

Cosmological and phenomenological implications of Wilsonian matter in realistic superstring derived models*

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Abstract

Superstring phenomenology aims at achieving two goals. The first is to reproduce the observed physics of the Standard Model. The second is to identify experimental signatures of superstring unification which, if observed, will provide further evidence for the validity of superstring theory. I discuss such potential signatures of superstring unification. I propose that proton lifetime constraints imply that the Standard Model gauge group must be obtained directly at the string level. In this case the unifying gauge group, for example $SO(10)$, is broken to the Standard Model gauge group by “Wilson lines”. The symmetry breaking by “Wilson line” has important implications. It gives rise to exotic massless states which cannot fit into multiplets of the original unifying gauge group. This is an important feature because it results in conservation laws which forbid the interaction of the exotic “Wilsonian” states with the Standard Model states. The “Wilsonian” matter states then have important phenomenological implications. I discuss two such implications: exotic “Wilsonian” states as dark matter candidates and “Wilsonian” matter as the messenger sector in gauge mediated dynamical SUSY breaking scenarios.

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Superstring phenomenology aims at achieving two goals. The first is to reproduce the observed physics. The second is to identify possible experimental signature of superstring unification which may provide further evidence for its validity. The first task is highly nontrivial and indeed only a handful of string models can claim to be potentially realistic. Indeed the number of constraints is large and satisfying all in one string model is an almost impossible challenge. A model which satisfies all of the constraints of the observed low energy physics, is likely to be more than an accident. Such a model, or class of models, will then serve as the laboratory for the search for exotic predictions of superstring unification. It will also serve as a laboratory in which we can address the important question of how the string vacuum is selected.

A few of the constraints that a realistic model of unification must satisfy are listed below.

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1. Gauge group $\longrightarrow SU(3) \times SU(2) \times U(1)_Y$
 2. Contains three generations
 3. Proton stable $(\tau_P > 10^{30+} \text{ years})$
 4. N=1 supersymmetry (or N=0)
 5. Contains Higgs doublets \oplus potentially realistic Yukawa couplings
 6. Agreement with $\underline{\sin^2 \theta_W}$ and $\underline{\alpha_s}$ at M_Z (+ other observables).
 7. Light left-handed neutrinos
 8. $SU(2) \times U(1)$ breaking
 9. SUSY breaking
 10. No flavor changing neutral currents
 11. No strong CP violation
 12. Exist family mixing and weak CP violation
 13. + ...
 14. + **GRAVITY**

The first question that we must ask is whether it is possible to construct a model which satisfies all of those criteria, or possibly a class of models which can accommodate most of these constraints. To date the most developed theory that can consistently unify gravity with the gauge interactions is string theory [1]. While alternatives may exist, it makes sense at this stage to try to use string theory to construct a model which satisfies the above requirements. Even if eventually string theory turns out not to be the fundamental theory of nature, a model which satisfies all of above constraints is likely to arise as an effective model from the true fundamental theory.

Semi-realistic string models were constructed by using various methods [2, 3, 4, 5, 6, 7, 8]. Of the above requirements the most difficult to satisfy in a realistic model is the constraint of proton stability. The reason is that supersymmetric models are in general plagued with dimension four and five operators which give rise to fast proton decay [9]. In supersymmetric point field theory models such operators may be

avoided by postulating the existence of some symmetries which forbid the dangerous operators. However, in string models we don't have this luxury. The desired symmetries either exist in the models or they do not. Often one can find that at some points in the moduli space of specific models the dangerous operators vanish due to some accidental global symmetry. This is not quite satisfactory for two reasons. First, in general the isolated points in the moduli space will not accommodate some other constraints, like potentially non-vanishing Yukawa couplings which can give rise to fermion masses. Second, in general global symmetries are badly broken in string models and the dangerous operators can be induced from non-vanishing nonrenormalizable terms which effectively reproduce the dangerous dimension four and five operators. The longevity of the proton lifetime imposes that such nonrenormalizable operators must be suppressed to very high orders. For these reasons, we would like the desired symmetry, which suppresses the proton decay operators, to be a robust gauge symmetry or a local discrete symmetry.

The problem of proton stability suggests that the allowed gauge groups at the string scale are restricted to very few choices. First, it is desirable to avoid giving a VEV of the order of the GUT or Planck scale to the right handed neutrino. Otherwise, effective dimension four operators may be induced. Second, if the symmetry is broken at the string level to $SO(6) \times SO(4)$ or directly to $SU(3) \times SU(2) \times U(1)^2$ then there is a superstring doublet-triplet splitting mechanism in which the triplets are projected from the massless spectrum by the GSO projections while the doublets remain in the light spectrum [10]. Therefore, if the symmetry at the string scale is broken to $SO(6) \times SO(4)$ or directly to the Standard Model gauge group the problems with proton decay can be avoided in a robust way. Thus, due to proton lifetime constraints, the preferred symmetries at the string level are $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)^2$.

Another very restrictive constraint on realistic superstring models is the requirement of agreement with $\sin^2 \theta_W$ and α_s at M_Z . This constraint is better known as the string scale gauge coupling unification problem. If we assume that the spectrum between the electroweak scale is that of the Minimal Supersymmetric Standard Model, then it is well known that the three gauge couplings intersect at a scale which is of the order 2×10^{16} GeV [11]. While the successful meeting of the couplings in the MSSM is intriguing, it is far from being well motivated. The MSSM is not a complete theory and clearly cannot accommodate all of the requirements listed above. Furthermore, there is nothing special about the spectrum of the MSSM. The assumption of a minimal spectrum is ad hoc and is not motivated from any fundamental principles. On the other hand, at tree level, string theory predicts that the gauge couplings are unified at a scale which is of the order 4×10^{17} GeV [12]. Thus, an order of magnitude separates the string unification scale from the MSSM unification scale. It would seem that in extrapolation of the gauge parameters over some fifteen orders of magnitude, a problem involving a single order of magnitude would have many possible solutions. Surprisingly, however, the problem is not easily resolved. In fact, most string models

