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Probing the Desert with Ultra–Energetic Neutrinos from the Sun and the Earth

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Abstract

Realistic superstring models generically give rise to exotic matter states, which arise due to the "Wilson–line" breaking of the non–Abelian unifying gauge symmetry. Often such states are protected by a gauge or local discrete symmetry and therefore may be stable or meta–stable. We study the possibility of a flux of high energy neutrinos coming from the sun and the earth due to the annihilation of such exotic string states. We also discuss the expected flux for other heavy stable particles – like the gluino LSP. We comment that the detection of ultra–energetic neutrinos from the sun and the earth imposes model independent constraints on the high energy cutoff, as for example in the recently entertained TeV scale Kaluza–Klein theories. We therefore propose that improved experimental resolution of the energy of the muons in neutrino detectors together with their correlation with neutrinos from the sun and the center of the earth will serve as a probe of the desert in Gravity Unified Theories.

1 Introduction

The spectacular confirmation of the Standard Model at LEP/SLC and other high energy colliders [1], as well as the proton longevity and suppression of left–handed neutrino masses, provide strong support for the grand desert scenario and unification. Indeed, the question of how to test this hypothesis in present and future experiments constitutes much of today's activity in high energy physics. Direct signatures at LHC/NLC will be able to probe the desert only up to a few TeV. On the other hand, it has been suggested that all present day experimental data including ultra–high energy cosmic rays above the Greisen–Zatsepin–Kuzmin cutoff, cannot rule out the possibility that the fundamental scale is near the TeV scale [2].

In this paper we examine the potential of probing the desert using ultra energetic neutrinos from the sun and the earth. Such ultra–energetic neutrinos arise if the dark matter in the galactic halo is composed of heavy long–lived matter, that is trapped in the sun and the earth and subsequently annihilates into neutrinos [3]. The magnitude of the expected signal depends on the competition between the trapping and annihilation rates of specific dark matter candidates, which in turn depends on the specific properties of such states. The well known case which has been studied extensively in the past is the trapping and annihilation of stable neutralinos which arise in supersymmetric theories with conserved R–parity [3, 4, 5].

High energy ($E > 1$ GeV) neutrinos from the sun and the earth can only be produced by the decay or annihilation of a heavy particle. For the simple reason that, unlike charged matter, neutrinos cannot be accelerated, the observation of ultra–energetic neutrinos whose production can be definitely correlated with the sun or the center of the earth, will signal new physics at the high energy scale. In addition, such an observation would also impose a model independent constraint on the fundamental cutoff scale.

Superstring models generically give rise to exotic matter states which arise because of the breaking of the non–Abelian unifying gauge symmetry, G, by Wilson–lines [6, 7, 8]. In many respects the unifying gauge symmetry is similar to the gauge group of four dimensional grand unification and the Wilson lines are similar to the Higgs bosons in the adjoint representation. However, there are some notable differences. The eigenvalues of the Wilson lines are quantized while the eigenvalues of the Higgs in the adjoint representation are continuous. Another important difference is that the breaking of the gauge symmetries by Wilson lines results in massless states that do not fit into multiplets of the original unbroken gauge symmetry. We refer to such states generically as exotic "Wilsonian" matter states. This is an important property as it may result in conserved quantum numbers that will indicate the stability of these massless "Wilsonian" states. The simplest example of this phenomenon is the existence of states with fractional electric charge in the massless spectrum of superstring models [7, 9, 10, 11]. Such states are stable due to electric charge conservation. As there exist strong constraints on their masses and abundance, states with fractional electric charge must be diluted away or extremely massive. The same "Wilson line" breaking mechanism, which produces matter with fractional electric charge, is also responsible for the existence of states which carry the "standard" charges under the Standard Model gauge group but which carry fractional charges under a different subgroup of the unifying gauge group. For example, if the group G is $SO(10)$ then the "Wilsonian" states may carry non–standard charges under the $U(1)_{Z'}$ symmetry, which is embedded in $SO(10)$ and is orthogonal to $U(1)_Y$. Such states can therefore be long–lived if the $U(1)_{Z'}$ gauge symmetry remains unbroken down to low energies, or if some residual local discrete symmetry is left unbroken after $U(1)_{Z'}$ symmetry breaking. The existence of heavy stable "Wilsonian" matter can be argued to be a "smoking gun" of string unification.

There are several examples of "Wilsonian matter" in the literature. The uniton is a dark matter candidate which arises in realistic heterotic string models [8]. We show that the limits due to energetic neutrinos from the sun and the earth are far more restrictive than the constraints arising from over-closure of the universe in the case of the uniton due to its strong interaction with matter. As another example of a strongly interacting stable particle, we also consider the constraints on recent proposals of gluino–LSP [12] (which has not been suggested as a dark matter candidate), and show that the constraints imposed from annihilation into energetic neutrinos in the sun are sufficiently restrictive to rule out this possibility. Other astrophysical constraints applicable to the gluino LSP were considered in [13]. Other candidates we consider are the *cryptons* of ref. [14] which are fractionally charged states confined by a hidden non-Abelian gauge symmetry, and other "Wilsonian matter" states which may consist of free fractionally charged states or Standard Model neutral states with fractional charges under the $U(1)_{Z'}$ symmetry. We discuss the potential for discovery of each of these states by energetic neutrinos from the sun.

Recently there has been considerable interest in the possibility that a very massive, long-lived particle can account for both dark matter and the observed ultra-high energy cosmic rays (above the GZK cutoff [15]) through its decays [16]. For example, a particle with a mass of 10^{13-14} GeV with a lifetime (depending on the relic density) greater than the age of the Universe could possibly explain these phenomena, although the more detailed computation in [17] indicates that the particle mass should be about 10^{11-12} GeV. The specific case of cryptons in this context was considered in [18, 17]. Though one would not expect the lifetime for such a massive particle to be sufficiently long, discrete symmetries have been proposed to achieve the required stability [19]. The light gluino candidates of [12] are too few in number to account for the dark matter in the galaxy, nevertheless, they have been proposed as a possible source for the ultra-high energy cosmic rays as do the very light gluinos discussed in [20]. The latter can not be tested by annihilation in the sun or the earth as they are too light and would produce neutrinos below threshold.

In what follows, we first describe the types of Wilsonian dark matter states that can be expected to arise in realistic string theory formulations and describe the specific candidates we consider. We then go on to compute the flux of neutrinos generated in the core of the sun and the earth and estimate the flux of upward-going muons which signal the annihilation into neutrinos. We find that that certain candidates such as the uniton which have strong interactions can be excluded as dark matter. As a source of the ultra-high energy cosmic-rays, the gluino–LSP candidate can also be excluded. For other candidates, such the crypton, the best one can do is place limits on its mass and interaction strength.

