MODEL-INDEPENDENT PROPERTIES AND COSMOLOGICAL IMPLICATIONS OF THE DILATON AND MODULI SECTORS OF 4-D STRINGS

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We show that if there is a realistic 4-d string, the dilaton and moduli supermultiplets will generically acquire a small mass $\sim O(m_{3/2})$, providing the only vacuum-independent evidence of low-energy physics in string theory beyond the supersymmetric standard model. The only assumptions behind this result are (i) softly broken supersymmetry at low energies with zero cosmological constant, (ii) these particles interact with gravitational strength and the scalar components have a flat potential in perturbation theory, which are well-known properties of string theories. (iii) They acquire a vev of the order of the Planck scale (as required for the correct value of the gauge coupling constants and the expected compactification scale) after supersymmetry gets broken. We explore the cosmological implications of these particles. Similar to the gravitino, the fermionic states may overclose the Universe if they are stable or destroy nucleosynthesis if they decay unless their masses belong to a certain range or inflation dilutes them. For the scalar states it is known that the problem cannot be entirely solved by inflation, since oscillations around the minimum of the potential, rather than thermal production, are the main source for their energy and can lead to a huge entropy generation at late times. We discuss some possible ways to alleviate this entropy problem, that favour low-temperature baryogenesis, and also comment on the possible role of these particles as dark matter candidates or as sources of the baryon asymmetry through their decay.

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1 Introduction

One of the major problems facing string theories is the lack of predictive power for low-energy physics. This problem is mostly due to the immense number of consistent supersymmetric vacua that can be constructed. Concentrating on model-independent properties of these vacua, it is well known that all of them include in the spectrum, besides the gravity sector, a dilaton ($S$) multiplet. The tree-level couplings of this field are well understood and are independent of the vacuum. In particular it couples to the gauge kinetic terms, thus its vev gives the gauge coupling constant at the string scale. It is also known to have vanishing potential to all orders in perturbation theory and to couple to every other light field only by non-renormalizable interactions.

Another generic property of 4-d string vacua is that each of them belongs to classes of models labelled by continuous parameters called moduli ($T_i$). These parameters are fields in the effective field theory whose potential is flat. The standard moduli fields characterize a particular compactification, e.g. the size and shape of a Calabi–Yau manifold or toroidal orbifold. But they are also known to exist even in models without a clear geometric interpretation like asymmetric orbifolds \footnote{Examples of models without moduli have been constructed, but none of them corresponds to a 4-d string vacuum \footnote{Strictly speaking, not all of the moduli fields have to take vev's of order 1. Our result applies only to those that have a vev of that order, having geometrical interpretation or not. The logical alternatives are that non-perturbative effects do not lift the degeneracy completely, thus leaving some flat directions and}}. Even though their existence is generic, their couplings are not as model-independent as those of the dilaton, except for having vanishing potential and non-renormalizable couplings to the other light fields.

Independent of the particular mechanism of supersymmetry breaking, any realistic string model is expected to lead to a low-energy theory with softly broken supersymmetry at low scale in order to solve the hierarchy problem. Experimental constraints also impose that the cosmological constant is essentially zero and that the gauge coupling constant at the unification scale is of order one. Also the size and shape (moduli) of the compact dimensions are expected to be of order the Planck scale.

These phenomenological requirements, together with the assumption that the supersymmetry breaking mechanism is responsible for fixing those vev’s, are enough to prove a very simple but strong result in string theory: The scalar and fermionic components of the dilaton and moduli superfields can take masses at most of the order of the gravitino mass, which is expected to be $\leq 1$ TeV in order to solve the hierarchy problem. This means that any realistic string theory (as defined above) should have at low-energies not only the supersymmetric standard model spectrum, but also the dilaton and moduli superfields \footnote{Strictly speaking, not all of the moduli fields have to take vev's of order 1. Our result applies only to those that have a vev of that order, having geometrical interpretation or not. The logical alternatives are that non-perturbative effects do not lift the degeneracy completely, thus leaving some flat directions and}. This fact is actually not surprising since before supersymmetry breaking these fields have flat potentials (and are therefore massless). The interesting observation
is that this can be claimed in a completely vacuum-independent way and is therefore a 'prediction' of any model satisfying the above hypothesis, which are almost imperative in any possible realistic string model (note, however, that in ordinary field theories the $S$ and $T_i$ fields are not necessarily present).

It is then natural to ask what the physical implications of these light particles are. Since all of their couplings to the observable sector are suppressed by powers of the Planck scale, it is perhaps worthless to look for direct phenomenological consequences. But they can play an important role in cosmology as very weakly interactive light particles do. In general, particles with these properties and masses in the TeV range are problematic for cosmology [3, 4, 5]. If they are stable, they can overclose the Universe, and if they decay late, they can dilute or alter the light elements abundance. For the fermionic fields this problem could be solved if there is a period of inflation, which would dilute their initial number density making them harmless. For scalar fields, the problem is more serious since after inflation they are usually sitting far from their zero-temperature minimum and they store a large amount of energy in the associated oscillations.

