2b, we show the capture rate for the sun, rescaled by the factor $\left(\frac{R_\oplus}{1A.U.}\right)^2 = 1.8 \cdot 10^{-9}$, to allow a clearer comparison with the result from the earth[3]. It is also shown (dashed line) the prediction for the spin dependent Z^0 interaction, neglecting the H^0 exchange. We have assumed $\alpha_H = 1$ and $\alpha_Z = 0.1$. These are reasonable values achieved in the minimal version of the supergravity model (see sect. 4).

The existence of a light Higgs boson has the effect of efficiently trapping neutralinos in the earth (Fig. 1a). Due to the presence of heavy elements in the sun, also the solar capture rate is sizably enhanced for $m_H = 10 - 20$ GeV (Fig. 1b). With these values of capture rates, one expects a visible flux of energetic neutrinos reaching the detectors.

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[2] For the computation of f_i and S_i the composition of the earth from ref. [4.8] was used, and for ϕ, the values 1.2 and 1.6 were used for the elements in the mantle and the core respectively [4.7].

[3] For the sun, the elements abundances from ref. [4.9] were used, and for Φ_i the values 3.16, 3.40 and 3.23 were taken for hydrogen, helium and the other elements respectively [4.7].

4.3-Neutrino fluxes and detection

The increase of neutralino population in the sun (or earth) due to the capture is balanced either by evaporation or by annihilation. The evaporation is significant only for neutralinos lighter than ~ 3.5 GeV for the sun [4.10,11] and ~ 11 GeV for the earth [4.7]. Such a light χ will not be considered here, since it may also lead to an overclosure of the universe [4.12]. The equilibrium between capture and annihilation is reached in a time which is generally much shorter than the solar age [4.13]. Therefore, we consider a condition of equilibrium, where the capture rate is twice the annihilation rate.

The differential flux of neutrinos, at a distance R from the source, is given by:

$$\frac{d\Phi}{dE_\nu} = \frac{\frac{1}{2}C}{4\pi R^2} \sum_f B_{\chi f}(m_\chi) \frac{dN_{f\nu}}{dE_\nu}. \tag{4.3.1}$$

The capture C is given in eq.(4.2.1) and $B_{\chi f}$ is the branching fraction of $\chi\chi \rightarrow f\bar{f}$. $\frac{dN_{f\nu}}{dE_\nu}$ is the energy distribution of the neutrino yield from $f\bar{f} \rightarrow \nu +$ "*anything*" and the sum runs over all possible particles f. Eq.(4.3.1) holds for any kind of neutrinos (and antineutrinos).

$B_{\chi f}$ depends on the interactions felt by the neutralino. Since χ is the lightest supersymmetric particle, f describes all standard fermions and Higgs bosons allowed by phase space. We assume that only one Higgs boson is light ($m_H < m_\chi$). In certain cases, the minimal supergravity model predicts also one light Higgs pseudoscalar. For simplicity, we disregard this possibility. We consider the "cold annihilation", where the velocity of the trapped neutralino is negligible and s-wave terms dominate. In this case, $\chi\chi \rightarrow H^0 H^0$ and $\chi\chi \rightarrow f\bar{f}$ through H^0 exchange (where f is an ordinary fermion) are forbidden and we are left only with $\chi\chi \rightarrow f\bar{f}$ through Z^0 exchange. Thus

$$B_{\chi f} \propto N_f m_f^2 \sqrt{1 - \frac{m_f^2}{m_\chi^2}}, \tag{4.3.2}$$

where N_f counts the number of colors of the fermion f.

Light hadrons (containing only u, d and s quarks) and light charged leptons (e, μ) will be stopped in the sun or the earth before they can decay. The subsequent neutrinos have an unobservable low energy spectrum [4.14]. Moreover, note that the production of light fermions is suppressed by the factor m_f^2 (see eq.(4.3.2)). On the contrary, the effects of energy loss or neutrino absorption for hadrons containing heavy quarks and for the τ lepton are negligible for neutralino masses lower than 100 GeV [4.14]. Assuming that the annihilation in a t quark pair is kinematically forbidden, only $f = \tau, c, b$ contribute in (4.3.1).

We will take for the neutrino distribution the approximate expression [4.14]:

$$\frac{dN_{f\nu}}{dE_\nu} = \frac{1}{m_\chi Z_{Ff}} \int_{\frac{E_\nu}{Z_{Ff} m_\chi}}^{1} dy \sum_0^4 a_{f\nu}^{(n)} y^{n-1}. \tag{4.3.3}$$

Eq. (3.3) holds when the heavy quark hadron is in the relativistic regime, *i.e.* $m_\chi \gg m_f$. The coefficients $a_{f\nu}^{(n)}$ and Z_{Ff}, taken from ref. [4.14], are listed in table I. They were obtained by a fit of Monte Carlo results considering the fragmentation of the jet produced by the heavy quarks and other final state QCD effects. This reduces the naive estimates of the mean neutrino energy in the decay of b and c quarks by factors of 2-3, since the heavy quark must share the "beam" energy m_χ with the others jet constituents and also because there are generally more than three particles in the final state. The coefficients in table I were obtained with ν_μ, and we will assume an equal yield also for ν_e (and for antineutrinos).

f	Z_F	a_0	a_1	a_2	a_3	a_4
τ	1.00	0.0005	-0.0104	2.1250	-2.0750	-0.0669
c	0.55	0.0097	-0.3917	4.6487	-8.6226	4.3834
b	0.71	-0.0629	5.8391	-22.5507	32.5397	-16.5512

Table I : The coefficients Z_{Ff} and $a_{f\nu}^{(n)}$ (in the case of ν_μ) for $f = \tau, c, b$, as taken from ref. [4.14].

Fig. 2 shows the differential flux of electron (or muon) neutrinos and antineutrinos, for some values of the neutralino mass. Fig. 2a refers to the terrestrial flux, assuming

$m_H = 10$ GeV and $\alpha_H = 1$. The rescaling to different values of m_H and α_H is staightforward (see eqs.(4.2.7), (4.3.1)). Fig. 2b refers to the solar flux, when $m_H = 10$ GeV (solid line) and when the Higgs exchange is neglected (dashed line). As before, we have taken $\alpha_H = 1$ and $\alpha_Z = 0.1$. Because of the fragmentation effect, the average heavy quark hadron energy is $\langle E_H \rangle = Z_F m_\chi$, being reduced with respect to the naive expectation $E_H = m_\chi$. This explains why the dominant b contribution to the neutrino spectrum drops below $E \simeq 0.71 m_\chi$, leaving only the energetic tail due to τ decays. This feature is clearly apparent in the shape of the curves of fig. 2.