2 Realistic free fermionic models

The realistic models in the free fermionic formulation are generated by a basis of boundary condition vectors for all world–sheet fermions [21]-[25]. The basis is constructed in two stages. The first stage consists of the NAHE set [21, 25], which is a set of five boundary condition basis vectors, $\{\mathbf{1}, S, b_1, b_2, b_3\}$. The gauge group after the NAHE set is $SO(10) \times SO(6)^3 \times E_8$ with $N = 1$ space–time supersymmetry. The NAHE set correspond to $Z_2 \times Z_2$ orbifold compactification [26]. The Neveu–Schwarz sector corresponds to the untwisted sector, and the sectors b_1, b_2 and b_3 to the three twisted sectors of the $Z_2 \times Z_2$ orbifold model. In addition to the gravity and gauge multiplets, the Neveu–Schwarz sector produces six multiplets in the 10 representation of $SO(10)$, and several $SO(10)$ singlets transforming under the flavor $SO(6)^3$ symmetries. The sectors b_1, b_2 and b_3 produce 48 spinorial 16's of $SO(10)$, sixteen each from the sectors b_1, b_2 and b_3. All the states at this level of the string model building are in GUT representations.

The second stage of the basis construction consists of adding three additional basis vectors to the NAHE set. These three additional basis vectors correspond to "Wilson lines" in the orbifold formulation. Three additional vectors are needed to reduce the number of generations to three, one from each sector b_1, b_2 and b_3. At the same time the additional boundary condition basis vectors break the gauge symmetries of the NAHE set. The $SO(10)$ symmetry is broken to one of its subgroups $SU(5) \times U(1)$, $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)_{B-L} \times U(1)_{T_{3_R}}$. To break the $SO(10)$ symmetry to $SU(3) \times SU(2) \times U(1)_C \times U(1)_L$* we first use a boundary condition basis vector which breaks $SO(10)$ to $SU(5) \times U(1)$ or $SO(6) \times SO(4)$. We then break the gauge group to $SU(3) \times SU(2) \times U(1)_C \times U(1)_L$. Since the superstring derived standard–like models contain the $SO(6) \times SO(4)$, as well as the $SU(5) \times U(1)$, breaking sectors, their massless spectra admits also the exotic representations that can appear in these models.

In the superstring standard–like models, the observable gauge group after application of the generalized GSO projections is $SU(3)_C \times U(1)_C \times SU(2)_L \times U(1)_L \times U(1)^3 \times U(1)^n$. The electromagnetic charge is given by

$$U(1)_{\mathrm{e.m.}} = T_{3_L} + U(1)_Y, \tag{2.1}$$

*$U(1)_C = \frac{3}{2}U(1)_{B-L}; U(1)_L = 2U(1)_{T_{3_R}}$.

where T_{3_L} is the diagonal generator of $SU(2)_L$, and $U(1)_Y$ is the weak hypercharge. The weak hypercharge is given by[†]

$$U(1)_Y = \frac{1}{3}U(1)_C + \frac{1}{2}U(1)_L \tag{2.2}$$

and the orthogonal combination is given by

$$U(1)_{Z'} = U(1)_C - U(1)_L. \tag{2.3}$$

The hidden E_8 gauge group is typically broken to $SU(5)_H \times SU(3)_H \times U(1)^2$, and the flavor $SO(6)^3$ symmetries are broken to factors of $U(1)$s.

The massless spectrum of the standard–like models contains three chiral generations from the sectors which are charged under the horizontal symmetries. Each of these consists of a 16 of $SO(10)$, decomposed under the final $SO(10)$ subgroup as

$$
\begin{aligned}
e_L^c &\equiv [(1, \ \frac{3}{2}); (1, \ 1)]_{(\ 1 \ , \ 1/2 \ , \ 1)} \ ; \ u_L^c \equiv [(\bar{3}, -\frac{1}{2}); (1, -1)]_{(-2/3,1/2,-2/3)}; && (2.4)\\
d_L^c &\equiv [(\bar{3}, -\frac{1}{2}); (1, \ 1)]_{(1/3,-3/2,1/3)} \ ; \ Q \equiv [(3, \frac{1}{2}); (2, \ 0)]_{(1/6,1/2,(2/3,-1/3))}; && (2.5)\\
N_L^c &\equiv [(1, \ \frac{3}{2}); (1, -1)]_{(\ 0 \ , \ 5/2 \ , \ 0)} \ ; \ L \equiv [(1, -\frac{3}{2}); (2, \ 0)]_{(-1/2,-3/2,(0,1))}, && (2.6)
\end{aligned}
$$

where we have used the notation

$$[(SU(3)_C \times U(1)_C); (SU(2)_L \times U(1)_L)]_{(Q_Y, Q_{Z'}, Q_{\rm e.m.})}, \tag{2.7}$$

and have written the electric charge of the two components for the doublets.

The matter states from the NS sector and the sectors b_1, b_2 and b_3 transform only under the observable gauge group. In the realistic free fermionic models, there is typically one additional sector that produces matter states transforming only under the observable gauge group [25]. These states complete the representations that we identify with possible representations of the Standard Model. In addition to the Standard Model states, semi–realistic superstring models may contain additional multiplets, in the 16 and $\overline{16}$ representation of $SO(10)$, in the vectorial 10 representation of $SO(10)$, or the 27 and $\overline{27}$ of E_6. Such states can pair up to form super-massive states. They can mix with, and decay into, the Standard Model representation unless some additional symmetry, which forbids their decay, is imposed. For example, in the flipped $SU(5)$ superstring models [21], two of the additional vectors which extend the NAHE set produce an additional 16 and $\overline{16}$ representation of $SO(10)$. These states are used in the flipped $SU(5)$ model to break the $SU(5) \times U(1)$ symmetry to $SU(3) \times SU(2) \times U(1)$.

[†] Note that we could have instead defined the weak hypercharge to be $U(1)_Y = \frac{1}{3}U(1)_C - \frac{1}{2}U(1)_L$. This amounts to the same redefinition of fields between the straight and flipped $SU(5)$. In this paper we will use the definition in Eq. 2.2.

In addition to the states mentioned above transforming solely under the observable gauge group, the sectors $b_j + 2\gamma$ produce matter states that fall into the 16 representation of the hidden $SO(16)$ gauge group decomposed under the final hidden gauge group. The states from the sectors $b_j + 2\gamma$ are $SO(10)$ singlets, but are charged under the flavor $U(1)$ symmetries. All the states above fit into standard representation of the grand unified group which may be, for example, $SO(10)$ or E_6, or are singlets of these groups. They carry the standard charges under the Standard Model gauge group or of its GUT extensions. The superstring models, however, contain additional states that cannot fit into multiplets of the original unifying gauge group. Such states are exotic stringy states and cannot fit into representations of the underlying $SO(10)$ symmetry group of the NAHE set. They result from the breaking of the $SO(10)$ gauge group at the string level via the boundary condition assignment. In the next section we enumerate the states that appear in free fermionic models.