These problems have been noticed before in the context of particular supergravity [5] and superstring [6] models, but according to the arguments above they are truly generic in string theory. We examine alternative mechanisms to solve these problems, finding that some of them could be naturally implemented in phenomenologically interesting string models. We find that the entropy problem is significantly alleviated if the inflaton decays at late times giving a small reheating temperature, and that the cosmological bounds then favour low-energy scenarios of baryogenesis (e.g. at the electroweak scale or below). Finally, we point out that the dilaton and moduli superfields may provide a dark matter candidate and, in $R$-parity violating models, they could actually trigger low-temperature baryogenesis [7].

## 2 Masses of the Dilaton and Moduli Superfields

Following the standard notation [8], the scalar potential for a $d = 4, N = 1$ supergravity theory can be conveniently written in $M_P$ units as

$$V = \left[ F_k F^j \mathcal{G}^k_j - 3e^{\mathcal{G}} \right] + \frac{1}{2} f_{\alpha \beta} D^\alpha D^\beta ,$$

where

$$D^\alpha = -\mathcal{G}_i T_i^{\alpha j} z_j , \quad F^j = e^{\mathcal{G}/2} (\mathcal{G}^{-1})^j_l \mathcal{G}^l$$

then massless fields or that the vev’s are smaller (larger) than the Planck scale for which the corresponding field is heavier (lighter). For an example in this class see [2].
Here $z_j$ denote the chiral fields, $T^\alpha$ are the gauge group generators, $G = K + \log |W|^2$ where $K$ is the Kähler potential and $W$ the superpotential, $f_{\alpha\beta}$ are the gauge kinetic functions (in superstrings $f_{\alpha\beta} = S\delta_{\alpha\beta}$ up to threshold corrections), and $G^i$ ($G_j$) denotes $\partial G / \partial z_i^* (\partial G / \partial z_j^*)$. $F_i$ and $D^\alpha$ are the $F$ and $D$ auxiliary fields. Supersymmetry is broken if at least one of the $F_i$ or $D^\alpha$ fields takes a vev. The scale of supersymmetry breakdown, $M_S$, is usually defined as $M^2_S = \langle F \rangle$ or $\langle D \rangle$ (depending on the type of breaking).

Let us now see why the masses of the dilaton and moduli fields are of order $m_{3/2}$. The assumptions to show this are the following

i) The $S$ and $T_i$ fields must take vev’s of the order of the Planck scale $M_P$. This is certainly mandatory for the $S$ field, whose vev gives (up to threshold corrections) the gauge coupling constant at the unifying string scale ($\langle \text{Re } S \rangle = g^{-2}_{\text{string}}$). If the model can be interpreted as coming from higher dimensional compactification, the size and shape of these compact dimensions are parametrized by $\langle T_i \rangle$ and are also expected to be of the order of the Planck scale.

ii) Supersymmetry is softly broken at a low scale with vanishing cosmological constant. More precisely, the mechanism for supersymmetry breakdown should yield a gravitino mass $m_{3/2} \lesssim 1$ TeV $\ll M_P$ (without unnatural fine-tunings), as is required from phenomenological reasons.

We emphasize here that we do not assume any particular scenario for achieving (i) and (ii). We just assume that if the observable world is the low-energy limit of a string theory, the previous conditions must be fulfilled. Besides points (i) and (ii), the following well-known property of string perturbation theory should be added:

iii) The fields $S$ and $T_i$ interact with gravitational strength and have a flat potential at string tree-level which remains flat to all orders of perturbation theory. This means that the only source of masses will come from non-perturbative effects which on the other hand should be responsible also for the supersymmetry breaking process. Hence, we will assume these fields to be massless in the absence of supersymmetry breaking.

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4It is conceivable that non-perturbative effects which trigger masses for the $S$ and $T_i$ fields without breaking supersymmetry could exist. However, supersymmetry has to be broken in any realistic model and non-perturbative effects would be responsible for that. The way to relax this assumption is the existence of a hierarchy of non-perturbative effects: Planck scale effects which fix all the vev’s without breaking supersymmetry and low-energy effects which break supersymmetry, as suggested for the dilaton in [10].
Once the assumptions are written, it is straightforward to find the order of magnitude of the masses. Roughly speaking, (see for instance [11]) after supersymmetry breaking, the mass splittings in a given supermultiplet are given by the product of the square of the supersymmetry breaking scale times the coupling strength between that field and the goldstino ($\sim M_S^2/M_P$) which is just the gravitino mass. Let us see this more explicitly\footnote{Notice that S and T$_i$ share the properties of the hidden sector fields in the general analysis of reference [12] in the context of N = 1 supergravity, although that analysis includes also cases such as linear superpotentials of the Polonyi type which do not hold for the moduli and dilaton fields. It is straightforward to prove our result using that formalism also.}

