Strings at the Intermediate Scale

or

is the Fermi Scale Dual to the Planck Scale?

C.P. Burgess\textsuperscript{a}, L.E. Ibáñez\textsuperscript{b} and F. Quevedo\textsuperscript{c}\footnote{On leave of absence from Instituto de Física, UNAM, México.}

\textsuperscript{a} Physics Department, McGill University, 3600 University St., Montréal, Québec, Canada, H3A 2T8.

\textsuperscript{b} Departamento de Física Teórica C-XI and Instituto de Física Teórica C-XVI, Universidad Autónoma de Madrid, Cantoblanco, 28049 Madrid, Spain.

\textsuperscript{c} D.A.M.T.P., Silver Street, Cambridge, CB3 9EW, England.

Abstract

We show that if the string scale is identified with the intermediate scale, \( M_s = \sqrt{M_W M_{\text{Planck}}} \approx 10^{11} \text{ GeV} \), then the notorious hierarchy, \( M_W / M_{\text{Planck}} \approx 10^{-16} \), can be explained using only \( M_c / M_s \approx 0.01 \sim \alpha_{\text{GUT}} \) as small input parameters, where \( M_c \) is the compactification scale. This is possible for weakly-coupled Type-I open-string vacua if the observed world is assumed to live in an \( N = 1 \) supersymmetric 3-brane sector coupled to a separate, hidden, 3-brane world which breaks supersymmetry, because for such a model \( M_W / M_{\text{Planck}} = \frac{1}{2} \alpha_{\text{GUT}}^2 (M_c / M_s)^6 \). We discuss some of the phenomenological issues presented by such an intermediate-scale string, showing that its benefits include: (i) the possibility of logarithmic gauge-coupling unification of the SM couplings at \( M_s \); (ii) a natural axionic solution to the strong-CP problem with a phenomenologically-acceptable Peccei-Quinn scale; (iii) experimentally-interesting neutrino masses, and more.


1 Introduction

It has recently been realized that the traditional connection between the string scale and the Planck mass, $M_s = \frac{1}{2} \sqrt{\alpha_{GUT} M_{\text{Planck}}}$ (as is found in perturbative heterotic string theory) need not apply to all string vacua [1, 2, 3, 4, 5, 6, 7, 8, 9, 10].

In this paper our purpose is to argue that the string scale is the geometric mean of the weak and Planck scales: $M_s = \sqrt{M_W M_{\text{Planck}}} \sim 10^{11} \text{ GeV}$. Our main argument in favour of this proposition is that its adoption permits the problematic hierarchy, $M_W/M_{\text{Planck}} \sim 10^{-16}$, to be completely understood without using any dimensionless inputs which are smaller than 1%. We believe this represents considerable improvement over other explanations of the hierarchy problem, which typically explain the small value of $M_W/M_{\text{Planck}}$ in terms of another kind of hierarchy. For instance, the recently-proposed intriguing possibility that $M_s$ might be as low as the TeV scale [2, 3, 4, 5, 6, 7, 8, 9, 10], requires a compactification scale, $M_c$, which satisfies $M_c/M_s \leq 10^{-5}$.

Besides ameliorating some of the phenomenological problems of the TeV-scale scenario (such as too-fast proton decay), we find a number of other good features follow if $M_s$ is identified with $10^{11} \text{ GeV}$. These include the possibility of logarithmic perturbative unification (at $M_s$) of the Standard Model (SM) gauge couplings; naturally-occurring axions having astrophysically-acceptable couplings and Peccei-Quinn (PQ) scales; potentially interesting neutrino masses, etc.