can immediately be discarded simply because they predict a value for $\sin^2 \theta_W$ at the string scale which is much smaller than the regular GUT prediction. Thus, these string models predict a $\sin^2 \theta(M_Z)$ which is much smaller than the experimentally observed value and cannot be adjusted by small correcting effects. This is the case because in most string models the weak hypercharge does not have the regular GUT embedding. Thus, in a realistic string models we would like the weak hypercharge to have the standard GUT embedding. One possibility is of course to consider string GUT models [14]. However, those will in general have problems with proton lifetime. Another proposal [15] is that nonperturbative string effects play an important role, and that one of the compactified dimensions is of the order of the GUT scale. In this case there may be additional GUT scale color triplets from the massive string spectrum and one has to ascertain that those do not cause rapid proton decay.

The class of realistic string models which are constructed in the free fermionic formulation [13] can satisfy both of these requirements. In fact, to date, these are the only string models that have been shown to satisfy both of these constraints. In these models an $SO(10)$ symmetry is broken at the string scale to $SU(5) \times U(1)$, $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)^2$ and the weak hypercharge has the standard $SO(10)$ embedding. Also there exist free fermionic models in which all the dangerous dimension four and five operators are suppressed to all orders of nonrenormalizable terms [10]. This is achieved in models in which the $SO(10)$ symmetry is broken to $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)^2$ and due to the superstring doublet–triplet splitting mechanism. Recently, an interesting study of the proton decay problem in the context of string models was done in Ref. [16]. It was found that precisely the sort of symmetries which appear in the free fermionic models are those needed to prevent proton decay and allow naturally suppressed neutrino masses.

It has been shown that the free fermionic models can potentially satisfy many of the other constraints that must be imposed on a realistic string model [17]. Thus, free fermionic models are candidates for a realistic string model. After identifying a class of potentially realistic string models the second question that we must address is whether there exist some generic signature of these models which, if observed, will provide further support for their validity.

In this talk I discuss such possible exotic signatures. In the free fermionic models, and in string models in general, one starts with a larger symmetry group which is subsequently broken to some intermediate unifying symmetry by the GSO projection. For example in the free fermionic models we start with a $SO(44)$ gauge group which is broken to $SO(10) \times SO(6)^3 \times E_8$. In the free fermionic models the breaking is achieved by defining boundary condition basis vectors for the world–sheet fermions which satisfy some string consistency constraints [13]. These intermediate gauge symmetry is broken further by means of additional boundary condition basis vectors. In particular the $SO(10)$ symmetry, in which the Standard Model is embedded is broken to one of its subgroups, $SU(5) \times U(1)$, $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)^2$. These additional breaking of the $SO(10)$ gauge symmetry gives rise to massless

states that cannot fit into multiplets of $SO(10)$. The additional vectors which break the $SO(10)$ symmetry correspond to Wilson lines in the orbifold models. I refer to the extra matter which arises from these sectors as Wilsonian matter. The breaking by Wilson lines may give rise to local discrete symmetries which forbid the interaction of the Wilsonian states with the Standard Model states. This unique stringy phenomena has important implications. It implies that the exotic Wilsonian matter states are stable and therefore may be good dark matter candidates.

To understand better the Wilsonian matter phenomena, it is useful to study the general structure of the realistic free fermionic models. As mentioned above the basis vectors which define the models are divided into two parts. The first part consists of the five vectors of the NAHE set [3, 8]. These basis vectors correspond to $Z_2 \times Z_2$ orbifold compactification. The Neveu-Schwarz (NS) sector corresponds to the untwisted sector of the orbifold models and produces the generators of the $SO(10) \times SO(6)^3 \times E_8$ gauge group. In addition to the spin two and spin one states the NS sector also produces three 10 representation of $SO(10)$ and several $SO(10)$ singlets. The three sectors b_1 , b_2 and b_3 correspond to the three twisted sectors of the $Z_2 \times Z_2$ orbifold model and produce 48 multiplets in the 16 representation of $SO(10)$.

The correspondence between this class of free fermionic models and orbifold models can be shown by constructing the same orbifold compactification which corresponds to some free fermionic models [18]. The simplest way to illustrate this is by adding to the NAHE set a basis vector X ,

$$X = \{\bar{\psi}_{1,\dots,5}, \bar{\eta}^1, \bar{\eta}^2, \bar{\eta}^3\} = 1,$$

which extends the symmetry to $E_6 \times SO(4)^3 \times E_8$, with $N = 1$ supersymmetry and 24 generations in the 27 representation of E_6 . The same model is constructed in the orbifold formulation by first constructing the toroidally compactified Narain lattice [19] which corresponds to the free fermionic point in the moduli space. The metric and antisymmetric tensors are give by

$$g_{ij} = \begin{pmatrix} 2 & -1 & 0 & 0 & 0 & 0 \\ -1 & 2 & -1 & 0 & 0 & 0 \\ 0 & -1 & 2 & -1 & 0 & 0 \\ 0 & 0 & -1 & 2 & -1 & -1 \\ 0 & 0 & 0 & -1 & 2 & 0 \\ 0 & 0 & 0 & -1 & 0 & 2 \end{pmatrix} \quad b_{ij} = \begin{cases} g_{ij} & i > j \\ 0 & i = j \\ -g_{ij} & i < j \end{cases}$$

where the metric is the Cartan matrix of $SO(12)$. At the point in the moduli space which corresponds to the free fermionic models, all the radii of the compactified dimensions are fixed at a specific value. At this point there are additional massless vectors bosons which arise due to the wrapping of the string on the compactified dimensions. At the specific point which correspond to the free fermionic models the gauge symmetry is enhanced to $SO(12) \times E_8 \times E_8$. If we now mod out by the $Z_2 \times Z_2$ discrete symmetry with standard embedding, we obtain a model with a

gauge group and spectrum which are the same as those which were obtained in the fermionic construction. This indeed demonstrates that the basis vectors of the NAHE set correspond to $Z_2 \times Z_2$ orbifold models. In the realistic free fermionic models we replace the vectors X , which extends the $SO(10)$ symmetry to E_6 , with the vector 2γ ,

$$2\gamma = \{\bar{\psi}_{1,\dots,5}, \bar{\eta}^1, \bar{\eta}^2, \bar{\eta}^3, \bar{\phi}_{1,\dots,4}\} = 1.$$