For simplicity of notation, let us collectively denote $\phi$ the scalar components of the $S$ and $T_i$ fields and $z$ the observable fields. Since the vev’s of the $\phi$ fields are of order the Planck scale it is convenient to define the dimensionless fields $\chi \equiv \phi/M_P$. The Kähler potential can be written as

$$K = M_P^2 K_0(\chi) + K_1(\chi) z^* z + \cdots$$

where $K_0$ and $K_1$ are arbitrary real functions. The important point here is that $K$ is of order $M_P^2$ and not higher, since the $\phi$ fields have only gravitational strength interactions. Hence $G_k^l = K_k^l \lesssim O(1)$ in (1).

In flat space the gravitino mass is given by

$$m_{3/2}^2 = e^G = e^K |W|^2$$

(3)

From (1), (3) and the cancellation of the cosmological constant, it is clear that

$$m_{3/2}^2 \sim M_S^4 M_P^2$$

(4)

Hence $V$, as given in eq.(1), is a sum of terms

$$V = \sum_a V_a$$

(5)

with $\langle V_a \rangle \lesssim m_{3/2}^2 M_P^2$ ($= M_S^4$). Some $V_a$ can be $\langle V_a \rangle = 0$ (or $\ll m_{3/2}^2 M_P^2$), but never $\langle V_a \rangle \gg m_{3/2}^2 M_P^2$, since this would imply a fine-tuning in order to cancel the cosmological constant. Of course, $\langle V_a \rangle = 0$ in the absence of supersymmetry breaking.

The $\phi$ masses are then given by

$$m_{\phi}^2 = \langle \partial^2 V/\partial\phi^2 \rangle = \sum_a \langle \partial^2 V_a/\partial\phi^2 \rangle = \frac{1}{M_P^2} \sum_a \langle \partial^2 V_a/\partial\chi^2 \rangle \sim M_P^{-2} m_{3/2}^2 M_P^2 \sim O(m_{3/2}^2)$$

(6)

The additional factor coming from the normalization of the kinetic term depends on $G_k^l$ and is $O(1)$. Notice that if the breaking of supersymmetry were explicit rather than spontaneous the same conclusion would hold as long as the terms in the potential are of order $M_S^4$.\footnote{Notice that S and T$_i$ share the properties of the hidden sector fields in the general analysis of reference [12] in the context of N = 1 supergravity, although that analysis includes also cases such as linear superpotentials of the Polonyi type which do not hold for the moduli and dilaton fields. It is straightforward to prove our result using that formalism also.}
Dilatino and Modulino Masses

As usual, the fermion component of the field whose $F$ (or $D$) auxiliary field takes a vev is the goldstino, which in principle is massless but is eaten by the gravitin through the super-Higgs effect. This field could perfectly be the dilatino $\tilde{S}$ or one of the modulinos $\tilde{T}_i$, or perhaps a certain combination of them. For the remaining fermionic components, $\psi_i$ of the chiral superfields $z_i$ the mass term is $[M_\psi]^{ij}\hat{\psi}_{Li}\psi_{Lj}$, where the fermionic mass matrix $M_\psi$ can be written as $[M_\psi]^{ij} = \sum_{n=1}^{4}[M_\psi^{(n)}]^{ij}$, with

\[
[M_\psi^{(1)}]^{ij} = -e^{K/2}|W|\left\{K_{ij} + \frac{1}{3}K_iK_j\right\},
\]

\[
[M_\psi^{(2)}]^{ij} = -e^{K/2}|W|\left\{\frac{K_iW_j + K_jW_i}{3W} - \frac{2W_iW_j}{3W^2}\right\},
\]

\[
[M_\psi^{(3)}]^{ij} = -e^{K/2}\sqrt{\frac{W^*}{W}}W^{ij},
\]

\[
[M_\psi^{(4)}]^{ij} = e^{G/2}G^l(G^{-1})^k\mathcal{G}_k^{ij}
\]

(7)

where $W_i = \partial W/\partial z_i$, $K^i = \partial K/\partial z_i$. Notice that the canonically normalized fermion fields $\hat{\psi}_{Ln}$ are given by $\psi_{Lj} = U_{jn}\Lambda_{nn}\hat{\psi}_{Ln}$, where $U$ is the unitary matrix diagonalizing the $K_j^i$ matrix and $\Lambda = diag(\lambda_n^{-1/2})$ with $\lambda_n$ the corresponding eigenvalues. Then the mass matrix for the $\hat{\psi}_{Ln}$ fields is given by

\[
\hat{M} = \Lambda U^+ M(\Lambda U^+)^T.
\]

(8)

Since in general the non-vanishing $K_j^i$ derivatives are $\lesssim O(1)$, $M$ and $\hat{M}$ are of the same order of magnitude. Of course, we are interested in the case where the $\psi_i, \psi_j$ fields are the dilatino $\tilde{S}$ and the modulinos $\tilde{T}_i$. Let us then analyse the magnitude of each $[M_\psi^{(n)}]^{ij}$ term individually.