Our arguments are based on the more general connection between $M_s$ and $M_{\text{Planck}}$ which occurs in perturbative Type-I string theory:

$$
\frac{M_c^{(p-6)}}{M_s^{(p-7)}} = \frac{\alpha_p M_{\text{Planck}}}{\sqrt{2}}
$$

where $p$ corresponds to the appropriate p-brane from which the gauge group (with coupling $\alpha_p$) originates. In addition, the condition that we remain within the realm of perturbation theory requires the corresponding $D = 10$ dilaton coupling, $\lambda_s$, to obey:

$$
\lambda_s = 2\alpha_p \left( \frac{M_s}{M_c} \right)^{p-3} = 2\sqrt{2} \frac{M_s^4}{M_c^3 M_{\text{Planck}}} \leq O(1).
$$

These two conditions require (in this simple isotropic case) the compactification scale, $M_c$, to not be much smaller than the string scale. Apart from this there is a remarkable freedom in choosing these scales.
There are, however, several natural options for these scales which suggest themselves. In order to explore these, consider for definiteness an embedding of the SM interactions within a set of coincident 3-branes.

1. The first possibility is to conform with perturbative heterotic tradition, and to place the string scale not far from the Planck scale.

2. A slightly better idea is to identify \( M_s \) with the GUT scale, \( M_X = M_s = 2 \times 10^{16} \) GeV, as indicated by the extrapolation of the low-energy couplings in the MSSM. This is consistent with the value of \( M_{\text{Planck}} \) inferred from (1.1), if we appropriately choose \( M_c \) very slightly below \( M_s \). Thus, in this way the old problem of reconciling gauge-coupling unification with perturbative heterotic strings is naturally solved. On the other hand, this scenario offers no explanation for the origin of the huge hierarchy of scales between the Fermi scale, \( M_W \), and \( M_{\text{Planck}} \), which must instead be blamed on some non-perturbative mechanism, like gaugino condensation.

3. Another alternative which has received recently much attention is the possibility of bringing \( M_s \) down to the weak scale: \( M_s \propto 1 \) TeV. This is a very intriguing possibility, raising as it does the possibility of testing string theory at accelerators. In this case, the choice \( M_c/M_s \leq 10^{-5} \) is required in order to obtain the correct value of \( M_{\text{Planck}} \). One trades in this way the standard hierarchy, \( M_W/M_{\text{Planck}} \), for this less extreme, but small, ratio. On the other hand, reconciling this scheme with constraints from cosmology, proton stability and gauge-coupling unification may prove non-trivial.

As stated above, we here propose to identify the string scale with the intermediate scale, \( M_s = \sqrt{M_W M_{\text{Planck}}} \propto 10^{11} \) GeV. This scale arises in a number of phenomenological settings, as we discuss below. The most interesting point of this scenario is the economy of its explanation of the \( M_W/M_{\text{Planck}} \) hierarchy, which here arises from the amplification of an initially very modest suppression, \( M_c/M_s \propto 10^{-2} \). This amplification occurs, without the need for any special hierarchy-generating mechanism such as gaugino condensation, due to the large powers of \( M_c/M_s \) which appears in eq. (1.1).

(After completing this paper we discovered some earlier work by Benakli [9] (see also [1]) which explores some of the advantages (including the connection with invisible
axions and neutrino masses) of a having the string scale of order $10^{11} - 10^{14}$ GeV. Different scenarios in possibilities for the generation of SUSY-breaking in the brane context have been also considered in refs. [8, 3, 11]. Similar studies in the context of the Horava-Witten M-theory scheme can be found in refs. [12].

To see how the $M_W/M_{\text{Planck}}$ hierarchy arises, we must determine how $M_W$ depends on the basic scales, $M_c$ and $M_s$. This dependence arises once supersymmetry breaks, which take to happen in a hidden sector of the model. Hidden sectors arise naturally within Type I string vacua. Let us consider for definiteness two separate sets of parallel 3-branes. We imagine ourselves to live on one set, containing the SM, while the other contains the hidden-sector interactions. If the positions of these 3-branes in the transverse space are sufficiently different, there are no massless states charged under both 3-brane groups. In what follows we consider an ideal situation of this type, in which we have the set of 3-branes containing the SM particles have unbroken $N = 1$ SUSY, and the distant (hidden) set of parallel 3-branes somehow completely break SUSY. We need not make any particular assumption about the nature of the SUSY breaking in that sector.