The background from atmospheric neutrinos has been discussed in refs. [4.15,16]. Contrarily to the signal shown in fig. 2, the background is concentrated at low energies and therefore dark matter neutrinos can be identified using spectral and angular information.

Let us turn to discuss the detection of the neutrino flux. The event rate for electrons (or muons) generated within the detector due to neutrino interaction is obtained by folding the flux (4.3.1) with the neutrino cross section:

$$R_e = \int_{E_{th}}^{m_\chi} dE_\nu \frac{d\Phi}{dE_\nu} \int_{E_{th}}^{E_\nu} dE_e \frac{d\sigma_{\nu+\bar{\nu}}}{dE_e}(E_\nu, E_e) \tag{4.3.4}$$

where E_e is the electron (muon) energy and E_{th} is the minimum energy necessary to identify it. For E_ν larger than the proton mass, the differential charged current cross section for neutrinos and antineutrinos per nucleon is [4.16]:

$$\frac{d\sigma_{\nu+\bar{\nu}}}{dE_e}(E_\nu, E_e) \simeq \sigma \left(1 + a\frac{E_e^2}{E_\nu^2}\right), \tag{4.3.5}$$

with $\sigma = 0.81 \cdot 10^{-38}$ cm^2 GeV^{-1} and $a = 0.93$. Replacing eq. (4.3.5) in eq. (4.3.4) and taking into account the number of target nucleons, we find:

$$R_e = 154 \frac{events}{Kt\ yr} \int_{E_{th}}^{m_\chi} dE_\nu \frac{d\Phi}{dE_\nu} E_\nu \mathcal{I}\left(\frac{E_{th}}{E_\nu}\right), \tag{4.3.6}$$

$$\mathcal{I}(x) = 1 + \frac{a}{3} - x - \frac{a}{3}x^3, \tag{4.3.7}$$

where the differential flux is expressed in cm^{-2} sec^{-1} GeV^{-1} and the energies in GeV. The solid line in fig. 3 describes the rate of events per kiloton year for the terrestrial and solar flux, taking $E_{th} = 2$ GeV. The same signal for electrons and muons is expected. However, the search for electrons is much more efficient, since the background is lower [4.16] (see also ref. [4.10]).

Energetic muon neutrinos can also be revealed through their interactions with the rock, producing muons which either pass through the detector or are stopped inside. The flux of muons with energy larger than E_{th} is given by [4.16]:

$$R_\mu = N_A \int_{E_{th}}^{m_\chi} dE_\nu \frac{d\Phi}{dE_\nu} \int_{E_{th}}^{E_\nu} dE_\mu \int_{E_\mu}^{E_\nu} dE'_\mu \frac{d\sigma_{\nu+\bar{\nu}}}{dE'_\mu}(E_\nu, E'_\mu) \frac{1}{\alpha(1+E_\mu/\epsilon)}, \tag{4.3.8}$$

where $\alpha \simeq 2$ MeV cm^2 / g and $\epsilon = 510$ GeV take into account the muon energy loss and N_A is the Avogadro's number. Substituting eq. (4.3.5) in eq. (4.3.8), we obtain the rate of muon events:

$$R_\mu = 77 \frac{events}{(10\ m)^2 yr} \int_{E_{th}}^{m_\chi} dE_\nu \frac{d\Phi}{dE_\nu} E_\nu^2 \mathcal{I}' \left(\frac{E_{th}}{E_\nu} \right), \tag{4.3.9}$$

$$\mathcal{I}'(x) = .73 - 1.3x + 0.5x^2 + .08x^4, \tag{4.3.10}$$

where a surface of 100 m^2 corresponds to a detector of 1 kiloton of water. This rate is plotted in fig. 3 (dashed line), taking $E_{th} = 2$ GeV. For the sun, eq. (4.3.9) has been multiplied by the factor 1/3. This takes into account the fact that the signal is detactable only during the night, because of the large background of muons from the atmosphere. This factor was omitted for the terrestrial signal, since it has been assumed that the neutralino annihilations occur in the core of the earth and the muons are always upgoing.

Fig. 3 shows that, for heavy neutralinos, *i.e.*, more energetic neutrinos, the event rate of muons produced in the rock outside the detector has a relative increase with respect to the rate of electrons generated within the detector. The large enhancement due to the light Higgs effect is apparent.

4.4-The couplings in the minimal supergravity model

In the previous sections, the parameters which specify the neutralino interactions, α_H and α_Z, were treated as free inputs, taking $\alpha_H = 1$ and $\alpha_Z = 0.1$ in the numerical calculation. However, they can be computed in terms of the fundamental parameters of the underlying supergravity theory.

Lets consider the softly broken supergravity model with the minimal choice of fields [4.3], where four neutralinos and two Higgs bosons are present. In this model:

$$\alpha_H = \left[(\gamma_3 sin\alpha + \gamma_4 cos\alpha)(\gamma_2 - \gamma_1 tan\theta_W)\left(2\frac{cos\alpha}{sin\beta} - \frac{sin\alpha}{cos\beta}\right)\right]^2 \qquad (4.4.1)$$

$$\alpha_Z = (\gamma_4^2 - \gamma_3^2)^2, \qquad (4.4.2)$$

where $tg\beta = \frac{v_2}{v_1}$ is the ratio of the two Higgs vacuum expectation values, and

$$cos2\alpha = c\left(\frac{c^2 + r^2 - 2r}{c^2 + r^2 - 2rc^2}\right) \qquad (4.4.3)$$

$$c \equiv cos2\beta, \qquad r \equiv \frac{m_H^2}{m_Z^2}. \qquad (4.4.4)$$

Finally, γ_i are the fractions of neutralino current states (b-ino, neutral w-ino, higgsinos) contained in the lightest neutralino mass eigenstate:

$$\chi = \gamma_1\tilde{B} + \gamma_2\tilde{W} + \gamma_3\tilde{H}_1 + \gamma_4\tilde{H}_2 \qquad (4.4.5)$$

For a pure photino ($\gamma_1 = cos\theta_W, \gamma_2 = sin\theta_W, \gamma_3 = \gamma_4 = 0$) or pure higgsino ($\gamma_1 = \gamma_2 = 0, \gamma_3 = sin\beta, \gamma_4 = cos\beta$) α_H vanishes, but these are rather extreme cases in the supersymmetric parameter space and, in general, sizable values of α_H are found.

The coefficients γ_i can be computed in terms of the gaugino mass M [4], the higgsino mass parameter μ and the ratio $\frac{v_2}{v_1}$, appearing in the neutralino mass matrix (see

[4] We are assuming a unification relation for gaugino masses and taking M as the SU(2) gaugino mass, renormalized at low energies.