The first term $[M_\psi^{(1)}]^{ij}$ is clearly of order $m_{3/2} = e^{K/2}|W|$ since the possible non-vanishing $K^{ij}, K^i, K^j$ derivatives are in general $\lesssim O(1)$ in Planck units. (Notice that for normal observable fields with zero vev’s these quantities are usually vanishing.) The second term $[M_\psi^{(2)}]_{ij}$ is $O(m_{3/2})$ for analogous reasons together with the fact that $W_i/W \lesssim O(1)$ in Planck units. To see this, note that an F-term, $F^k$, can be written (see eq.(3)) as

\[
F^k = e^{K/2}|W|(K^{-1})^k_{\ j}\left(K^j + \frac{W^j}{W}\right)
\]

(9)

where the derivatives of $K$ are $\lesssim O(1)$ and $\langle F_k \rangle \lesssim M_pm_{3/2}$. Thus $W_i/W \lesssim O(1)$. The third term $[M_\psi^{(3)}]^{ij}$ contains the contribution to the fermion masses which in principle is not triggered by supersymmetry breaking (notice that this term is not proportional to $m_{3/2}$). On the other hand, in the absence of supersymmetry breaking we know that
m_\tilde{S} = m_\tilde{f}_i = 0$. Consequently, for the dilaton and the moduli, $W^{ij}$ can only be different from zero as an effect of supersymmetry breaking. Following a reasoning similar to that of estimating the value of $\partial^2 V / \partial \phi^2$, it is easy to see that $W^{ij} = O(W) M_F^{-2} = O(m_{3/2})$, and thus $[M^{(3)}_W]^{ij} = O(m_{3/2})$. Finally, $[M^{(4)}_W]^{ij}$ can be written as $[M^{(4)}_W]^{ij} = F^k K^{ij}$, where $K^{ij} \sim O(1)$ in Planck units and $\langle F^k \rangle \sim M_P m_{3/2}$. Therefore, $[M^{(4)}_W]^{ij} = O(m_{3/2})$.

To summarize, under the assumptions (i)–(iii) the masses of the scalar and fermionic components of the dilaton and the moduli are $O(m_{3/2})$.

### A Class of Models

We illustrate the previous results now in a class of orbifold models in which all of the moduli are frozen except for the overall one ($T$). The Kähler potential in Planck scale units is

$$K(S, T, \Phi_\alpha) = K_0(S, T) + \sum_\alpha K^\alpha_1(T)|\Phi_\alpha|^2 + O(|\Phi_\alpha|^4)$$

(10)

where

$$K_0(S, T) = -\log Y - 3 \log(T_R)$$

$$K^\alpha_1 = (T_R)^{n_\alpha}$$

(11)

Here $Y = S_R + \delta \log(T_R)$, where $\phi_R = \phi + \bar{\phi}$, $\phi = S, T$. Also $\delta \equiv \frac{\delta G_S}{4\pi}$ is a model dependent constant coming from the one-loop anomaly cancelling Green–Schwarz counterterms and $n_\alpha$ represent the modular weights of the given matter superfield and depend on the sector (twisted or untwisted) of the corresponding field. For the superpotential we will take an arbitrary function $W(S, T, \Phi_\alpha) = W^{\text{pert}}(T, \Phi_\alpha) + W^{\text{np}}(S, T, \Phi_\alpha)$ where, as we know, the perturbative part does not depend on $S$ and vanishes with $\Phi_\alpha$, and the non-perturbative part is unknown. To write the explicit expression for the fermion masses, we find it convenient to work with a general $W$ which for vanishing matter fields will be only the non-perturbative part $W = W^{\text{np}}(S, T)$.

The mass matrix (7) for the fermionic components of the $S$ and $T$ fields is then given by

$$M^{SS} = m_{3/2} \left[ \frac{1}{Y^2} - \frac{2}{Y} \frac{W^S}{W} - (C^{SS} + \frac{1}{3} B^2) \right]$$

$$M^{TS} = -m_{3/2} \left[ C^{ST} + \frac{B}{3} D + \frac{\delta B}{Y(T_R)} (1 + \frac{3}{A}) + \frac{\delta W^T}{Y^2 A W} \right]$$

$$M^{TT} = m_{3/2} \left[ E \left( 1 - \frac{2}{(T_R)^2} \frac{W^T}{A W} \right) - \left( C^{TT} + \frac{D^2}{3} \right) + \frac{\delta B}{(T_R)^2} (1 - 6 \frac{\delta}{A Y}) \right]$$

(12)