3 Exotic matter

The boundary condition basis vectors, beyond the NAHE set, are used to break the unifying $SO(10)$ gauge group. Since these sectors correspond to "Wilson lines" they give rise to massless states that do not fit into representations of the original $SO(10)$ symmetry. As a result the massless spectrum contains states with fractional charges under the unbroken $U(1)$ generators of the original $SO(10)$ gauge group. This is a new and common feature of superstring models. In many examples the exotic states appear in vector–like representations and can acquire a heavy mass. The "Wilsonian" matter phenomenon is an important feature as it may result in discrete symmetries that may prevent the decay of the exotic massive states into the Standard Model states, and therefore give rise to dark matter candidates. The following exotic matter representations can appear in free fermionic level one models. The states are classified according to the unbroken $SO(10)$ subgroup in each sector [8, 25].

From the $\underline{SO(6) \times SO(4)}$ type sectors we obtain the following exotic states.

- Color triplets : $[(3, \frac{1}{2}); (1, 0)]_{(1/6,1/2,1/6)} \quad ; \quad [(\bar{3}, -\frac{1}{2}); (1, 0)]_{(-1/6,-1/2,-1/6)}$
- Electroweak doublets : $[(1, 0); (2, 0)]_{(0,0,\pm 1/2)}$
- Fractionally charged $SU(3)_C \times SU(2)_L$ singlets :

$$[(1, 0); (1, \pm 1)]_{(\pm 1/2,\mp 1/2,\pm 1/2)} \quad ; \quad [(1, \pm 3/2); (1, 0)]_{(\pm 1/2,\pm 1/2,\pm 1/2)} \qquad (3.1)$$

The color triplets bind with light quarks to form mesons and baryons with fractional electric charges $\pm 1/2$ and $\pm 3/2$. The $\underline{SO(6) \times SO(4)}$ type states can appear in the Pati–Salam type models [23] or in Standard–like models.

From sectors which break the $SO(10)$ symmetry into $\underline{SU(5) \times U(1)}$ we obtain exotic states with fractional electric charge $\pm 1/2$

- Fractionally charged $SU(3)_C \times SU(2)_L$ singlets :

$$[(1, \pm 3/4); (1, \pm 1/2)]_{(\pm 1/2, \pm 1/4, \pm 1/2)} \tag{3.2}$$

In general the fractionally charged states may transform under a non–Abelian hidden gauge group in which case the fractionally charged states may be confined. For example, in the "revamped" flipped $SU(5)$ model [21] the states with fractional charge $\pm 1/2$ transform as 4 and $\bar{4}$ of the hidden $SU(4)$ gauge group. The states with the charges in eq. (3.2) are called the "cryptons" and may form good dark matter candidates [9] if the lightest confined state is electrically neutral. In the "revamped" flipped $SU(5)$ model it has been argued that the lightest state is the "tetron", which contains four fundamental constituents. In other models, states with the charges of eq. (3.2) may be singlets of all the non–Abelian group factors.

Finally in the superstring derived standard–like models we may obtain exotic states from sectors which are combinations of the $\underline{SO(6) \times SO(4)}$ breaking vectors and $\underline{SU(5) \times U(1)}$ breaking vectors. These states therefore arise only in the $\underline{SU(3) \times SU(2) \times U(1)^2}$ type models. These states then carry the standard charges under the Standard Model gauge group but carry fractional charges under the $U(1)_{Z'}$ gauge group. The following exotic states are obtained:

- color triplets :

$$[(3, \frac{1}{4}); (1, \frac{1}{2})]_{(-1/3, -1/4, -1/3)} \quad ; \quad [(\bar{3}, -\frac{1}{4}); (1, \frac{1}{2})]_{(1/3, 1/4, 1/3)} \tag{3.3}$$

Due to its potential role in string gauge coupling unification [27], this state is referred to as "the uniton" [8]. The uniton forms bound states with the lightest up and down quarks and gives rise to ultra–heavy mesons. In ref. [8] it has been shown that the lightest meson can be the electrically neutral state.

- electroweak doublets : $[(1, \pm\frac{3}{4}); (2, \pm\frac{1}{2})]_{(\pm 1/2, \pm 1/4, (1,0);(0,-1))}$

Unlike the previous electroweak doublets, these electroweak doublets carry the regular charges under the standard model gauge group but carry "fractional" charge under the $U(1)_{Z'}$ symmetry. Finally, in the superstring derived standard–like models we also obtain states which are Standard Model singlets but carry "fractional" charges under the $U(1)_{Z'}$ symmetry.

- Standard model singlets with "fractional" $U(1)_{Z'}$ charge :

$$[(1, \pm\frac{3}{4}); (1, \mp\frac{1}{2})]_{(0, \pm 5/4, 0)} \tag{3.4}$$

These states may transform under a non–Abelian hidden gauge group or may be singlets of all the non–Abelian group factors. This type of Standard Model singlet appears in all the known free fermionic standard–like models.

There are several important issues to examine with regard to the exotic states. Since some of these states carry fractional charges, it is desirable to make them sufficiently heavy or sufficiently rare. A priori, in a generic string model, it is not at all guaranteed that the states with fractional electric charge can be decoupled or confined [28]. Therefore, their presence imposes an highly non–trivial constraint on potentially viable string vacua. In the NAHE–based free fermionic models, all the exotic matter states appear in vector–like representations. They can therefore obtain mass terms from renormalizable or higher order terms in the superpotential. We must then study the renormalizable and nonrenormalizable superpotential in the specific models. The cubic level and higher order terms in the superpotential are extracted by applying the rules of ref. [29]. The problem of fractionally charged states was investigated in ref. [11, 30] for the model of ref. [22]. By examining the fractionally charged states and the trilinear superpotential, it is observed that all the fractionally charged states receive a Planck scale mass by giving a VEV to the neutral singlets $\bar{\phi}_4, \bar{\phi}'_4, \phi_4, \phi'_4$ which imposes the additional F flatness constraint $(\phi_4\bar{\phi}'_4 + \bar{\phi}_4\phi'_4) = 0$. The other exotic states which are Standard Model singlets do not receive mass by this choice of flat direction. Therefore, at this level of the superpotential, the fractionally charged states can decouple from the remaining light spectrum. Similarly, the issue of fractionally charged states in the model of ref. [24] was studied in ref. [31] where it was found that all the fractionally charged states receive large mass from renormalizable or nonrenormalizable terms. Similar results were also found in the case of the Gepner models [10]. The second issue that must be examined with regard to the exotic "Wilsonian" matter is the interactions with the Standard Model states. The fractional charges of the exotic states under the unbroken $U(1)$ generators of the $SO(10)$ gauge group, may result in conserved discrete symmetry which forbid their decay to the lighter Standard Model states [8].