M (GeV)	μ (GeV)	$\frac{v_2}{v_1}$	m_χ (GeV)	α_Z	α_H $(m_H{=}10GeV)$	α_H $(m_H{=}30GeV)$
200	60	2	11	.10	1.9	2.0
50	50	2	25	.53	1.7	1.4
100	-50	2	42	.18	0.2	0.1
200	150	2	60	.01	1.7	1.8
200	60	8	27	.19	9.7	10.3
100	-100	8	42	.06	5.3	3.6
200	150	8	74	.02	8.3	8.9

TableII: The couplings α_Z, α_H and the neutralino mass m_χ, for representative values of M, μ and $\frac{v_2}{v_1}$.

chapter 1). In table II, numerical evaluations of eq. (4.4.1)-(4.4.2) and the value of the neutralino mass for representative choices of M, μ and $\frac{v_2}{v_1}$ are given.

α_H turns out to be typically of order 1, even if larger values are achieved for large $\frac{v_2}{v_1}$. This is due to the fact that the third factor in eq. (4.1) increases approximately linearly with $tg\beta$ and therefore $\alpha_H \propto \left(\frac{v_2}{v_1}\right)^2$. Large $\frac{v_2}{v_1}$ values are generally predicted by a correct radiatively induced electroweak symmetry breaking, for heavy top quark. In this case, the predictions on the Higgs contributions shown in the figures (with $\alpha_H = 1$) are rather conservative and further enhancements are achievable.

α_Z is typically of order 0.1. It is always less than 1 because of orthogonality and it becomes smaller as the gaugino component dominates over the higgsino component. In this case, the spin dependent interaction mediated by the Z^0 is further suppressed.

In the minimal supergravity model, the mass of the scalar Higgs is bounded by:

$$m_H \leq M_Z|cos2\beta|, \tag{4.4.6}$$

justifying the interest for a light Higgs boson. Apart from another neutral scalar and one charged Higgs which are predicted to be heavy, the supergravity model contains a Higgs pseudoscalar with mass $m_H\sqrt{\frac{1-r}{c^2-r}} > m_H$. For large $\frac{v_2}{v_1}$, this becomes quite close to m_H and can become important in the computation of the neutrino energy spectrum.

The search for extra-atmospheric energetic neutrinos imposes bounds on the parameter space (M, μ, v_2/v_1) relevant to the neutralino, as main component of the galactic halo. In particular, an experimental sensitivity of 1 event/Kt yr for contained electron events and 10 events/100 m^2 yr for the flux of through-going muons (reasonable experimental bounds) can exclude a region which is slightly wider than the one derived from direct dark matter searches (see fig. 2 of ref. [4.4]). The shape of the excluded regions are very similar, since they both depend on the same coefficient α_H, eq. (4.4.1), and follow its behaviour. However, the indirect search for energetic neutrinos can turn out to be more efficient than the direct dark matter detection in studying the supersymmetric parameter space, in presence of a light Higgs boson. Finally, the search for dark matter neutralinos has also an interesting connection and complementarity with the planned experiments at the e^+e^- and hadron colliders.

4.5-Conclusions

In this chapter, the effect of a Higgs boson with a mass in the range 10-30 GeV on the capture of dark matter neutralinos by the sun and the earth has been considered, studying the neutrino flux induced by their subsequent annihilations. It should be remembered that the presence of light Higgs bosons and an effective supersymmetry at the Fermi scale are deeply connected. The description of light scalars free from the naturalness problem [4.17] needs the introduction of supersymmetry and, in the context of the most satisfactory models, supersymmetry necessarily implies a light Higgs [4.5]. The simultaneous presence of supersymmetry and a light Higgs boson easily leads to a scenario where the lightest neutralino is a good cold dark matter candidate with coherent interaction with nuclei [4.4]. Through this interaction, a large number of neutralinos is trapped by the sun and the earth. The relative capture rates were shown in fig. 1. In absence of spin independent interactions, essentially no capture is predicted for the earth, while the large fraction of hydrogen of the sun can still trap the neutralinos. Nevertheless, due to the presence of heavy elements, a sizable enhancement of the solar capture is expected for $m_H \leq 20$ GeV. The subsequent annihilation of trapped neutralinos provides a flux of neutrinos which can reach the underground proton decay detectors. The neutrino energy spectrum has been shown in fig. 2. The energetic tail is expected to be visible. These neutrinos can be detected through their charged current interactions which produce electrons or muons (τ production is kinematically suppressed). At first sight, electron search looks more promising because of the lower background. However, muon neutrinos can produce, through their interactions with the rock outside the detector, a muon shower reaching the apparatus. In fig. 3, the two event rates were compared, showing that the latter is relatively enhanced for heavy neutralinos. A final remark is that a typical neutralino coupling, as predicted by the minimal supergravity model was assumed in this investigation. Different couplings can be obtained by scaling the results. In section 4, the behaviour of these couplings as functions of the fundamental parameters of the minimal supergravity model was studied. The search for energetic neutrinos sets severe constraints on the parameter space relevant to a neutralino, as main component of the galactic halo.

Though, the existence of a Higgs boson lighter than about 30 GeV has important consequences for the detection of supersymmetric dark matter. This seems very interesting in view of the fact that this Higgs mass range is expected to be fully tested at SLC and LEP.