\(^6\)Notice that this justifies the procedure of ref. \[^{14}\] where intermediate mass hidden fields are integrated out whereas the moduli fields are kept in the low-energy theory.
where $A \equiv 3 + \delta/Y$, $B \equiv W^S/W - 1/Y$, $C^{I,J} \equiv W^{I,J}/W - W^I W^J/W^2$, $D \equiv W^T/W - A/(T_R)$ and $E = \frac{A}{(T_R)^2} (1 + \frac{\delta^2}{A Y})$. For the correctly normalized fields $\tilde{\phi}_1, \tilde{\phi}_2$ the mass matrix is given by (8), where in our case

$$U^+ = \left( \begin{array}{cc} \frac{K_S^T}{\sqrt{(\lambda_1 - K_S^2) + (K_S^T)^2}} & \frac{\lambda_1 - K_S^2}{\sqrt{(\lambda_1 - K_S^2) + (K_S^T)^2}} \\ \frac{\lambda_2 - K_T^2}{\sqrt{(\lambda_2 - K_T^2) + (K_T^T)^2}} & \frac{\lambda_1 - K_T^2}{\sqrt{(\lambda_2 - K_T^2) + (K_T^T)^2}} \end{array} \right)$$

and $\lambda_{1,2} = \frac{1}{2} \left( \frac{1}{Y^2} + E \right) \pm \sqrt{\frac{1}{Y^2} + E - \frac{2}{Y^2} (A - \frac{\delta^2}{A Y})}$. It is clear that the masses are of $O(m_{3/2})$ and the precise value will depend on the non-perturbative superpotential. Notice that if $\delta = 0$ and $W_{np} \neq W_{np}(T)$, then $M^{TT} = 0$. For the bosonic fields we have also calculated the masses and found them of the same order. For simplicity we will present here only the final result in the limit $\delta = 0$ and for the particular case that the superpotential factorizes $W(S, T) = \Omega(S) \Gamma(T)$. In this case the physically normalized mass matrix does not have off-diagonal $(S, T)$ components and its eigenvalues are, for the dilaton field:

$$M_{S^\pm}^2 = m_{3/2}^2 S^2_R \left( |\frac{2 \Omega^S}{\Omega} + S_R C^{SS}|^2 + |S_R C^{SS}|^2 + (2 C^{SS} S_R B + h.c.) \right)$$

$$\pm 2 |S_R B (2 C^{SS} + S_R C^{SS}) + \frac{2 \Omega^S}{\Omega} (\frac{\Omega^S}{\Omega} + S_R C^{SS})|,$n

wheras for the modulus we find:

$$M_{T^\pm}^2 = m_{3/2}^2 \frac{T_R^2}{9} \left( |\frac{2 \Gamma^T}{\Gamma} + T_R C^{TT}|^2 + |T_R C^{TT}|^2 - (2 C^{TT} D + h.c.) \right)$$

$$\pm 2 |\frac{2 \Gamma^T}{\Gamma} + T_R C^{TT} - D (2 C^{TT} + C^{TT})|.$$n

Notice that unlike the observable fields for which usually one member of the supermultiplet acquires a mass of $O(m_{3/2})$ and the other vanishes (since their masses are protected by gauge invariance) for these fields both members of the supermultiplet have similar masses of $O(m_{3/2})$, unless there is a cancellation in any of the equations above or, as we said before, the non-perturbative effects do not lift the flatness of the potentials and some of the fields remain massless. This is the case if the superpotential is independent of any of the fields, as can be easily verified.

\footnote{An apparent discrepancy with the calculation of \cite{15} in this limit is due to the fact that we are using the mass matrix after the super-Higgs effect, which is reflected in the $1/3$ factors in eqs.\cite{15}.}
3 Cosmological Implications

It is natural to ask what the physical implications of the lightness of the dilaton and moduli sector are. First of all it is necessary to note that the couplings of these particles to the observable sector are suppressed by powers of the Planck scale, so it is perhaps useless to look for direct phenomenological consequences unless they happen to be very light and could mediate long range forces [10]. However, the dilaton and the moduli can play an important role in cosmology as very weakly interacting light particles do.

Let us first consider the fermionic components ($\tilde{S}$ and $\tilde{T}_i$) of these fields. Their cosmology is very similar to that of the gravitino [4, 17]. In particular, if they were stable and in the absence of inflation, they would overclose the Universe unless their masses are in the range $m_{\tilde{S}}, m_{\tilde{T}_i} \lesssim 1$ keV [4]. A gravitino mass of this order is an interesting possibility which has received some attention recently [18] but most standard phenomenological scenarios assume $m_{3/2} \sim \text{TeV}$, which seems to be in conflict with cosmology. Nevertheless, these particles are actually expected to decay (unless they were the lightest supersymmetric particles) with a rate $\Gamma_{\tilde{S}} \sim m_{\tilde{S}}^3 M_p^2 (\phi = S, T_i)$. They should decay before nucleosynthesis, otherwise their decay products give unacceptable alteration of the primordial $^4\text{He}$ and D abundances. This imposes a lower bound on their masses [4, 17]

$$m_{\tilde{S}}, m_{\tilde{T}_i} > O(10) \text{ TeV}$$

This bound is not very comfortable, but may be fulfilled since $m_{\tilde{S}}, m_{\tilde{T}_i} = O(m_{3/2})$ includes the possibility of one or two orders of magnitude higher than $m_{3/2}$.