With this substitution the gauge group is $SO(10) \times U(1)^3 \times SO(4)^3 \times SO(16)$ with $N = 4$ supersymmetry and 24 generation in the 16 of $SO(10)$. The vectors $b_j + 2\gamma$ now give 24 multiplets in the 16 representation of the hidden $SO(16)$ gauge group. The modification from the vector X to the vector 2γ can be regarded as a transition from a $(2, 2)$ model to a $(2, 0)$ model, and can also be achieved by redefining the GSO phases.

At the level of the NAHE set the observable gauge symmetry is $SO(10) \times SO(6)^3$ and the number of generations is 48, sixteen from each of the sectors b_1 , b_2 and b_3 . The number of generations is reduced to three by adding three additional boundary condition basis vectors to the NAHE set. Each one of the sectors b_1 , b_2 and b_3 produces exactly one generation. We observe that the emergence of three generations in the realistic free fermionic models is deeply rooted in the underlying $Z_2 \times Z_2$ orbifold structure. Each one of the generations is in the 16 representation of $SO(10)$ decomposed under the final subgroup of the original $SO(10)$ gauge group. Thus, we see that the Standard Model spectrum in this models has an underlying $SO(10)$ unification structure and the weak hypercharge has the standard $SO(10)$ embedding. This outcome is very important as these models will share some of the desirable features of $SO(10)$ GUT unification. The only difference is that here the $SO(10)$ symmetry is broken at the string theory level rather than at the field theory level. This distinction has of course highly nontrivial consequences. In particular, as mentioned above with regard to the proton lifetime problem.

The breaking of the $SO(10)$ gauge symmetry can be done with the same boundary condition basis vectors which reduce the number of generations. The $SO(10)$ symmetry is broken to one of its subgroups. This is achieved by the assignment of boundary conditions to the set $\bar{\psi}^{1,\dots,5}$:

$$b\{\bar{\psi}_{\frac{1}{2}}^{1\dots 5}\} = \left\{\frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2}\right\} \Rightarrow SO(10) \rightarrow SU(5) \times U(1), \quad (1)$$

$$b\{\bar{\psi}_{\frac{1}{2}}^{1\dots 5}\} = \{1 \ 1 \ 1 \ 0 \ 0\} \Rightarrow SO(10) \rightarrow SO(6) \times SO(4). \quad (2)$$

To break the $SO(10)$ symmetry to $SU(3) \times SU(2) \times U(1)_C \times U(1)_L^*$ both steps, (1) and (2), are used, in two separate basis vectors. The flavor $SO(6)^3$ symmetries are also broken by the additional basis vectors to horizontal $U(1)$ symmetries. Similarly, the hidden E_8 gauge group is broken as well to one of its subgroups. For example in the standard-like models of refs. [6, 7], $SO(6)^{1,2,3} \rightarrow U(1)^{1,2,3} \times U(1)^{4,5,6}$, and

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

From the sector $b_1 + b_2 + \alpha + \beta$

Higgs doublets

$$h_{\alpha\beta}{}_{-\frac{1}{2}, -\frac{1}{2}, 0, 0, 0, 0} \quad \bar{h}_{\alpha\beta}{}_{\frac{1}{2}, \frac{1}{2}, 0, 0, 0, 0}$$

$SO(10)$ singlets with $U(1)$ charges $\Phi_{\alpha\beta}$, Φ_1^\pm , Φ_2^\pm , Φ_3^\pm

All these states fit into “standard” reps. of $SO(10)$

As seen above the sectors which are produced by the NAHE set and the combination $\alpha + \beta$ produce states which fit into standard $SO(10)$ representations or are $SO(10)$ singlets. In addition to the “standard” spectrum from the sectors above, there exist in the fermionic models “exotic” spectrum which cannot fit into $SO(10)$ multiplets. These spectrum arises from sectors which are combinations of the NAHE basis vectors and the basis vectors $\{\alpha, \beta, \gamma\}$. These combinations produce the exotic matter in vector-like representations. In general, unlike the “standard” spectrum, the “exotic” spectrum is highly model dependent. We can however classify the exotic matter according the pattern of the $SO(10)$ symmetry breaking by the specific sectors. Each of these sectors breaks the $SO(10)$ symmetry to $SU(5) \times U(1)$, $SO(6) \times SO(4)$ or $SU(3) \times SU(2) \times U(1)^2$. Thus, in $SU(5) \times U(1)$ models only one type of the exotic states can appear. Similarly, in $SO(6) \times SO(4)$ type models another type of exotic states can appear. Finally, the $SU(3) \times SU(2) \times U(1)^2$ type models contain both the $SU(5) \times U(1)$ and $SO(6) \times SO(4)$ type states as well as states which are unique to $SU(3) \times SU(2) \times U(1)^2$ type models.

Using the notation

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

where $U(1)_Y = 1/3U(1)_C + 1/2U(1)_L$, $U(1)_{Z'} = U(1)_C - U(1)_L$ and $U(1)_{e.m.} =$

$T_{3L} + U(1)_Y$.