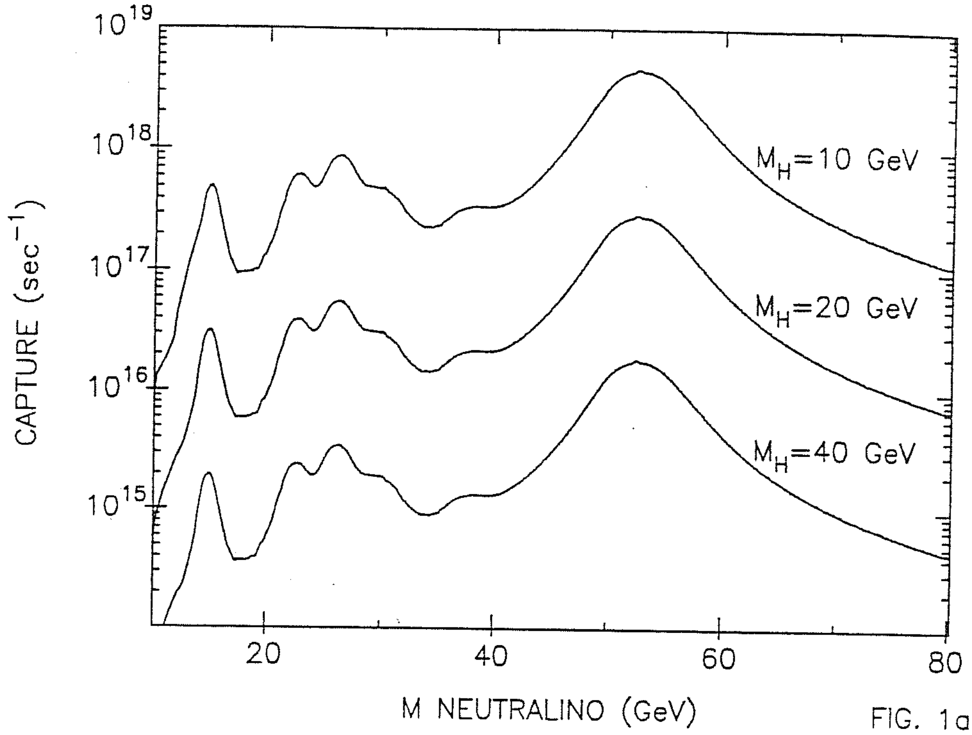

FIG. 1a

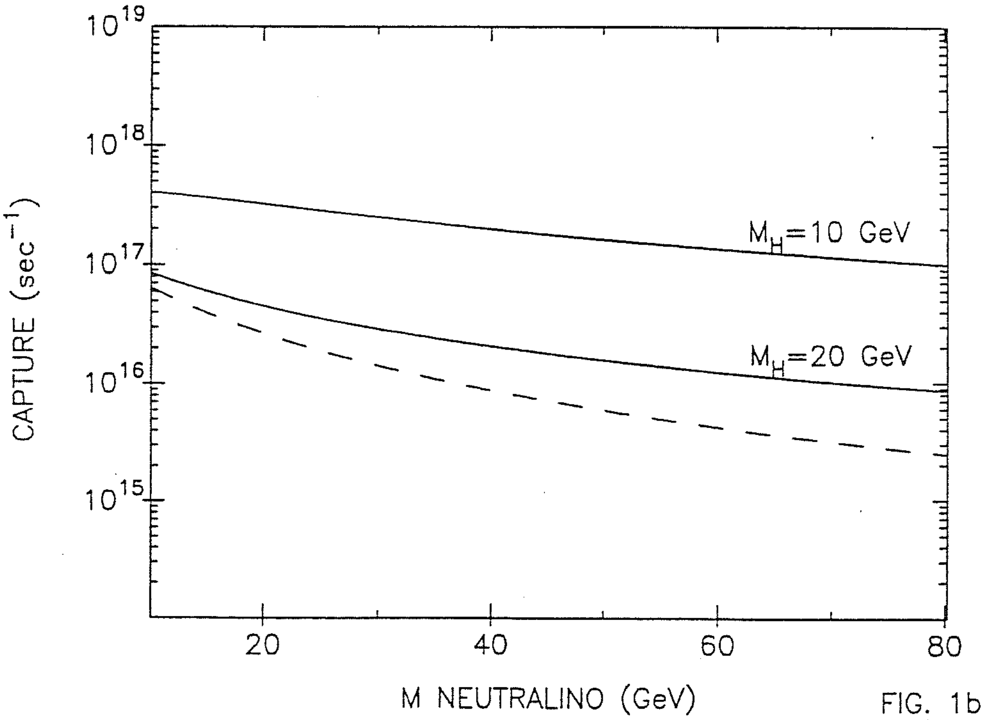

FIG. 1b

Fig.1. Neutralino capture rates by the earth (a) and the sun (b) versus m_χ, for different values of the Higgs mass, taking $\alpha_H = 1$ and $\alpha_Z = 0.1$. The dashed line shows the case were the Higgs effect has been neglected. The solar capture has been scaled by $(R_\oplus/1A.U.)^2 = 1.8 \cdot 10^{-9}$.