The previous cosmological problems associated with light, very weakly interacting (but decaying) fermionic fields are obviated if the Universe undergoes a period of inflation that dilutes them, provided that the reheating temperature after inflation satisfies

$$T_{RH} \lesssim 10^8 \left( \frac{100 \text{ GeV}}{m_{3/2}} \right) \text{ GeV}$$

in order for them not to be regenerated in large amounts. In this case, low-temperature mechanisms for baryogenesis, which have recently received much attention, become strongly favoured (for a recent review see [19]). On the other hand, it is worth mentioning a mechanism [7] in which baryogenesis is driven by the late decay of a particle such as the gravitino or the axino. It is based on the possible existence of baryon-number-violating terms in the superpotential of the form $u_L^c d_L^c d_L^c$ (for the third generation) and exploits the sources of CP violation appearing in supersymmetric models. This mechanism can also be implemented for the dilatino and modulino fields, since they have axino-like couplings
with gravitational strength, and could work under condition (16) but requires a reheating temperature \( O(M_P) \) or no inflation at all (as is the case for the gravitino).

Let us now turn to the scalar (dilaton and moduli) fields. They present much more severe problems than their supersymmetric partners since in general, when they start their relevant cosmological evolution, they are far from their zero-temperature minimum and the excessive energy associated to their oscillations around the minimum tend to be problematic. This well known situation \([3]\) has been called in the literature the ‘Polonyi problem’ or the ‘entropy crisis’ and has been noticed in the context of particular supergravity or superstring models \([3]\). However, according to the arguments of the previous section, this is a truly generic problem in string theory.

The dilaton and moduli fields are expected to be initially shifted from their zero-temperature minimum either due to the effect of thermal fluctuations or of quantum fluctuations \([20]\) during inflation or also due to the fact that their coupling to the inflaton will generally modify, during inflation, the value corresponding to the minimum of the potential \([21]\). The shift produced by these effects may even be as large as \( O(M_P) \) but its magnitude depends on the particular form of the scalar potential and should be estimated case by case.

The evolution of the \( \phi \) field (canonically normalized dilaton or moduli) is determined by

\[
\ddot{\phi} + (3H + \Gamma_{\phi})\dot{\phi} + \frac{\partial V}{\partial \phi} = 0
\]

(18)

where \( \Gamma_{\phi} \) is the \( \phi \) width \( (\sim m_{\phi}^2 M_P^{-2}) \) and \( H \) is the Hubble constant. Although \( V \) is not known, if we consider just oscillations around the minimum taking \( V \sim \frac{1}{2} m_{\phi}^2 \phi^2 \) one sees that the friction dominates until the time \( t_{in} \sim m_{\phi}^{-1} (T_{in} \sim \sqrt{m_{\phi} M_P} \) in the case of radiation domination), after which the field oscillates and the density evolves as \( T^3 \) \([3]\). Hence, taking as \( \phi_{in} \) the value of the shift in the fields at the end of inflation (and neglecting its evolution up to \( t_{in} \)), the initial density \( \rho_{\phi}(T_{in}) \sim m_{\phi}^2 \phi_{in}^2 \) increases with respect to the density in radiation as

\[
\frac{\rho_{\phi}(T)}{\rho_{rad}(T)} = \frac{\rho_{\phi}(T_{in})}{\rho_{rad}(T_{in})} \frac{T_{in}}{T}
\]

(19)

If \( \phi \) were stable \( (\Gamma_{\phi} \sim 0) \), the constraint that \( \phi \) does not overclose the Universe today imposes

\[
\phi_{in} \leq 10^{-10} M_P (m_{\phi}/100 \text{GeV})^{-1/4},
\]

(20)

which is much smaller than the typical shifts in the fields expected due to the mechanisms mentioned above.
However, $\phi$ will generally decay. For instance, the dilaton $S$ is coupled to all the gauge bosons of the theory through the term $Re \int F_{\mu\nu}F^{\mu\nu} + Im \int F_{\mu\nu}\tilde{F}^{\mu\nu}$, where $f = S$ and threshold corrections. Also, the moduli fields appear in the threshold corrections to $f$ and/or in the Yukawa couplings between the charged fields. Therefore, $\rho_\phi$ will eventually be converted in radiation. As long as $\phi_{in} \geq 10^{-8} M_P \sqrt{m_\phi/\text{TeV}}$ the field $\phi$ will dominate the energy density at the moment it decays, that will correspond to a temperature $T_D \approx m_\phi^{11/6} \phi_{in}^{-2/3} M_P^{-1/6}$ and will reheat the Universe to $T_{RH}$ given by

$$T_{RH} \sim m_\phi^{3/2} M_P^{-1/2}.$$  \hspace{1cm} (21)