The following exotic states appear in the free fermionic models

Exotic matter

$SO(6) \times SO(4)$ type states

$$[(\bar{3}, \frac{1}{2}); (1, 0)]_{(1/6, 1/2, 1/6)} \quad ; \quad [(\bar{3}, -\frac{1}{2}); (1, 0)]_{(-1/6, -1/2, -1/6)}$$

$$[(1, 0); (2, 0)]_{(0, 0, \pm 1/2)}$$

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

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

$SU(5) \times U(1)$ type states

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

$SU(3) \times SU(2) \times U(1)^2$ type states

$$\begin{aligned}
& [(3, \frac{1}{4}); (1, \frac{1}{2})]_{(-1/3, -1/4, -1/3)} \quad ; \quad [(\bar{3}, -\frac{1}{4}); (1, \frac{1}{2})]_{(1/3, 1/4, 1/3)} \\
& [(1, \pm\frac{3}{4}); (2, \pm\frac{1}{2})]_{(\pm 1/2, \pm 1/4, (1,0); (0,-1))} \\
& [(1, \pm\frac{3}{4}); (1, \mp\frac{1}{2})]_{(0, \pm 5/4, 0)}
\end{aligned}$$

Thus the exotic states which appear in the $SU(5) \times U(1)$ and $SO(6) \times SO(4)$ type sectors are fractionally charged with electric charges $\pm 1/2$. The $SU(3) \times SU(2) \times U(1)^2$ type sectors give rise to matter states which have the “standard” charges under the Standard Model gauge group but carry “fractional” charges under the $U(1)_{Z'}$ gauge symmetry.

To examine the phenomenology of these exotic states we have to study their interactions. The interaction terms in the superpotential are obtained by calculating the correlators between vertex operators [20]

$$\langle V_1^f V_2^f V_3^b \dots V_N^b \rangle$$

The non-vanishing correlators must be invariant under all the symmetries and the string selection rules.

The free fermionic models contain an anomalous $U(1)$ gauge symmetry. The anomalous $U(1)$ generates a Fayet–Iliopoulos D-term which breaks supersymmetry at the Planck scale [21]. The anomaly is canceled by assigning VEVs to some Standard Model singlets in the massless string spectrum. In general, these singlets are charged also with respect to other $U(1)$ symmetries of a given string model. Therefore, requiring that all the D-terms vanish imposes a nontrivial set of constraints on the allowed VEVs. In addition we require that all the F-terms vanish which imposes that the superpotential and all of its derivatives vanish to all order of nonrenormalizable terms. Some of the higher order nonrenormalizable terms become effective renormalizable operators due to the VEVs which are used to cancel the anomalous $U(1)$ D-term equation.

The implications of the “Wilsonian” matter states were studied in detail in several examples of realistic free fermionic standard-like models. The first example is the model of ref. [24]. In this model the $SO(10)$ gauge group is broken to $SU(3) \times SU(2) \times U(1)^2$. It contains three generations and electroweak Higgs doublet which can generate phenomenologically realistic fermion mass spectrum. The top Yukawa has been calculated from cubic order term in the superpotential. The bottom quark and tau lepton Yukawa were calculated from quartic order terms and there exist mass and generation mixing terms for the lighter two generations from higher order terms. In this model there are no dimension four operators which can mediate proton decay due to a custodial gauge symmetry [25]. All the “standard” color triplets are projected out from the massless spectrum. As a result all the dimension five operators vanish as well. Thus, in this model proton decay may only arise from the massive string spectrum. This model contains exotic “Wilsonian” color triplets and electroweak doublets in vector-like representations. String-scale gauge coupling unification in this model is compatible with low energy data, provided that the “Wilsonian” color triplets and electroweak doublets exist at the appropriate mass thresholds.

The exotic “Wilsonian” matter states appear in the free fermionic models in vector-like representations, and obtain mass terms from cubic level or nonrenormalizable terms in the superpotential.

The first example of exotic “Wilsonian” matter states are the states with fractional electric charge which appear in $SU(5) \times U(1)$ and $SO(6) \times SO(4)$ type sectors. As there are strong limits on the existence of free fractionally charged states, this states must be very massive, confined into integrally charged states, or diluted away. For example, in the flipped $SU(5)$ model [3, 22] all the fractionally charged states transform under a non-Abelian hidden gauge group and are therefore confined. In ref. [23] it was shown that all the fractionally charged states in the model of ref. [4] get mass terms at the cubic level of the superpotential by giving VEVs to four Standard Model singlets in the spectrum of that model along a Flat F and D direction. Thus, in this model all the fractionally charged states receive Planck scale mass and decouple from the light spectrum. Finally, the fractionally charged states may of course be diluted away. In which case their abundance today may be too small to be detected experimentally.

In addition to the fractionally charged states, the free fermionic standard-like models contain “Wilsonian” states which carry the regular charges under the Standard Model gauge group but carry “fractional” charges under the $U(1)_{Z'}$ symmetry, which exist in $SO(10)$. These states can be color triplets, electroweak doublets, or Standard Model singlets and may be good candidates for dark matter. The first example of such states are the color triplets. The existence of color triplets at intermediate energy scale is motivated from the problem of gauge coupling unification. The analysis of ref. [26] showed that heavy string thresholds [27], light SUSY thresholds, intermediate gauge structure, hypercharge normalization [28] do not resolve the problem. The analysis there suggests that the only way to resolve the problem is the existence of additional intermediate matter thresholds, beyond the spectrum of the MSSM [29]. The additional color triplet thresholds are in general much lighter than the additional electroweak doublets. The extra color triplets and electroweak doublets do appear in some string models [24]. Specifically, in the model of ref. [24] the extra states appear from “Wilsonian” sectors and are therefore candidates for “Wilsonian” dark matter. Due to its role in the string gauge coupling unification problem, this type of color triplet is referred to as the uniton [30].