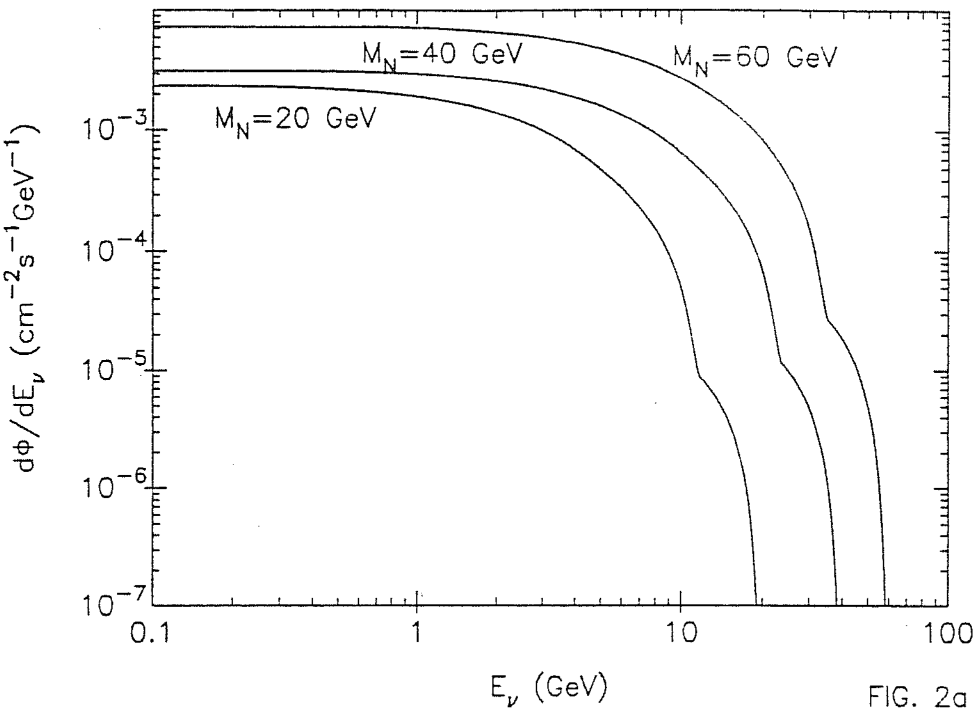

FIG. 2a

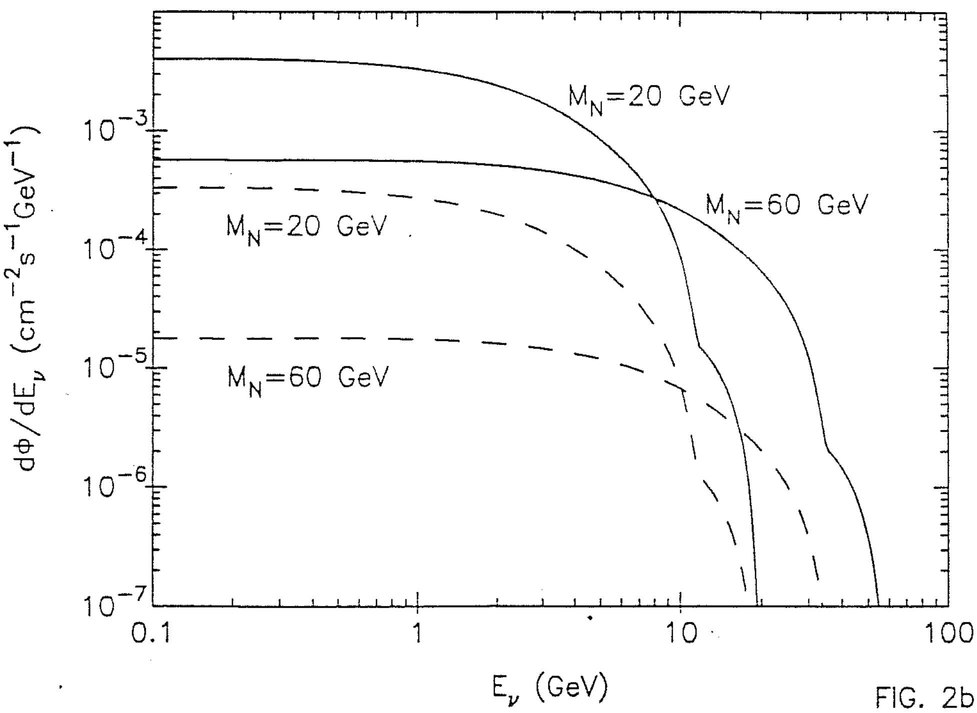

FIG. 2b

Fig.2. The electron (or muon) neutrino and antineutrino energy spectrum from the earth (a) and the sun (b), for different choices of the neutralino mass. We have taken $\alpha_H = 1$ and $\alpha_Z = 0.1$. The solid lines refer to $m_H = 10$ GeV while the dashed lines show the case where the Higgs effect has been neglected.

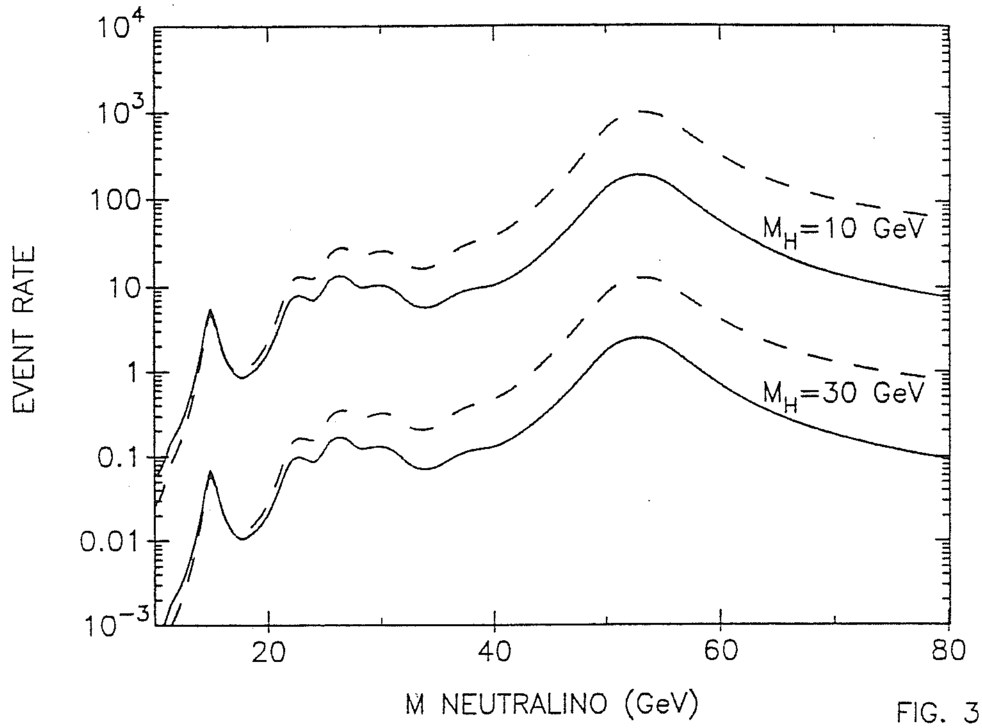

FIG. 3a

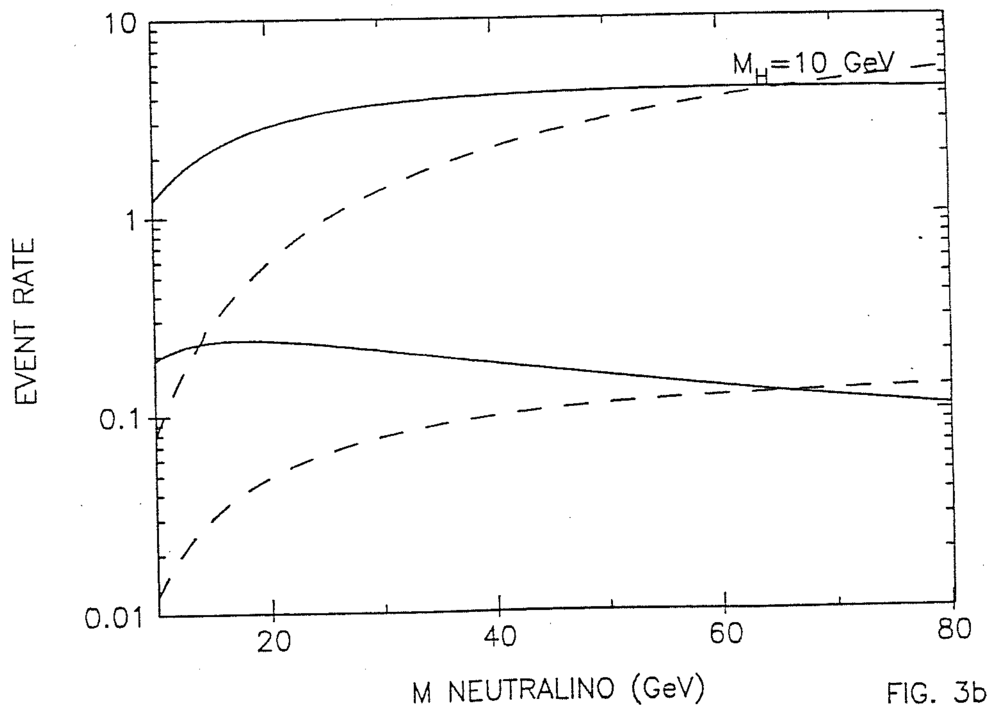

FIG. 3b

Fig.3. Solid lines: rate of events per kiloton year for electrons (or muons) produced within the detector (R_e). Dashed lines: rate of events per 100 m^2 year for muons entering the detector and either going through or being stopped (R_μ). We show the result for the earth (a) and the sun (b) with $m_H = 10$ GeV and with no Higgs contribution. For the sun, R_μ is scaled by the factor 1/3. We have taken $\alpha_H = 1$ and $\alpha_Z = 0.1$.