4 Ultra energetic neutrinos from the sun and the earth

We next examine the predicted flux of energetic neutrinos from the sun and the earth for several exotic stringy states as well as other meta–stable states, which have been proposed in the literature. The exotic states which we study are : 1) The "uniton" from eq. (3.3); 2) the Gluino–LSP of ref. [12]; 3) the "crypton" from eq. (3.2); 4) the exotic Standard Model singlet from eq. (3.4).

Before computing the detailed neutrino fluxes, it will be useful to discuss the generic features of the indirect dark matter detection with superheavy particles of mass M_X. General scaling properties can be established and will depend mainly on the type of forces inducing the interaction of dark matter particles X with matter. Ultimately, for every type of force, there is a maximal mass M_X which can be probed by annihilation in the sun or the earth.

We concentrate our analysis on the trapping and annihilation of heavy particles inside the cores of the sun and the earth [32], and the detection of the energetic neutrinos by underground detectors. Another, very important issue is the relation between direct and indirect searches for ultra heavy particles. If the density of massive particles in the galactic halo is fixed by the dark matter density ($\simeq 0.3$ GeV cm^{-3}), then the scaling of the event rate for the direct search of a dark matter candidate is proportional to σ/M_X which appears to be more advantageous than the $1/M_X^2$ fall off of the trapping rate. For large M_X, direct detection is based on registering the recoil energy of the nuclei during the elastic collision with the dark matter particle. Since the recoil energy is not sensitive to the scale M_X, when X is much heavier then the nucleus, direct detection does not allow one to gain sensitivity due to the increasing mass of the particle. On the other hand, the probability, P, for the conversion of neutrinos to muons does increase rapidly with mass (for masses $\gtrsim$ 100 GeV, the increase is stronger than linear), and indirect detection is more sensitive to the presence of dark matter. At very high masses, the conversion probability flattens out and eventually direct detection wins. We will discuss the probability function $P(M_X)$ in more detail below. Here, we will restrict our attention to indirect detection by annihilation in the cores of the sun and the earth.

Energetic neutrinos from the sun and the earth have already been discussed as a possible signal for neutralino dark matter [3, 4, 5]. Therefore, from the experimental perspective there are two relevant issues. The first is the expected flux of energetic neutrinos for each dark matter candidate and the feasibility of observing it in contemporary detectors. The second issue is the experimental prospect of distinguishing among the various dark matter candidates. The energetic neutrinos are detected by observing the muons that they produce upon interacting with matter in the vicinity of the detector. Currently the experiments can only determine the flux of muons above a certain cutoff. Therefore, the prospect for experimentally distinguishing among the candidates depends on the ability of experiments to record the actual energy of the muon or to apply a variable energy threshold.

The mechanisms responsible for the production of very energetic neutrinos inside the earth and the sun are very similar. In what follows we concentrate in details on the trapping and annihilation of heavy particles in the core of the sun and simply present similar estimates for the earth. Ultra energetic neutrinos from the sun arise if the ultra heavy particles from the galactic halo are trapped in the sun and subsequently annihilate into neutrinos. The expected flux depends on the competition between the trapping and annihilation rates. Following ref. [3], we express the initial density, n_X, of ultra heavy particles of mass, M_X, in the halo as fraction of the required dark matter density

$$M_X n_X = \delta_X \bar{v}^2 / 6\pi G a^2, \tag{4.1}$$

where $\bar{v} \simeq 300$ km/s is the r.m.s. velocity of a halo X in the solar neighborhood and $a = 10$ kpc. The coefficient $\delta_X \leq 1$ is kept as a free parameter, reflecting the partial contribution of particles X into the total dark matter density.

The elastic scattering on nuclei, $X + N \;\; \rightarrow \;\; X + N$, leads to the gravitational trapping of X particles inside the sun. The trapping rate is given by

$$\Gamma_T = (7.3 \times 10^{28} \mathrm{sec}^{-1}) \delta_X \frac{1\mathrm{GeV}}{M_X} \sum_N \sigma_{N,36} Y_N f_{E,N}, \tag{4.2}$$

where $\sigma_{N,36}$ is the elastic cross section in units of $10^{-36}\mathrm{cm}^2$ and Y_N is fraction of nuclei N by number in the sun. The most important factor here is $f_{E,N}$, the kinematic parameter, accounting for the fraction of the halo particles which loose enough energy to get trapped. In the limit of very heavy M_X, f reduces to [33]

$$f_{E,N} \simeq \frac{43 m_N}{M_X} \tag{4.3}$$

Therefore, the scaling of the trapping rate with M_X (for very large M_X) goes as

$$\Gamma_T \sim \frac{\sigma_{\mathrm{elastic}}}{M_X^2} \tag{4.4}$$

If the cross section is constant the trapping rate exhibits a quadratic fall off behavior.

Another important quantity, determining the flux of energetic neutrinos is the annihilation rate given by [3]

$$\Gamma_A = (5 \times 10^{54} s^{-1}) \left(\frac{n_{X_\odot} m_P}{\rho_\odot} \right)^2 \langle \sigma v \rangle_{A,26} W. \tag{4.5}$$

Here $\langle \sigma v \rangle_{A,26} = \langle \sigma v \rangle_A / 10^{-26}\mathrm{cm}^3\mathrm{s}^{-1}$ is the low–energy annihilation cross section times relative velocity, $n_{X_\odot}$ is the mean X particle density inside the sun, $\rho_\odot$ is the mean mass density of the sun. W is the result of the non-uniform distribution of X particles inside the sun,

$$W = \frac{\int n_{X_\odot}^2(r) dV}{n_{X_\odot}^2 V_\odot} = V_\odot \frac{\int n_{X_\odot}^2(r) dV}{(\int n_{X_\odot}(r) dV)^2} \tag{4.6}$$

From ref. [34], we take

$$W \simeq 250 \left(\frac{M_X}{1\mathrm{GeV}} \right)^{3/2} \tag{4.7}$$

The annihilation rate therefore scales as $(n_{X_\odot})^2 / (M_X)^{1/2}$ assuming a natural scaling $\langle \sigma v \rangle \sim 1/M_X^2$.