The uniton has the “standard” charges under the standard model gauge group and “fractional” $U(1)_{Z'}$ charge. The possible interactions of the uniton with the Standard Model states are:

$$\begin{aligned}
& LQ\bar{D}, u_L^c e_L^c D, QQD, u_L^c d_L^c \bar{D}, d_L^c N_L^c D, \\
& QDh \\
& \bar{D}\bar{D}u_L^c
\end{aligned}$$

The terms in the first line above are of the form $b_i b_j D \phi^{n-3}$. Where b_i and b_j are states from the sectors b_1 , b_2 and b_3 , D is the uniton and ϕ^{n-3} is a string of Standard Model singlets which insures that a given term satisfies all the string selection rules. because of the “fractional” $U(1)_{Z'}$ charge of the uniton all the terms above break

$U(1)_{Z'}$. Therefore, the string ϕ^{n-3} should contain “fractional” $U(1)_{Z'}$ charge. In this model the only Standard Model singlets with “fractional” $U(1)_{Z'}$ charge are triplets of $SU(3)_H$. The same is true for the interaction terms from the second and third lines. Therefore, if we assume that the $SU(3)_H$ gauge group is unbroken, then all the interaction terms with the Standard Model states vanish to all orders of nonrenormalizable terms. In this case the uniton is stable.

The uniton is stable and may be a candidate for the dark matter. Under the Standard Model gauge group the uniton has the same charges as a down quark and is strongly interacting. It forms bound meson states with the regular up and down quarks: U^\pm, U^0 . An important question is which of the states, the charged or the neutral, is the lighter state. The mass difference is determined by the current mass difference and the interaction mass difference. The mass difference between the charged and neutral state cannot be calculated reliably because of nonperturbative color interactions. In ref. [30] we studied the mass splitting by using heavy quark effective theory as well as potential models. We argued that with our present understanding of QCD and our present knowledge of the experimental data, it is impossible to conclude which of the mesons is the lighter one. Thus, we argued that at present there exists a window in the parameter space for which $M_{U^\pm} > M_{U^0}$. In this case the uniton may be a good dark matter candidate.

I now discuss the cosmological and astrophysical bounds on the uniton [30]. The uniton is a strongly interacting particle. In the early universe it remains in thermal equilibrium until it becomes non-relativistic. The uniton decouples from the thermal bath when its annihilation rate falls behind the expansion rate of the universe. In the non-relativistic limit, $T/M < 1$, the uniton annihilation rate is given by

$$\Gamma = \langle \sigma |v| \rangle n_{eq} \simeq \frac{\pi N \alpha_s^2}{M^2} n_{eq},$$

where M is the mass of the uniton, α_s is the strong coupling at decoupling, n_{eq} is the number density of the uniton at equilibrium,

$$n_{eq} = g_{\text{eff}} \left(\frac{mT}{2\pi} \right)^{3/2} \exp(-m/T) \quad (\text{non - relativistic}),$$

and N is a summation over all the available annihilation channels and is given by

$$N = \sum_f a_f$$

The amplitudes a_f are obtained by calculating the annihilation cross section of the uniton to all the strongly interacting particles, which include the six flavors of quarks (fig. 1) and squarks (fig. 2) and the gluons (fig. 3) and the gluinos (fig. 4). The final states are taken to be massless, yielding $a = 4/3$ for quarks; $a = 14/27$ for gluons; $a = 2/3$ for squarks and $a = 64/27$ for gluinos. In the calculation of the cross section one has to: “properly” include the ghost contribution to remove the unphysical polarization contribution in the annihilation into gluons (fig. 3).

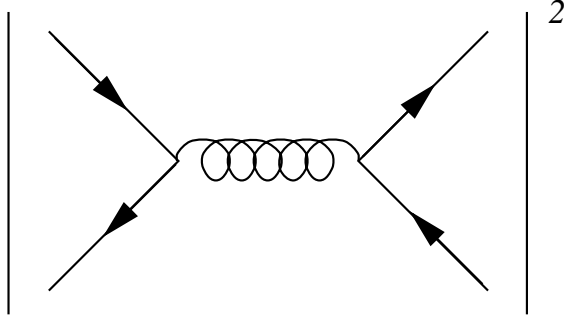


Figure 1: $Q\bar{Q} \longrightarrow q\bar{q}$ annihilation

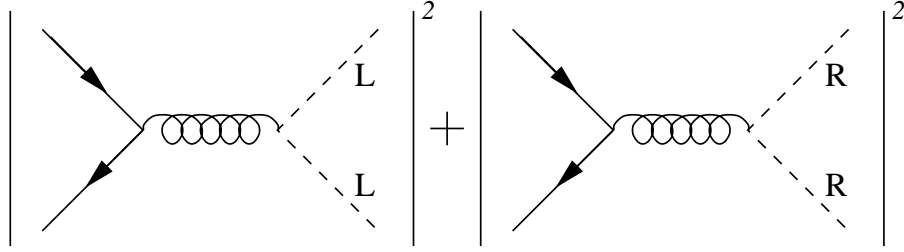


Figure 2: $Q\bar{Q} \longrightarrow \tilde{q}\tilde{q}^*$ annihilation

In the expanding universe, the evolution equation of the particle density in co-moving volume is

$$\frac{dY}{dx} = -\lambda x^{-2}(Y^2 - Y_{eq}^2).$$

Here $Y = n/s$, $x \equiv M/T$ and

$$\lambda = \left. \frac{x \langle \sigma|v| \rangle s}{H} \right|_{x=1} = 0.83 N \alpha_s^2 \frac{g_* s}{\sqrt{g_*}} \frac{m_{pl}}{M}. \quad (3)$$