(Notice that $T_{RH}$ is essentially independent of $\phi_{in}$, provided the Universe has been $\phi$-dominated and the total density is $\sim \rho_c$.) Now, the decay products of $\phi$ will destroy the $^4\text{He}$ and $D$ nuclei, and thus the successful nucleosynthesis predictions, unless $T_{RH} > 1$ MeV, since then the nucleosynthesis process can be re-created. This implies

$$m_\phi > O(10) \text{ TeV} \hspace{1cm} (22)$$

This bound is, of course, similar to that of eq.(16) for the dilatino and modulinos (and for the gravitino), but however it cannot be escaped with the help of inflation. On the other hand, the entropy increase when $\phi$ decays, $\Delta$, is given by

$$\Delta = \left( \frac{T_{RH}}{T_D} \right)^3 \sim \frac{\phi_{in}^2}{m_\phi M_P}.$$  \hspace{1cm} (23)

If $\Delta$ is very large, it would erase any pre-existing baryon asymmetry. Denoting $\Delta_{max}$ the maximum tolerable entropy production (this will depend on the specific baryogenesis mechanism, but $\Delta_{max} \sim 10^5$ might be acceptable) we obtain the constraint

$$\phi_{in}^2 < \Delta_{max} M_P m_\phi \hspace{1cm} (24)$$

We want to note here that a possible way out of this dilemma would be to have the baryogenesis generated by the $\phi$ decays, that would take place just before nucleosynthesis. This could be done for instance if the $\phi$ decay into gaugino pairs is allowed (which seems reasonable in view of (22)) and implementing then a mechanism similar to that of ref.[7] just discussed. Unlike the baryogenesis scenarios based on decays of fermions (gravitino, dilatino, etc.) those with the $\phi$ decays would not imply any severe constraint on the reheating after inflation since it is easier to dominate the energy density at the moment of decay.

This situation can however be ameliorated: the previous discussion assumed that at the moment in which $\phi$ starts to oscillate, the inflaton has already decayed and produced the main reheating of the Universe. However, this is not necessary, especially in view
of the many scenarios of low-temperature baryogenesis being considered nowadays. If
the inflaton $\varphi$ were instead to decay at $t_{RH} \gg t_{in}$, the situation will significantly change
due to the fact that between those times both the inflaton and $\phi$ will oscillate and their
energies will both evolve as $a^{-3}$, where $a$ is the scale factor. Hence, using the fact that,
at $t_{in}$, $H = m_{\phi}$ so that $\rho_{\varphi} = 3M_{P}^{2}m_{\phi}^{2}/(8\pi)$, we have
\[
\frac{\rho_{\phi}}{\rho_{\varphi}}(t_{RH}) \simeq \frac{8\pi \phi_{in}^{2}}{3M_{P}^{2}} .
\] (25)
Only after the inflaton energy is converted into radiation ($\rho_{\varphi}(t_{RH}) \sim T_{RH}^{4}$) the relative
contribution of $\phi$ to the density starts to increase as
\[
\frac{\rho_{\phi}}{\rho_{\varphi}}(T) \simeq \frac{8\pi \phi_{in}^{2} T_{RH}}{3M_{P}^{2} T} .
\] (26)
and, for a stable $\phi$, the condition that at $T = 3$ K the Universe is not overclosed becomes
\[
\phi_{in} \leq \sqrt{\frac{1 \text{TeV}}{T_{RH}}} 10^{-6} M_{P} ,
\] (27)
that is much less severe than the previous bound if the reheating temperature is of the
order of the weak scale.

In the more plausible situation in which $\phi$ decays, by a similar reasoning we obtain
that $\phi$ will be dominating the density of the Universe at the decay time only if
\[
\phi_{in} \geq \left( \frac{m_{\phi}}{10 \text{ TeV}} \right)^{3/4} \left( \frac{T_{RH}}{100 \text{ GeV}} \right)^{-1/2} 10^{-3} M_{P} ,
\] (28)
and in that case the Universe has a further reheating to $T'_{RH} \simeq \sqrt{m_{\phi}^{2}/M}$ with an increase
in the entropy by a fraction
\[
\Delta \simeq \frac{\phi_{in}^{2} T_{RH}}{(M_{P} m_{\phi})^{3/2}} .
\] (29)
The analogue of eq. (24) is $\phi_{in}^{2} < \Delta_{max}(M_{P} m_{\phi})^{3/2}/T_{RH}$ relaxing the bounds for small
enough values of $T_{RH}$. For the indicative values used before (remember that $m_{\phi} > 10$
TeV eliminates the problems associated to nucleosynthesis and $T_{RH} \sim 100$ GeV could
allow for electroweak baryogenesis to occur) $\Delta$ is still not unreasonably large even for
$\phi_{in} = O(M_{P})$ ($\Delta \leq 10^{5}$), although clearly for those values of $\phi_{in}$ the approximation of a
quadratic potential goes badly wrong and the detailed form of the potential would have
to be taken into account.