Usually, the annihilation rate is a much slower process if the density of X particles inside the sun $n_{X_\odot}$ is of the order of their density n_X in the halo. This situation is unstable, leading to the accumulation of dark matter candidates in the sun up to the level when the annihilation and trapping rates become equal

$$\Gamma_A = \Gamma_T \tag{4.8}$$

However, for large masses, the time needed to reach the equilibrium condition,

$$\tau_{eq} = (n_{X\odot} M_\odot/\rho_\odot)/\Gamma_T \simeq 2\times 10^{13}\mathrm{sec}\ \delta_X^{-1/2}\left(\frac{M_X}{1\mathrm{GeV}}\right)^{5/4}\frac{1\mathrm{GeV}}{\sqrt{M_X^2\langle\sigma v\rangle_{A,26}\sigma_{p,36}}} \qquad (4.9)$$

may turn out to be larger than the age of the sun $\tau_\odot$. For a constant elastic cross section, and an annihilation cross section which drops off as M_X^{-2}, this quantity scales as $\tau_{eq} \sim M_X^{5/4}$. The condition $\tau_{eq} < \tau_\odot$ must be checked for every particular dark matter candidate.

If equilibrium is reached, the flux of neutrinos from the annihilation inside the core of the sun is determined by the trapping rate and by an average number of neutrinos N_{eff},

$$\phi_{\nu\odot} = \frac{1}{2}\Gamma_T N_{eff}/4\pi(1\mathrm{A.U.})^2 \left\{ \begin{array}{cc} 1 & \mathrm{for}\ \ \tau_\odot \gg \tau_{eq} \\ (\tau_\odot/\tau_{eq})^2 & \mathrm{for}\ \ \tau_\odot \ll \tau_{eq} \end{array} \right. \simeq$$

$$(560cm^{-2}s^{-1})N_{eff}\delta_X\sigma_{p,36}\frac{\mathrm{GeV}^2}{M_X^2}\left\{ \begin{array}{cc} 1 & \mathrm{for}\ \ \tau_\odot \gg \tau_{eq} \\ (\tau_\odot/\tau_{eq})^2 & \mathrm{for}\ \ \tau_\odot \ll \tau_{eq} \end{array} \right. , \qquad (4.10)$$

where we have kept only the contribution of the proton in the elastic cross section. N_{eff} should be interpreted as an everage number of neutrinos *escaping* from the surface of the sun and produced as a consequence of a single annihilation event inside the core.

Using very similar assumptions our estimate for the flux of the neutrinos generated in the center of the earth is

$$\phi_{\nu\oplus} = \frac{1}{2}\Gamma_T N_{eff}/4\pi R_\oplus^2 \left\{ \begin{array}{cc} 1 & \mathrm{for}\ \ \tau_\oplus \gg \tau_{eq} \\ (\tau_\oplus/\tau_{eq})^2 & \mathrm{for}\ \ \tau_\oplus \ll \tau_{eq} \end{array} \right. \simeq$$

$$(0.05cm^{-2}s^{-1})N_{eff}\delta_X\sum_N Y_N\sigma_{N,36}\frac{\mathrm{GeV}^2}{M_X^2}\left\{ \begin{array}{cc} 1 & \mathrm{for}\ \ \tau_\oplus \gg \tau_{eq} \\ (\tau_\oplus/\tau_{eq})^2 & \mathrm{for}\ \ \tau_\oplus \ll \tau_{eq} \end{array} \right. \qquad (4.11)$$

where, of course, τ_{eq} and N_{eff} must be adjusted from those in eq. (4.10).

A more precise treatment cannot be achieved in a model independent way and therefore, we next turn to specific candidates.

5 High energy neutrinos produced by superstring relics

The initial flux of neutrinos produced close to the center of the sun or the earth can be estimated relatively easily. This flux must contain a significant portion of neutrinos with energies $\sim M_X$. The subsequent fate of highly energetic neutrinos is very different and very model and energy dependent. In the formulae for neutrino fluxes (4.10–4.11), N_{eff} is the number of neutrinos produced in a single annihilation event times the probability for high energy neutrinos to reach the vicinity of a detector.

In general, an $X\bar{X}$–pair can annihilate into various SM particles including neutrinos. If neutrinos are the direct product of the decay, we should expect to have significant absorption of the neutrinos with energies larger than few hundred GeV [35] in the core of the sun. This occurs due to the conversion of neutrinos into muons as the result of charged current interactions and leads to the exponential loss of the signal as muons are efficiently stopped by the solar media and most of them decay at rest. It does not mean, however, that the effective number of the neutrinos escaping from the surface of the sun is negligibly small even though the absorption length is smaller than the radius of the sun. Other annihilation products of an $X\bar{X}$–pair will contribute to N_{eff}. In fact, the hadronic products of $X\bar{X}$-annihilation can be an important source of neutrinos. The decay of hadrons containing different flavors lead to a broad spectrum of neutrinos as hadronic lifetimes are significantly different. The average energy of the neutrinos produced in the decay of heavy quarks except for the t–quark is much smaller than M_X and determined to be of order a few hundred GeV [35]. For these neutrinos, absorption inside the sun does not lead to a drammatic attenuation of the signal. This means that in the limit $M_X \gg 1$ TeV, the hadronic decay channels are mostly responsible for the production of neutrinos which can reach the surface of the sun without being converted into muons. Inside the Earth the absorption of the neutrinos is not the problem as the attenuation of the signal is very small up to very high energies, $M_X \sim 10^9$ GeV [37].

For our numerical estimates we will assume the following crude pattern for the spectrum. The flux of the neutrinos coming from the sun originating from the decay of X-particles with $M_X \gg 1$ TeV is downgraded in the solar media by approximately two orders of magnitude and the maximum of neutrino energy distribution is of the order of few hundred GeV. If the energy of the annhilating particles is less than 1 TeV, N_{eff} can be larger, $N_{eff} \sim O(1)$, especially if the direct channels of decay into neutrinos are open. For the earth we assume that the spectrum of neutrinos reaching the vicinity of a detector is basically the same as it is produced near the center and the fraction of very energetic neutrinos is not suppressed.

The most important charactersitic, determining the flux of neutrinos is the elastic cross section, as it can vary by orders of magnitude and even scale differently with M_X for different candidate choices. The behavior of the elastic cross section on ordinary matter for the different types of dark matter candidates we consider (stable string relics, fractionally charged states etc.) depends on several generic features of these particles. Let us assume that heavy particles X does not carry an open charge with respect to SM gauge group but is composed of constituents which are charged with respect to standard $SU(3) \times SU(2) \times U(1)$. The confining force may be the usual strong/electromagnetic interaction or a new gauge interaction as is the case for "cryptons". The size of the wave function for such a system will be determined by the *lightest* constituent and the nature of the confining force. If all constituents are heavy, of the order M_X, then the interaction between such a particle and normal matter will at least scale as $1/M_X$ and the cross section will fall off as $1/M_X^2$. The presence of a "light component" inside this particle (light quark, electrons etc.) will lead to a constant cross section, which does not vanish in the limit, $M_X \to \infty$.