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CHAPTER 5

A Supersymmetric Solution to the Solar Neutrino and Dark Matter Problems

As it has been pointed in chapter 1, both the solar neutrino and dark matter problems could be simultaneously solved by the cosmion, a Weakly Interacting Massive Particle fulfilling several restrictive constraints [5.1]. First of all, the cosmion should be a (quasi)–stable particle with a sizable cosmological relic density in order to account for the missing mass. An efficient solar capture and energy transport requires the effective cosmion scattering cross section in the sun (σ_s) to be about:

$$\sigma_s \simeq 4.10^{-36} cm^2 \tag{5.1}$$

Moreover, the cosmion annihilation cross section in the sun (σ_A) must be suppressed according to

$$< \sigma_A v > \lesssim 10^{-4} \sigma_s \simeq 4.10^{-40} cm^2 \tag{5.2}$$

where v is the relative velocity of the annihilating particles. Finally, the cosmion should have a mass in the range of 4–10 GeV, so that it does not evaporate from the sun and leads to an efficient thermal transport.

The elementary particle realization of this suggestive astrophysical scenario is not straightforward. The standard neutrino, photino, higgsino and sneutrino do not satisfy the above requirements[5.2]. In particular, $\sigma_s \simeq 4.10^{-36} cm^2$ is difficult to be achieved, being about two orders of magnitude larger than a typical weak cross section. Moreover, for self charge–conjugate states, the inequality (5.2) seems problematic. Generally, a process contributing to σ_s leads also, in the crossed channel, to a contribution to σ_A of the same order of magnitude. Cosmions with a conserved quantum number and cosmic asymmetry between cosmions and anti–cosmions do not suffer from this problem. In the cosmion models proposed so far [5.3-6], it is in general this last idea that is employed to suppress annihilations without fulfilling (5.2).

In this chapter, it will be shown that the neutralino, in the presence of a light ($\simeq 1 - 2$ GeV) Higgs boson, can behave as a cosmion in spite of its Majorana nature [5.7]. A Higgs with mass of 1-2 GeV lies within the still experimentally allowed window discussed in ref. [5.8]. The light Higgs exchange provides the large neutralino capture cross section in the sun, eq. (5.1), as was discussed in **1.3**. The annihilation channels due to the Higgs coupling turn out to be p–wave suppressed. Because of the low neutralino velocity in the sun, this naturally accounts for eq. (5.2). With an appropriate choice of parameters, the neutralino annihilation due to the other interactions present in the theory can also be suppressed, and eq. (5.2) is then satisfied. It will be shown how the whole scheme can be consistently included in the minimal supersymmetric standard model (for a review, see ref. [5.9]). Although the attention will be focused in this model, the couplings will be parameterized in such a way that the discussion can be extended to other versions of supersymmetric models.

In the minimal supersymmetric model, one finds four neutralinos, which are the mass eigenstates, mixtures of b–ino, neutral w–ino and two higgsinos. Due to R–parity conservation, the lightest neutralino (χ) is a possible dark matter candidate. Its mass and its combination in terms of the four weak eigenstates

$$\chi = \gamma_1 \tilde{b} + \gamma_2 \tilde{w} + \gamma_3 \tilde{h}_1 + \gamma_4 \tilde{h}_2 \ , \tag{5.3}$$

are functions of the following parameters of the theory: (i) the $SU(2)$ gaugino mass M[1]; (ii) the Higgs mixing mass μ; (iii) the ratio $v_2/v_1 = \tan\beta$ of the vacuum expectation values of the two neutral Higgs bosons[2]. See Chapter 1 for more details.

Let us first parameterize the scalar Higgs interaction with fermionic matter (f) and the neutralino (χ) as follows:

$$\mathcal{L}_H = \frac{g}{2}\, Q_H H^0 \bar{\chi}\chi - \frac{g}{2} K_f\, \frac{m_f}{M_w}\, h^0 \bar{f} f \tag{5.4}$$

In the context of the minimal model, we find

$$K_f = \begin{cases} \cos\alpha/\sin\beta & \text{for } f = \text{up-type} \\ -\sin\alpha/\cos\beta & \text{for } f = \text{down-type} \end{cases} \tag{5.5}$$

$$Q_H = \gamma_3 \sin\alpha + \gamma_4 \cos\alpha \tag{5.6}$$

$$\cos 2\alpha = \cos 2\beta \left(\frac{\cos^2\beta + r^2 - 2r}{\cos^2 2\beta + r^2 - 2r\cos^2 2\beta} \right) \qquad r \equiv \frac{m^2_{H^0}}{m^2_z} \tag{5.7}$$

The scalar Higgs exchange[5.10] mediates a coherent interaction between the neutralino and matter[5.11], according to the graph of fig. 1.

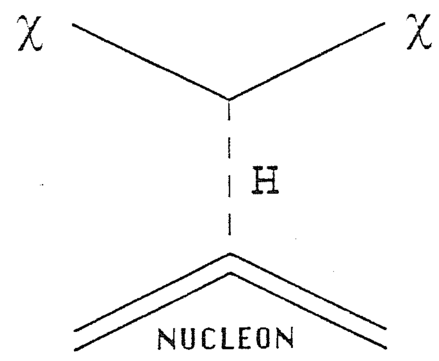

Fig. 1: Higgs mediated neutralino interaction with nuclei

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[1] We assume the usual unification relation for the gaugino masses.

[2] v_2 is proportional to the up quark mass, while v_1 is proportional to the down quark mass

$$< \sigma_A (\chi\chi - H^0 H^0) \, v >= \frac{g^4}{24\pi} \frac{Q_H^4}{m_\chi^2} \frac{(9 - 8x + 2x^2)(1-x)^{1/2}}{(2-x)^4} v^2 \quad (5.12)$$

where $x \equiv \frac{m_{H^0}^2}{m_\chi^2}$. The v^2 factor in eqs. (5.11)–(5.12) is the signal of the above–mentioned p–wave suppression.