Here the entropy s is $(2\pi^2/45)g_* m^3 x^{-3}$. The decoupling condition $dY/dx \simeq 0$ gives [31]

$$x_{dec} = \ln[(2+c)\lambda ac] - \frac{1}{2} \ln\{\ln[(2+c)\lambda ac]\},$$

where $a = 0.145(g/g_{*s})$ and c is $Y(T_{dec})/Y_{eq}(T_{dec})$, which is of order one. We approximately estimate the decoupling temperature to be of the form

$$T_{dec} \simeq \frac{M}{\ln(m_{pl}/M)}.$$

The uniton density at the present universe is

$$Y_0 = \frac{3.79 x_{dec}}{\sqrt{g_*} m_{pl} M \langle \sigma|v| \rangle},$$

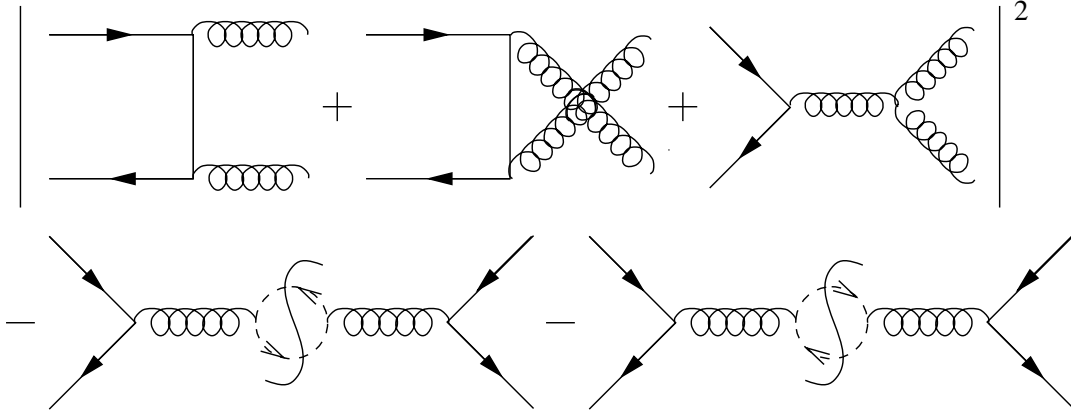


Figure 3: $Q\bar{Q} \rightarrow gg$ annihilation

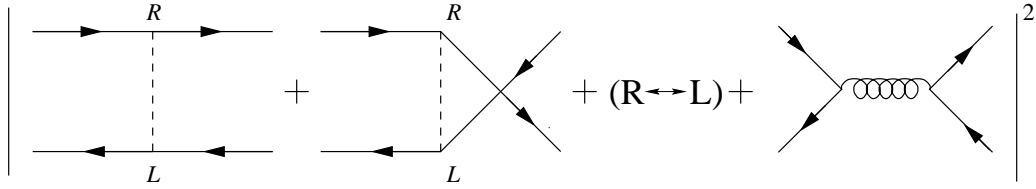


Figure 4: $Q\bar{Q} \rightarrow \tilde{g}\tilde{g}$ annihilation

where we set $g_* = g_{*s}$, since the decoupling temperature is high. Since the relic energy density of a massive decoupled particle is $\rho_0 = M s_0 Y_0$, we can estimate the ratio of energy density to the critical energy density at the present universe to be

$$\Omega_0 h^2 \equiv \frac{\rho h^2}{\rho_c} \simeq 10^9 \frac{\ln(m_{pl}/M) M^2}{N \alpha_s^2 \sqrt{g_*} m_{pl}} \text{GeV}^{-1}.$$

The cosmological data indicates that $0.1 < \Omega h^2 < 1$. Using this condition we get an upper bound on the mass of the uniton

$$M < 10^5 \alpha_s (N \sqrt{g_*} \ln(m_{pl}/M))^{1/2} \text{GeV} \approx 10^5 \text{ GeV}. \quad (4)$$

If we assume the presence of inflation the bound on the uniton mass is modified. In the case of inflation and with $T_R < T_{dec}$, the uniton is diluted and is regenerated after reheating by out-of-equilibrium production. Since the uniton is completely diluted after the inflation, the relic density at the reheating temperature is 0. We can approximate it as

$$\frac{dY}{dx} = \lambda x^{-2} Y_{eq}^2,$$

with $Y_{eq} = 0.145 g/g_* x^{3/2} e^{-x}$. Integrating this relation from the reheating temperature to the present temperature we get

$$Y_0 = \frac{\lambda g^2}{2} \left(\frac{0.145}{g_*} \right)^2 \left(x_r + \frac{1}{2} \right) e^{-2x_r},$$

where $x_r \equiv M/T_R$. T_R is the reheating temperature, and

$$\Omega_0 h^2 \simeq 9 \times 10^3 N \alpha_s^2 g^2 \frac{m_{pl}}{M} \left(\frac{200}{g_*} \right)^{1.5} \left(x_r + \frac{1}{2} \right) e^{-2x_r}.$$

We can estimate the bound on the mass,

$$M > T_R \left[25 + \ln(\sqrt{M/T_R}) \right]. \quad (5)$$

In this case the bound on the unton mass depends on T_R . It is noted that there are three windows (in the parameter space M/σ_p , with σ_p denoting the scattering cross section on protons) for strongly interacting dark matter (such as U_0) [32] which possibly meet our requirements. The first window is in the relatively low mass range ($10 \text{ GeV} < M < 10^4 \text{ GeV}$) and in the range $10^{-24} < \sigma_p < 10^{-20} \text{ cm}^2$. In other two windows it is required to have $10^5 \text{ GeV} < M < 10^7 \text{ GeV}$ and $M > 10^{10} \text{ GeV}$, respectively, assuming a cross section, in both cases, less than 10^{-25} cm^2 . These constraints include bounds from various experiments (such as experiments performed using solid state cosmic-ray detectors and plastic track cosmic-ray detectors) and from cosmological consideration (such as the galactic halo infall rate and the life-time of neutron stars [33]).