The first case we consider here is the uniton. A large neutrino flux from the sun can arise in the case of uniton, as its interaction with ordinary matter is mediated by strong interactions. Depending on the type of uniton (its total spin, internal quantum numbers, etc.) the interaction with ordinary matter will be mediated by π, σ, ρ mesons. The natural estimate for the cross section is

$$\sigma_{(uniton,N)} \sim (200\mathrm{MeV})^{-2} \sim 10^{-26}\mathrm{cm}^2. \tag{5.1}$$

The independence of the elastic cross section with respect to M_X in the case of uniton is obvious. The heavy component of the exotic hadron (X particle) participates in the scattering only as a spectator. The interaction is determined by the light component(s) of the exotic hadron wave function and its strength is set by the its size and by the energies of the excited states, that is by the quantities related to Λ_{QCD}. The independence of the elastic cross section from M_X can be understood from the simple analogy with cross section of scattering of heavy atoms at low energies. The latter is determined by the size of the atoms, the wave functions of outer electrons, atomic polarizabilities, etc. and does not depend on the masses of the heavy nuclei.

For the annihilation cross section we take $\langle \sigma v \rangle_A \sim \alpha_s^2/M_X^2$. The time needed to reach the equilibrium density of unitons inside the sun is

$$\tau_{eq,uniton} \sim 6 \times 10^4 \mathrm{sec}\ \delta_X^{-1/2} \left(\frac{M_X}{1\mathrm{GeV}}\right)^{5/4}, \tag{5.2}$$

If $\delta_X \sim 1$, equilibrium is reached so long as $M_X \lesssim 10^{10} GeV$. For X-particles trapped inside the earth the time to reach equilibrium is $\sim 5 \cdot 10^4$ times larger which can be realized for $M_X \lesssim 10^6 GeV$.

Assuming equilibrium, we can easily obtain the neutrino flux by plugging the estimate (5.1) into the expression for the neutrino flux (4.10)

$$\begin{aligned} \phi_{\nu\odot} &\sim 5.6 \times 10^{12} \mathrm{cm}^{-2}\mathrm{s}^{-1} N_{eff} \delta_X \left(\frac{1\mathrm{GeV}}{M_X}\right)^2 \\ \phi_{\nu\oplus} &\sim 5 \times 10^{8} \mathrm{cm}^{-2}\mathrm{s}^{-1} N_{eff} \delta_X \left(\frac{1\mathrm{GeV}}{M_X}\right)^2, \end{aligned} \tag{5.3}$$

If $N_{eff} \sim O(10^{-2})$ and $\delta_X \sim O(1)$ these fluxes are enormous, which results in strong constraints on the mass and/or density of unitons. To compare it with existing experimental data we must include the effectiveness of the conversion of highly energetic muon neutrinos into muons. The flux of the muons can be calculated [36, 37] by including the probability P for a neutrino to convert into a muon and the muon reaching the detector. This probability rises as E^2 for energies below ~ 1 TeV. At 1 TeV, $P(E = 1\mathrm{TeV}) \sim 10^{-6}$ and flattens at higher energies [36, 37]. The actual value of P at high energies differs in these two works but for our purposes it is reasonable to take $P(E > 10^7\text{GeV}) \sim 10^{-3}$, which is a conservative estimate since the probability is certainly higher at higher energies.

Comparing the resulting flux of the energetic muons with the existing bounds from the KAMIOKANDE, MACRO and BAKSAN collaborations, $\phi_\mu \lesssim 1.4 \times 10^{-14}$ $\text{cm}^{-2}\text{s}^{-1}$ [38, 39, 40], and assuming a maximal abundance of unitons and, we deduce the bound on the uniton mass M_X:

$$M_X > 10^9 - 10^{10}\text{GeV} \tag{5.4}$$

This bound is obtained from the sun generated flux, assuming $N_{eff} \sim 0.01$ and $P \sim 10^{-6}$ and it is roughly consistent with the equilibrium conditions. The earth generated signal is less restrictive largely because the equilibrium condition is violated for $M_X > 10^6$ GeV.

In contrast to eq. (5.4), bounds on the uniton mass from its relic density due to annihilation in the early Universe require $M_X < 10^7 - 10^8$ GeV [8]. Our constraint (5.4) therefore, rules out the uniton as a dark matter candidate. It is interesting to note that the bound (5.4) is similar to the mass range required for a superheavy relic decay to produce the ultra-high energy cosmic-rays. Thus, if a strongly interacting, very massive, dark matter particle is the source of the ultra-high energy cosmic-rays, then a signal in the form of neutrinos from the sun and the earth should arise as the indirect detection experiments improve their sensitivity.

If we relax the constraint on the mass by making the cosmological abundance, δ_X, small, we can use estimates similar to the ones made above for the gluino-LSP scenario which has been advocated recently in [12]. In this case, the cosmological abundance of gluino–containing hadrons was estimated to be at the level of $\delta_X \sim 10^{-9}(M_X/1\text{GeV})$. The cross-section of the gluino-hadron with ordinary matter is quite large and can be estimated to be 10^{-26} cm^2 as in the case with uniton. For a gluino mass of order 100 GeV, this abundance would lead to a muon flux from the sun five orders of magnitude above current experimental limit, even if we assume a low value for $N_{eff} \sim 0.01$. Thus, here too, the indirect detection limits provide a strong constraint and are complementary to the astophysical constraints discussed in [13]. The gluino-LSP bound state is forbidden to roam freely in the galactic halo and as such can not provide a source for the ultra-high energy cosmic-rays unless it is infact unstable and therefore not the LSP (or R-parity is broken). As we noted earlier, our constraints do not apply to the light gluino scenario proposed in [20], since the neutrino energies will be below threshold for the existing detectors.

Let us consider now the case of "cryptons", neutral objects composed from the constituents which bear $U(1)_{em}$ charge ("cryptons") held together by a confining force of non–SM origin. It is clear that their interaction with ordinary matter is suppressed because of the absence of light constituents in their structure. The absence of an overall charge with respect to the electromagnetic gauge group does not preclude these particle from possessing an anomalous magnetic moment provided that they have spin and X and $\bar{X}$ are different. The size of the anomalous magnetic moment depends on the details of the force confining the charged constituents together. It is natural to assume, however, that the size of the anomalous magnetic moment is similar to an "exotic Bohr magneton", $e/(2M_X)$. In this case the fall off of the cross section with M_X is just quadratic,

$$\sigma_p \sim \frac{\alpha^2}{M_X^2} \tag{5.5}$$

Still, this decrease is very sharp (especially if we are interested in very massive states) and the corresponding neutrino flux is low. Inserting (5.5) into the neutrino flux from the sun (4.10), we find

$$\phi_{\nu\odot} \sim 2.2 \times 10^{11}\mathrm{cm}^{-2}\mathrm{s}^{-1}\alpha^2 N_{eff}\delta_X \left(\frac{1\mathrm{GeV}}{M_X}\right)^4. \tag{5.6}$$

The earth signal is additionally suppressed as most of the nuclei inside the earth are spinless. Normalizing the probability of conversion to muons at the neutrino energy scale of 1 TeV, $P = 10^{-6}P_6$, we find the following constraint

$$N_{eff}\delta_X P_6 \left(\frac{1\mathrm{TeV}}{M_X}\right)^4 < 1.2 \times 10^{-3} \tag{5.7}$$

Equilibrium is reached for M_X lighter than a few TeV, if we take annihilation and elastic scattering cross sections to be of the same order, $\langle\sigma_A v\rangle \sim \sigma_p$. Though (5.7) is significantly weaker than the bound on the uniton, it is not altogether trivial.