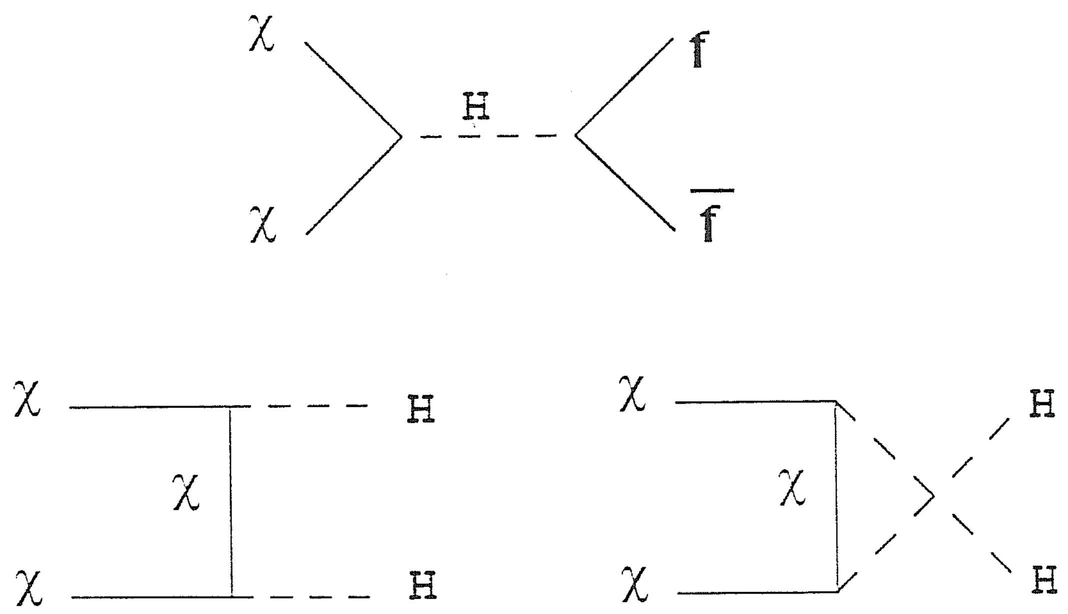

Fig. 2: Neutralino annihilation channels involving interaction (5.4)

The annihilation rates (5.11) and (5.12) have to satisfy eq. (5.2). However, a more stringent constraint comes from the requirement of a correct relic abundance of neutralinos as dark matter candidates. Their contribution to the present energy density of the universe if approximately given by

$$\Omega_\chi h^2 \simeq \frac{2.10^{-37} cm^2}{< \sigma_A v >_f}, \quad (5.13)$$

where h is the Hubble constant in units of 100 km sec^{-1} Mpc^{-1}. The annihilation rate in eq. (5.13) is evaluated at the neutralino freeze–out temperature.

Now, since the neutralino mass density should exceed the baryon cosmic density, without overclosing the universe, it results

$$0.025 \lesssim \Omega_\chi h^2 \lesssim 1. \tag{5.14}$$

Because of the p–wave suppression, it is possible to simply rescale the annihilation rate and, from eq. (5.13), obtain that (5.14) implies

$$2.10^{-43} cm^2 \lesssim < \sigma_A v > v_f^2 \lesssim 9.10^{-42} cm^2, \tag{5.15}$$

where $< \sigma_A v >$ has to be evaluated in the sun ($v \simeq 10^{-3}$) and v_f is the neutralino freeze–out velocity.

Note that if eq. (5.15) holds, then eq. (5.2) is automatically satisfied. Using the expressions (5.11–12) for $< \sigma_A v >$, eq. (5.15) roughly corresponds to the bound:

$$10^{-3} \lesssim Q_H^2 \lesssim 10^{-1}, \tag{5.16}$$

Therefore, it has been shown that if the Higgs–neutralino coupling constant Q_H lies in the range (5.16), then (i) a Higgs scalar with mass $m_{H^0} \simeq 1 - 2$ GeV mediates the correct solar capture cross section, eq. (5.1); (ii) the neutralino annihilation in the sun due to the presence of the light Higgs is safely suppressed; (iii) the neutralino has the correct relic density to account for dark matter.

Obviously, the neutralino feels all the interactions of the full supersymmetric theory, besides those in eq. (5.4). Therefore, it will be discussed now how the suppression of the annihilation rate, eq. (5.2), can be achieved in the complete model.

Squarks and sleptons can mediate the process $\chi\chi \to f\bar{f}$ through t–channel exchange. If they are heavy enough, the annihilation is suppressed, according to eq. (5.2). A limit on their masses can be derived assuming that all squarks and sleptons are degenerate and taking the neutralino as a pure photino state[4]. This is a good approximation, since the couplings of the higgsino components are suppressed by factors

[4] For more general formulae, see ref. [5.11].

of $\frac{m_f}{m_w}$. In the non–relativistic limit, the exchange of s fermions with mass $\tilde{m}$ gives

$$< \sigma_A \left(\tilde{\gamma}\tilde{\gamma} \rightarrow f\bar{f}\right) v >= \frac{8\pi\alpha^2}{\tilde{m}^4} \sum_f c_f\, Q_f^4\, m_f^2 \sqrt{1-\frac{m_f^2}{m_\chi^2}} \tag{5.17}$$

Q_f and c_f are respectively the electric charge and the color factor for the fermions allowed by phase space. The annihilation rate (5.17) satisfies the constraint (5.2), if:

$$\tilde{m} \gtrsim 280\ GeV. \tag{5.18}$$

This is a fairly reasonable value for s fermion masses in standard supergravity models.

Z^0–exchange in the s–channel can also contribute to $\chi\chi \rightarrow f\bar{f}$. In the non–relativistic limit, the annihilation rate is:

$$< \sigma_A \left(\chi\chi \rightarrow f\bar{f}\right) v >= \frac{G_F^2}{4\pi}\, Q_z^2 \sum_f c_f\, m_f^2 \sqrt{1-\frac{m_f^2}{m_\chi^2}} \tag{5.19}$$

Q_Z is the coupling constant of the $Z^0\chi\chi$ interaction term:

$$\frac{g}{4\cos\theta_w}\, Q_z\, Z_\mu\, \bar{\chi}\, \gamma^\mu\, \gamma_5\, \chi \tag{5.20}$$

and, in the minimal model, is given by

$$Q_z = \gamma_3^2 - \gamma_4^2, \tag{5.21}$$

where the parameters γ_i have been defined in eq. (5.3).

Eq. (5.19) shows the well–known p–wave suppression for the annihilation into massless fermions, due to the Majorana nature of neutralinos[5.12]. However, this suppression is, for heavy fermions, not sufficient to accomplish eq. (5.2). It is then necessary to require

$$Q_z^2 \lesssim 10^{-3} \tag{5.22}$$

Even in the minimal supersymmetric model, the proposed scheme is consistent. In other words, it will be shown that the existence of a neutralino with mass 4–10 GeV and conditions (5.16) and (5.22) can be simultaneously fulfilled. This is true in two different situations.