Hence, although the problem associated to the scalar field oscillations is severe, we
have seen that there are ways to alleviate it. Furthermore, in the context of supergravity
models, there are two interesting proposals to solve this ‘entropy crisis’ \cite{21, 22}. The first method \cite{21} is based on a scenario in which supersymmetry is broken through an O’Raifeartaigh superpotential. Then, the scalar field has almost vanishing vev and a rather large mass ($\sim (m_{3/2} M_P)^{1/2}$) at zero-temperature, thus the contributions to the scalar potential during inflation do not change appreciably the position of the minimum. Consequently, the energy stored in the scalar field is very small, avoiding cosmological problems. Besides the fact that this mechanism does not address the problem of quantum fluctuations, it cannot be implemented for the $S$ and $T$ fields since they must have $O(M_P)$ (non-vanishing) zero-temperature vev’s and their masses should be $O(m_{3/2})$.

A further mechanism that has been proposed \cite{22} is to couple the problematic fields to heavy ones that decay promptly into radiation. In this way the coupled evolution of the scalar fields may allow for a transfer of the energy stored in the oscillations to the radiation. This method might be accomodated in the context of string theories for the $S$ and $T$ field since there are heavy fields with mass terms in the perturbative superpotential that, in general, will depend on the moduli fields. Likewise, the factor $e^K$ in front of the scalar potential provides a non-suppressed coupling between the dilaton and the massive fields. This simply reflects the fact that all the physical couplings are proportional to the string coupling constant. There is still an additional source of couplings between the dilaton and massive fields. It is known that if the four-dimensional gauge group of the string has an ‘anomalous’ $U(1)$ factor, a Fayet–Iliopoulos $D$-term is generated in string perturbation theory \cite{23}. The corresponding term in the scalar potential has the form

$$V_{FI} \sim \frac{1}{S_R} \left| \frac{Tr Q^a}{4\pi^2} \frac{1}{S_R} + \sum_i Q^a_i |C_i|^2 \right|^2$$

where $Q^a_i$ are the anomalous charges of the scalar fields $C_i$. Thus, there appear effective mass terms of the form $\sim (S + \bar{S})^{-2} |C_i|^2$. Furthermore, in this scenario there appear masses for the dilatino and Im $S$, through their couplings to the ‘anomalous’ photino and photon respectively, thus the cosmological problems associated with these fields seem to be reduced (although there remains a combination of dilaton and matter fields which remains massless at this level and to which the previous analysis applies). It is quite suggestive that most of the phenomenologically interesting superstring models have been constructed within this framework \cite{24}. For the remaining fields, the possibility that the transfer of energy mentioned previously be efficient would be worth of a detailed analysis, although that is beyond the scope of the present paper.

Let us also mention that another danger coming from the dilaton and moduli fields being initially sat far from the zero-temperature minimum is that they might never fall into it if there is a barrier in between. (This problem has been recently stressed in ref.\cite{25}.) This can typically happen for the real part of the dilaton ($Re S$) if it initially
sits at \( Re \, S \gg Re \, S_o \), since for \( Re \, S \rightarrow \infty \) all the gauge and gravitational interactions are switched off, and thus supersymmetry is restored and \( V \rightarrow 0 \). Therefore it is necessary to assume that the initial value of the generic scalar field, \( \phi_{in} \), is basically in the slope leading to the zero-temperature minimum. A further problem is that for the typical potentials arising from gaugino condensation, the kinetic energy of the dilaton may be so large as to prevent inflation to take place \([25]\), unless some mechanism exist to push \( \phi \) close to the minimum (this problem may be only characteristic of those potentials and not necessarily generic in string theory).

Finally we will mention that the dilaton and moduli-sector fields may also provide candidates for dark matter. As is apparent from the previous discussion, if the fermionic fields \( \tilde{\phi} \) were stable they would close the Universe for \( m_{\tilde{\phi}} \sim 1 \, \text{keV} \) in the absence of inflation, or if the \( \tilde{\phi} \) mass is set to a \( T_{RH} \)-dependent value if there is an inflationary epoch, in a similar way as for the gravitino \([28]\) (for instance \( m_{\phi} \simeq 100 \, \text{GeV} \) for \( T_{RH} \simeq 10^9 \, \text{GeV} \)). If some of the scalar fields \( \phi \) were stable, they would close the Universe if the bounds in eqs. (20) or (27) are saturated. In that case, the missing energy of the Universe would be stored in the coherent oscillations of the \( \phi \)-field in a similar way to the more standard axionic dark matter.

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