The last case we consider is that of the pure SM singlets, charged with respect to an extra $U(1)$ group. The interaction of these particles with ordinary matter will depend on the mass of the gauge boson associated with this additional group. The size of the elastic cross section mediated by exotic Z' can be estimated as

$$\sigma_p \sim \frac{g_X'^2}{4\pi}\frac{g_p'^2}{4\pi}\frac{m_p^2}{M_{Z'}^4}. \tag{5.8}$$

Here g_X' and g_p' denote total charges of exotic and ordinary matter with respect to this additional $U(1)$ group. If Z' is very heavy, of the order M_X, then this cross section is very small. $M_{Z'}$ can be kept as a free parameter, not necessarily related to M_X. However, direct searches for Z''s place a lower limit on its mass at the level of 500 GeV, which is still too heavy to produce a large elastic cross section. Assuming $M_{Z'} < M_X$, we obtain the following constraint on on the mass of Z' and the combination of charges,

$$N_{eff}\delta_X P_6 \frac{g_X'^2 g_p'^2}{e^4}\frac{1\mathrm{TeV}^6}{M_{Z'}^4 M_X^2} < 1.2 \times 10^{-3} \tag{5.9}$$

Numerically, the limits derived for the two cases, "cryptons" and SM singlets with Z'–charge, are similar to those placed on a conventional neutralino dark matter scenario, suggesting that the current sensitivity to the masses of these particles cannot be much better than few TeV.

6 Discussions and conclusions

Perhaps more exciting than setting a limit on exotic states is the possibility of detecting a high energy neutrino signal from the sun or the earth. The detection of an ultra high energy neutrino from the earth or the sun would in fact imply a lower limit to the fundamental cut–off scale. If the origin of the neutrino can be correlated with the position of the sun or with the direction to the center of the earth, one would be forced to conclude that such a neutrino could only be produced via the annihilation (or decay) of a massive particle with $m \sim E_\nu$. In models with large extra dimensions, the fundamental cut–off scale is necessarily below the 4-dimensional Planck scale as given by

$$L^n M^{2+n} = M_P^2$$

for the compactification of a $4+n$ dimensional theory where M is the fundamental cut–off scale and L is the characteristic size of the extra dimensions. It is clear that in such a model, the maximum particle mass is also M. For example, in the $4+n$ dimensional theory, the maximum mass of a scalar particle is $m \sim M$ (above M, the full quantum theory would be required). The scalar mass term $m^2\hat{\phi}^2$ where $\hat{\phi}$ is the $4+n$-dimensional scalar, becomes $m^2\phi^2$, where $\phi = L^n\hat{\phi}$ after the scalar kinetic term is renormalized. Thus the limit $m \leq M$ is maintained in four dimensions. The limit also applies to the KK excitations of a lighter field. In 4-dimensions, the KK masses of the excitations scale as $m_{KK} \sim n/L$. However the number of such states is limited by the maximum momentum in the extra dimension which is M. Therefore, the highest mass state we can consider is $N/L \sim M$. The same argument holds for the maximum mass of a fermion as well. Therefore, any detection of a high energy neutrino for which we can be sure was produced in the decay or annihilation of a particle with an ultra heavy mass M_X, sets a firm lower bound on the fundamental scale, $M \geq M_X$.

If we are interested in setting a strong constraint on the fundamental cut–off scale (that is one which is significantly stronger than the current experimental limit of $M \gtrsim 1$ TeV), a detector with sufficiently high angular resolution is needed. The large water detectors such as Kamiokande have an angular resolution of order 1°. In this case, energies above 10's of GeV can not be distinguished and an interesting constraint on the fundamental cut–off scale can not be set. In contrast, a detector such as Soudan II, which has an angular resolution of order 0.1°, will be able to correlate only very energetic ($E \gg 1$ TeV) neutrinos with the centers of the sun or the earth. Of course, such a constraint requires a positive detection of ultra-high energy neutrinos from the sun and/or the earth.

In summary, we have argued that the search for the highly energetic neutrinos coming from the annihilation of very massive particles inside the sun and the earth can be used to probe energy ranges above the electroweak scale.

The flux of highly energetic neutrinos is crucial to the indirect detection of dark matter. The flux has a universal scaling behavior, $1/M_X^2$, if the magnitude of the elastic cross section is constant. However, the actual fall off (with M_X) of the induced muon flux is much milder, since the probability of the conversion of neutrinos to muons rises. The elastic cross section with ordinary matter is of course candidate specific, and its magnitude depends on the nature of forces mediating the interaction and on the presence of a "light component" in the wave function of the exotic state.

In the case of the uniton and gluino-LSP, there are strong interactions which induce the elastic cross section with the ordinary matter and the exotic particle has a light component u, d or A^a_μ in it. As a result, the trapping rate is large leading to very large neutrino signal unless the the mass of X particle is bigger than 10^9 GeV or their initial density is much smaller than one needs for dark matter energy density. In the case of the gluino-LSP candidate, even the low density expected for these particles does not preclude a large detection rate of neutrinos from the sun.

We believe that the dedicated search for ultra-high energy neutrinos coming from the sun can be very important in setting a lower bound on the "fundamental scale". A signal, correlated with the positions of the sun and/or the center of the earth, can only be the consequence of the decay or annihilation of heavy particles inside the sun/earth. At $M_X > 1$ TeV, neutrinos from the center of the earth will be more energetic and a significant fraction of the spectrum will have energies of order M_X. In contrast, the maximum of the neutrino spectrum coming from the sun appears to be shifted towards energies less than 1 TeV. Once the neutrino signal is identified, one can use the shape of the energy spectrum of the muons and try to reconstruct the value of the mass for these particles which can be well above the reach of accelerator physics. The existence of heavy particles with masses M_X would imply that the fundamental (string/quantum gravity) scale is at least as large as M_X. It may serve as a new way of constraining scenarios with large extra dimensions and low (as low as 1 TeV) fundamental Plank scale. However, this constraint does require a positive signal rather than a flux limit.