The unton is only one example of the exotic Wilsonian matter states in the string derived models. There are several other examples with different properties. For example the models contain color triplets which are weak singlets with weak hypercharge $\pm 1/6$. These color triplets form bound mesonic and hadronic states with the regular up and down quarks which have fractional charge in multiples of $\pm 1/2$. At the same time the same models also contain weak doublets and singlets with fractional electric charge $\pm 1/2$. Thus, these fractionally charged baryons and leptons will form neutral bound states, and the binding energy depends on their masses. The evolution of these states in the early universe is similar to that of the unton since the three gauge couplings are of approximately the same order. The existence, however, of both baryons and leptons with fractional electric charge $\pm 1/2$ offers new scenarios that may weaken the existing limits on fractional charged matter. This is indeed an exciting possibility that merits further investigation.

Another example of a Wilsonian matter state is the state which is a Standard Model singlet with ‘‘fractional’’ charge under the $U(1)_{Z'}$. This state is similar to the right handed neutrino but has half the $U(1)_{Z'}$ of the right handed neutrino. This type of states appear only in the superstring derived standard-like models. For example, in the model of refs. [24] such states appear from the sectors $b_{1,2} + b_3 + \beta \pm \gamma$. In this model these states transform as 3 and $\bar{3}$ of a hidden $SU(3)_H$ gauge group. Their interactions with the Standard Model states vanish to all orders of nonrenormalizable terms if the $SU(3)_H$ is left unbroken. Thus, these states interact with the Standard Model states only by the exchange of the $U(1)_{Z'}$ vector boson. The $U(1)_{Z'}$ gauge symmetry has to be broken somewhere in between the weak scale and the Planck

scale. Thus, depending on the scale of $U(1)_{Z'}$ symmetry breaking there are several possible dark matter scenarios for the Wilsonian singlet.

1. $M \gg M_{Z'}$ without inflation. In this case the Wilsonian singlet can annihilate into the Standard Model fermions and into their superpartners, and into the Z' vector boson and its superpartner. The calculation is similar to the corresponding calculation for the unimon and we obtain a similar limit.

$$M \leq 10^5 \text{ GeV}$$

2. $M \gg M_{Z'}$ with inflation and $T_R > M_{Z'}$. This case again is similar to the case of the unimon with inflation and again we obtain a similar limit

$$M > T_R \left[25 + \frac{1}{2} \ln \left(\frac{M}{T_R} \right) \right].$$

3. $M \ll M_{Z'}$ without inflation \rightarrow relativistic at decoupling

In this case the Wilsonian singlet is a WIMP and it can only annihilate into the Standard Model fermions and their superpartners via the $U(1)_{Z'}$ interactions which are suppressed by $1/M_{Z'}$. Decoupling occurs when the Wilsonian singlet is still relativistic. The number density in the comoving volume is then estimated to be

$$Y_0 \equiv \frac{n_{EQ}}{s} = 0.278 \frac{g_{\text{eff}}}{g_{*s}(T_{\text{dec}})} \simeq 1.2 \times 10^{-3}.$$

where the particle content of the MSSM is assumed and $T_{\text{dec}} > 1 \text{ TeV}$. We then obtain,

$$M < 3 \text{ keV}$$

4. $M \ll M_{Z'}$ with inflation.

In this case the Wilsonian singlet can be heavy. Inflation will dilute the Wilsonian singlet and they will be regenerated after reheating. There are two regions to consider, $T_R < M$ and $T_R > M$. In the first case the W_s decouples when it is non-relativistic. In this limit we obtain a bound on the mass of the Wilsonian singlet which is similar to the previous bounds in an inflationary scenario,

$$M > T_R \left[25 + \frac{1}{2} \ln \left(\frac{M^5}{M_{Z'}^4 T_R} \right) \right], \quad T_R < M.$$

In the relativistic limit ($T_R > M$) the cross section is temperature dependent. In this case we approximate the thermal average of the energy square by

$$\langle s \rangle = 4 \langle E^2 \rangle \simeq \left[\frac{5}{4} \right] 40 T^2.$$

and the cross section is given by is given by

$$\sigma|v| \simeq \frac{8}{3} N_{Z'} \pi \frac{s}{M_{Z'}^4}, \quad \text{if } T_R > M. \quad (6)$$

Again we integrate the Boltzmann equation for the number density in comoving volume from the reheating temperature to the present universe. In this case we obtain a relation between three unknown parameters, $(M, M_{Z'}, T_R)$,

$$M < \frac{M_{Z'}^4}{T_R^3} 6.9 \times 10^{-25} \left(\frac{g_*}{200} \right)^{1.5} \frac{1}{N_{Z'} g_{\text{eff}}^2}, \quad T_R > M.$$

Next, I turn to discuss the Wilsonian matter states in the context of a superstring motivated dynamical SUSY breaking scenario [35]. In the dynamical SUSY breaking scenarios, supersymmetry breaking is generated dynamically at a relatively low scale and is transmitted to the observable sector by the gauge interactions of the Standard Model [36]. Supersymmetry, in these scenarios, is broken nonperturbatively in a hidden sector and the breaking is mediated to the observable sector by a messenger sector. The universality of the Standard Model gauge interactions results in generation blind mass parameters for the supersymmetric scalar spectrum. Consequently, in these scenarios supersymmetric flavor changing neutral currents are naturally suppressed. A crucial assumption in this regard is the absence of interaction terms between the messenger sector and the Standard Model states.

However, as seen above this is precisely what happens in the case of the Wilsonian matter states. Namely, the “fractional” charges of the “Wilsonian matter states result in unbroken local discrete symmetries which forbid the coupling of the Wilsonian states to the Standard Model states. Thus, the Wilsonian matter states are natural candidates for the messenger sector in the dynamical SUSY breaking scenarios.