1^0 Solution: The neutralino is a "quasi–photino" state. Then, the higgsino components γ_3, γ_4 are small and Q_z is suppressed, see eq. (5.21). Also Q_H goes to zero as the higgsino components vanish, but Q_H is linear in γ_3, γ_4 (see eq. (5.5)) while Q_z is quadratic. This allows to satisfy simultaneously eqs. (5.16) and (5.22). In general, large values of v_2/v_1 are needed, in agreement with the expectation of the radiative electroweak symmetry breaking in supergravity models with a heavy top quark. Since the right hand side of eq. (5.9) grows about quadratically with v_2/v_1, even if Q_H is small, α_H is of order one, and $m_H = 1 \div 2$ GeV is expected from eq. (5.10). Under these conditions, the main annihilation channel for primordial neutralinos is $\chi\chi \to b\bar{b}$ through Higgs exchange, which correctly leads to $\Omega_\chi \simeq 0.1 \div 1$. The "quasi–photino" solution is obtained when the mass parameter M is small with respect to μ and M_z, which automatically leads to a light neutralino. If the usual unification relation for gaugino masses is assumed, the interesting feature that gluinos lie just behind their present experimental limit results.

2^0 Solution: The neutralino parameters are close to the condition

$$\mu = \left(\frac{\cos^2 \theta_w}{M} + \frac{\sin^2 \theta_w}{M'} \right) \sin 2\beta M_z^2 \tag{5.23}$$

where M' is the $U(1)$ gaugino mass. When this relation holds, the neutralino mass matrix becomes singular [5.14]. In the vicinity of equality (5.23), one neutralino state is light. Now, the higgsino components are not necessarily small, but they are related by

$$\gamma_3 = \tan\beta\gamma_4 \tag{5.24}$$

If $v_2/v_1 \simeq 1$, eq. (5.24) entails the suppression of Q_z. Q_H is not suppressed and turns out to be much larger than in the previous case. Therefore, $\chi\chi \to H^0 H^0$ is the main annihilation channel of primordial neutralinos. This rate can be rather large leaving the neutralino relic abundance near the lower end of the allowed region (eq. (5.14)). Again, α_H is of order one and we obtain $m_H \simeq 1 \div 2$ GeV.

Both solutions can possibly lead to an uncomfortable light chargino (the charged weakly interacting supersymmetric particle). However, if the neutralino is sufficiently close in mass to the chargino, the light chargino is not experimentally ruled out, since its visible signals are very soft.

It has been checked numerically that both previous solutions provide a consistent picture of the neutralino as a cosmion in the minimal supersymmetric model. In Table I, a sample of neutralino parameters leading to a correct scenario are shown. The values in the table reflect the main features of the two different solutions just illustrated.

Solution	M	μ	v_2/v_1	M_χ	Q^2_H	Q^2_z	α_H
1^0	10	-150	20	5	3.10^{-3}	9.10^{-4}	1.2
	15	-120	10	8	2.10^{-3}	8.10^{-4}	0.2
2^0	80	150	1.1	4	2.10^{-1}	5.10^{-4}	1.8
	105	130	1.1	10	2.10^{-1}	7.10^{-4}	2.0

Table I: *Values of the neutralino parameters $M, \mu, v_2/v_1$ which lead to a correct cosmion solution.*

Turning now to the consideration of light pseudoscalars, these particles can be important since they have the right CP quantum number for contributing to the annihilation of neutralinos in s–wave. The relevant processes are $\chi\chi \to f\bar{f}$ through pseudoscalar exchange or scalar–pseudoscalar pair production via Z^0 or neutralino exchange.

The minimal supersymmetric model contains one pseudoscalar. Since its coupling with the neutralino is deeply connected to the scalar–neutralino coupling, the only way to get rid of the pseudoscalar contribution is by taking it sufficiently heavy. From $\chi\chi \to f\bar{f}$ mediated by the pseudoscalar, eq. (5.2) implies that its mass should be larger than about $200 \div 300$ GeV.

In the case of Solution 1, the large v_2/v_1 and the small m_H necessarily imply that the pseudoscalar is approximately degenerate with the scalar and therefore very light. If one insists in large values of v_2/v_1, one has to enlarge the Higgs sector (extra gauge singlet, more Higgs doublets,...) in order to avoid the relation between scalar and pseudoscalar masses. On the contrary, the scenario envisaged in solution 2 predicts exactly the correct Higgs mass spectrum. If $v_2/v_1 \simeq 1$, one Higgs scalar is forced to be very light (few GeV), while the pseudoscalar gets very heavy (hundreds GeV).

The last remark concerns the limits on light Higgs bosons. Theoretically, there is no lower bound on the Higgs mass in a two doublet model. As discussed in ref. [5.8], due to theoretical uncertainties, experimental data on B decays are not able to rule out a Higgs of 1–2 GeV. It is also premature to exclude such a Higgs from the unsuccessful search for $\Upsilon \to H^0\gamma$ [5.15], since higher order QCD corrections can easily affect the result. However, if $v_2/v_1 \gg 1$, the light Higgs has an enhanced coupling with down-type quarks and thus it can certainly be excluded by the CUSB experiment [5.15]. Therefore, Solution 1 has to be discarded, while Solution 2 is fully consistent.

In conclusion, a scenario in which the supersymmetric neutralino is a good dark matter candidate and simultaneously solves the solar neutrino problem has been described. The scheme works as follows.

If the neutralinos are sufficiently concentrated in the sun, they affect the solar core temperature and reduce the flux of the observed ^{8}B neutrinos. In order to do so, the neutralino should have a mass of 4–10 GeV, large cross section and small annihilation rate in the sun.

The exchange of a Higgs scalar with mass 1–2 GeV yields a capture cross section in the correct range. Such a light Higgs has not yet been experimentally ruled out[5.8]. The contribution to the annihilation from the same Higgs interaction vanishes in s–wave and thus it is strongly suppressed.

This p–wave suppression is less effective for primordial neutralinos which are much hotter than the neutralinos captured by the sun. Then, the annihilation from Higgs interaction leads to a neutralino relic density which correctly accounts for dark matter.

The contributions to the annihilation coming from forces mediated by squarks, sleptons and Higgs pseudoscalars are small enough if these particles are heavier than few hundreds GeV. The Z^0–exchange is suppressed by inhibiting the $Z^0\chi\chi$ coupling.

This scheme can be implemented in the minimal supersymmetric standard model if the neutralino mass parameters approximately satisfy eq. (5.23) with $v_2/v_1 \simeq 1$. Then, the neutralino is light and has suppressed interactions with the Z^0. The appropriate Higgs mass spectrum is also automatically predicted.

Experiments looking for dark matter, light Higgs bosons and supersymmetric particles can all very well test the validity of the scenario proposed here.

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