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References

[1] For a recent review see *e.g.*: L. Nodulman, hep-ex/9810035.

[2] N. Arkani–Hamed, S. Dimopoulos and G. Dvali, *Phys. Rev.* **D59** (1999) 086004.

[3] J. Silk, K.A. Olive, and M. Srednicki, *Phys. Rev. Lett.* **55** (1985) 257;
M. Srednicki, K.A. Olive, and J. Silk, *Nucl. Phys.* **B279** (1987) 804.

[4] See *e.g.* G. Jungman, M. Kamionkowski and K. Griest, *Phys. Rep.* **267** (1996) 195.

[5] L. Bergstrom, J. Edsjo, and P. Gondolo, *Phys. Rev.* **D58** (1998) 103519.

[6] Y. Hosotani, *Phys. Lett.* **B126** (1983) 309; *Phys. Lett.* **B129** (1983) 193.

[7] X.G. Wen and E. Witten, *Nucl. Phys.* **B261** (1985) 651;
G.G. Athanasiu, J.J. Atick, Michael Dine and Willy Fischler, *Phys. Lett.* **B214** (1988) 55;
A. Schellekens, *Phys. Lett.* **B237** (1990) 363.

[8] S. Chang, C. Coriano and A.E. Faraggi, *Phys. Lett.* **B397** (1997) 76; *Nucl. Phys.* **B477** (1996) 65.

[9] J. Ellis, J.L. Lopez and D.V. Nanopoulos, *Phys. Lett.* **B247** (1990) 257.

[10] P. Huet, *Nucl. Phys.* **B350** (1991) 375.

[11] A.E. Faraggi, *Phys. Rev.* **D46** (1992) 3204.

[12] S. Raby, *Phys. Rev.* **D56** (1997) 2852; *Phys. Lett.* **B422** (1998) 158;
S. Raby and K. Tobe, *Nucl. Phys.* **B539** (1999) 3; H. Baer, K. Cheung and J.F. Gunion, *Phys. Rev.* **D59** (1999) 075002.

[13] R.N. Mohapatra and V.L. Teplitz, *Phys. Rev. Lett.* **81** (1998) 3079.

[14] J. Ellis, J.L. Lopez and D.V. Nanopoulos, *Phys. Lett.* **B247** (1990) 257.

[15] K. Greisen, *Phys. Rev. Lett.* **16** (1966) 748;
G.T. Zatsepin and V.A. Kuzmin, *Pisma Zh. Eksp. Theor. Fiz.* **4** (1966) 114.

[16] V. Berezinsky, M. Kachelriess, and A. Vilenkin, *Phys. Rev. Lett.* **79** (1997) 4302;
V. Berezinsky, hep-ph/9802351; astro-ph/9811268;
V.A. Kuzmin and V.A. Rubakov, *Phys. Atom. Nucl.* **61** (1998) 1028.

[17] M. Birkel and S. Sarkar, *Astropart. Phys.* **9** (1998) 297.

[18] K. Benakli, J. Ellis, and D.V. Nanopoulos, *Phys. Rev.* **D59** (1999) 047301.

[19] K. Hamaguchi, Y. Nomura, and T. Yanagida, *Phys. Rev.* **D58** (1998) 103503; *Phys. Rev.* **D59** (1999) 063503;
K. Hamaguchi, K.I. Izawa, Y. Nomura, and T. Yanagida, hep-ph/9903207.

[20] D.J.H. Chung, G.R. Farrar, and E.W. Kolb, *Phys. Rev.* **D57** (1998) 4606.

[21] I. Antoniadis, J. Ellis, J. Hagelin, and D.V. Nanopoulos, *Phys. Lett.* **B231** (1989) 65;
J. Lopez, D.V. Nanopoulos, and K. Yuan, *Nucl. Phys.* **B399** (1993) 654.

[22] A.E. Faraggi, D.V. Nanopoulos and K. Yuan, *Nucl. Phys.* **B335** (1990) 347.

[23] I. Antoniadis, G.K. Leontaris and J. Rizos, *Phys. Lett.* **B245** (1990) 161;
G.K. Leontaris, *Phys. Lett.* **B372** (1996) 212;
G.K. Leontaris and J. Rizos, hep-th/9901098 .

[24] A.E. Faraggi, *Phys. Lett.* **B278** (1992) 131; *Phys. Lett.* **B274** (1992) 47.

[25] A.E. Faraggi, *Nucl. Phys.* **B387** (1992) 239; hep-th/9708112.

[26] A.E. Faraggi, *Nucl. Phys.* **B407** (1993) 57; *Phys. Lett.* **B326** (1994) 62.

[27] A.E. Faraggi, *Phys. Lett.* **B302** (1993) 202;
K.R. Dienes and A.E. Faraggi, *Phys. Rev. Lett.* **75** (1995) 2646; *Nucl. Phys.* **B457** (1995) 409.

[28] S. Chaudhuri, G. Hockney and J. Lykken, *Nucl. Phys.* **B469** (1996) 357;
G.B. Cleaver *et. al.* , *Nucl. Phys.* **B525** (1998) 3; *Nucl. Phys.* **B545** (1998) 47; *Phys. Rev.* **D59** (1999) 055005; *Phys. Rev.* **D59** (1999) 115003.

[29] S. Kalara, J. Lopez, and D.V. Nanopoulos, *Nucl. Phys.* **B353** (1991) 650.

[30] G.B. Cleaver, A.E. Faraggi and D.V. Nanopoulos, hep-ph/9811427; hep-ph/9904301.

[31] A.E. Faraggi, *Nucl. Phys.* **B403** (1993) 101.

[32] L. Krauss, M. Srednicki, and F. Wilczek, *Phys. Rev.* **D33** (1986) 2079;
K. Freese, *Phys. Lett.* **B167** (1986) 295.

[33] W.H. Press and D.N. Spergel, *Ap.J.* **296** (1985) 679.

[34] K. Griest and D. Seckel, *Nucl. Phys.* **B283** (1987) 681.

[35] S. Ritz and D. Seckel, *Nucl. Phys.* **B304** (1988) 877.

[36] T.K. Gaisser and T. Stanev, *Phys. Rev.* **D31** (1985) 2770.

[37] M. Honda and M. Mori, *Prog. Theor. Phys.* **78** (1987) 963.

[38] M. Mori *et al.*, *Phys. Rev.* **D48** (1993) 5505.

[39] M. Ambrosio *et al.*, hep-ex/9812020.

[40] M.M. Boliev, *et al.*, *Nucl. Phys. Proc. Supp.* **B48** (1996) 83.