A superstring motivated dynamical SUSY breaking scenario was recently proposed [35] in the model of ref. [24]. In this model the NS sector produces the generators of the $SU(5)_H \times SU(3)_H \times U(1)^2$ hidden gauge group. The hidden gauge group contains two non-Abelian factors, $SU(5)_H$ and $SU(3)_H$. In this model it is shown [35] that the requirement of a phenomenologically acceptable generation mixing requires that $SU(5)_H$ is broken near the Planck scale while $SU(3)_H$ is left unbroken. In this case the nonperturbative interactions in the hidden $SU(3)$ gauge group may indeed be generated at a relatively low scale, $\Lambda_3 \approx 100$ TeV, in accordance with the low-energy gauge-mediated dynamical SUSY breaking scenarios. As argued in ref. [37] a non-vanishing F -term may be generated in the direction of one of the gauge singlets, ξ_i , due to the hidden matter condensates. The analysis there was done for the model of ref. [6]. However, because of the similarities between the models I assume that a similar F -term can be generated in this model as well. The model of ref. [24] contains two pairs of color triplets of the uniton type, $\{D_1, \bar{D}_1, D_2, \bar{D}_2\}$. In the superpotential we find the couplings $\xi D_1 \bar{D}_1 + \xi D_2 \bar{D}_2$. As shown above these uniton

states do not have superpotential terms with the Standard Model states. Therefore, this uniton states are natural candidates for the messenger sector in the dynamical SUSY breaking scenarios.

The problem of superstring gauge coupling unification motivates a predictive hypothesis with regard to the messenger sector. Dynamical SUSY breaking scenarios in the context of the MSSM require the existence of both color triplets and electroweak doublets in order not to spoil the unification of the gauge couplings. It requires that the messenger sector states fall into complete representations of $SU(5)$ and that the color triplets and electroweak doublets are almost degenerate in mass. These constraints makes the dynamical SUSY breaking scenarios in the context of the MSSM somewhat ad hoc and unattractive. However, in the context of the string derived models, extra color triplets and doublets are in fact required to obtain unification of the gauge couplings at the string unification scale rather than at the MSSM unification scale. Thus, in the context of the string models the existence of the messenger sector is well motivated. Furthermore, in general, the mass scale of the color triplets has to be much lighter than the mass scale of the electroweak doublets. This motivates the hypothesis that the messenger sector consists solely of color triplets. Moreover, as shown previously the uniton dark matter scenario requires that the uniton mass is of the order $M \approx 100$ TeV, which is precisely the mass scale which is required for the uniton to play the role of the messenger sector in dynamical SUSY breaking scenarios. The string of lucky strikes does not end here. For if we assume the existence of color triplets as well as electroweak doublets, then the charged lepton and the color triplets are almost degenerate in mass. In this case both the color triplets and the charged lepton are stable. While the color triplets, as was argued above, can confine to form neutral bound states, the charged leptons cannot. Stable charged leptons may be in contradiction with the existence of neutron stars [33]. This string of lucky coincidences may be more than an accident, and motivates the hypothesis that the messenger sector consists solely of color triplets.

The hypothesis that the messenger sector consists solely of color triplets results in very specific predictions for the supersymmetric spectrum. In the dynamical SUSY breaking scenarios the gaugino masses are obtained by one-loop exchange of the messenger sector states and are given by,

$$M_i(\Lambda) = c_i \frac{\alpha_i(\Lambda)}{4\pi} \Lambda \quad (7)$$

where Λ is the SUSY breaking scale, c_i are coefficients which depend on the messenger sector, and $\alpha_i(\Lambda)$ are the Standard Model coupling constants at the scale Λ . The scalar masses arise from two-loop diagrams and are given by

$$m^2(\Lambda) = 2\Lambda^2 \left\{ C_3 \left[\frac{\alpha_3(\Lambda)}{4\pi} \right]^2 + C_2 \left[\frac{\alpha_2(\Lambda)}{4\pi} \right]^2 + \frac{3}{5} \left(\frac{Y}{2} \right)^2 \left[\frac{\alpha_1(\Lambda)}{4\pi} \right]^2 \right\} \quad (8)$$

where the weak hypercharge has the standard $SO(10)$ normalization $U(1)_Y = 3/5U(1)_1$ and $C_3 = 4/3$ for color triplet scalars and zero for sleptons and $C_2 = 3/4$ for electroweak doublets and zero for singlets.

With the hypothesis that the messenger sector consists only of color triplets

$$M_2 \equiv 0.$$

The chargino mass matrix is given by

$$M_{\tilde{C}} = \begin{pmatrix} \tilde{M}_2 & M_W \sqrt{2} \sin \beta \\ M_W \sqrt{2} \cos \beta & \mu \end{pmatrix}, \quad (9)$$

With this hypothesis \tilde{M}_2 in eq. (9) is equal to zero and β and μ are taken as free parameters. Thus, in this scenario one of the eigenvalues of the chargino mass matrix is predicted to be below the W -boson mass. Imposing the current experimental limits on the supersymmetric spectrum [38], the lightest chargino mass is predicted to be in the range

$$M_{\chi^\pm} \approx 56 - 65 \text{ GeV}.$$

Similarly, from eq. (8) it is seen that in this case the sneutrino is the lightest superparticle.

To conclude the Wilsonian matter states have important cosmological and phenomenological implication. They give rise to natural dark matter candidates whose stability is protected by a local discrete symmetry. This is an important advantage of the Wilsonian dark over other dark matter candidates whose stability relies on the existence of global symmetries. The Wilsonian matter states in the string derived models offer exciting possibilities for confronting string inspired scenarios with experimental data. For example, the hypothesis of dynamical SUSY breaking with color triplets solely predicts a light chargino and will be tested at LEP2. The existence of stable Wilsonian matter states may be tested in dark matter searches and in searches for rare elements.

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