SM particles do not directly feel the breaking of SUSY which takes place in the hidden 3-brane sector. The effects for the SM of SUSY-breaking are only transmitted through the influence of closed-string sector fields, which are the only ones which can move into the bulk of spacetime and couple to both kinds of 3-brane sectors. Thus the SUSY-breaking felt by the SM fields are automatically suppressed by powers of $M_{\text{Planck}}$. In particular, if no particular suppression of SUSY-breaking in the hidden sector is assumed, one expects that the SUSY-breaking soft terms felt in the SM sector are of order:

$$M_W \sim m_{3/2} \sim \frac{F}{M_{\text{Planck}}} \sim \frac{M_s^2}{M_{\text{Planck}}}$$

Consider next, for simplicity, that all six of the compact dimensions share a common overall compactification scale, $M_c$. From the above formulae one has:

$$\frac{M_W}{M_{\text{Planck}}} = \frac{\alpha_3^2}{2} \left( \frac{M_c}{M_s} \right)^6$$

Or, more specifically:

$$M_W = \frac{\alpha_3 M_c^3}{\sqrt{2} M_s^2}; \quad M_{\text{Planck}} = \frac{\sqrt{2} M_s^4}{\alpha_3 M_c^2}.$$
Notice that SUSY-breaking disappears as $M_c \to 0$ and the distance between visible and hidden 3-branes go to infinity. Furthermore, if one takes $M_s/M_c \propto 160$ and $\alpha_3 = 1/24$ one indeed obtains the desired hierarchy $M_W/M_{\text{Planck}} = 10^{-16}$. As claimed, this small ratio arises from the large power of $M_c/M_s$ which amplifies a modest input value for $M_c/M_s$. It is remarkable how such a large hierarchy of 16 orders of magnitude can so naturally appear from such a modest initial suppression $M_c/M_s \propto 10^{-2}$. We regard this initial ‘hierarchy’ of $10^{-2}$ to be no hierarchy at all, since such small numbers are easy to obtain. Furthermore, we know that small numbers of this order have to appear elsewhere in the theory anyhow, such as if we are to understand the intergeneration ratio of Yukawa couplings.

Several remarks are worth recording before passing on to the phenomenological implications of this scenario.

1. A similar argument works equally well for the $p = 9$ case, in which similar formulae are obtained, but with the replacements $\alpha_3 \leftrightarrow \alpha_9$ and $M_s/M_c \leftrightarrow M_c/M_s$. In this case the required input factor would be the inverse of that just discussed, $M_c/M_s \propto 160$. The condition that we remain within Type I perturbation theory further requires $\lambda_s = 2\alpha_p(M_s/M_c)^{(p-3)} < 1$. This is satisfied for both 3-branes and 9-branes. In the 9-brane case, however the Type I dilaton coupling $\lambda_s$ would have to be very small and the scheme becomes less natural.

2. Next, we remark in passing on another fascinating connection between the two derived scales, $M_W$ and $M_{\text{Planck}}$, which are generated from the two fundamental scales of the model, $M_s$ and $M_c$. If we write $\alpha' \sim 1/M_s^2$ then $M_W$ and $M_{\text{Planck}}$ are related one to the other by $M_W = 1/(\alpha'M_{\text{Planck}})$. This is reminiscent of a ‘T-duality’ relationship, raising the tantalizing speculation that perhaps in some sense the physics at the Fermi scale might turn out to be ‘T-dual’ to the physics at the Planck scale.

3. Finally, we ask whether a similar explanation of the hierarchy problem is possible within the Horava-Witten representation of strongly coupling heterotic strings. We find the situation to be different in this case, for which the following relations obtain [1]:

$$M_{\text{Planck}}^2 = \frac{M^9_m}{M_c^6 M_\rho} \quad \alpha_9 = (\sqrt{2\pi})^{2/3} \left(\frac{M_c}{M_m}\right)^6, \quad (1.6)$$
where $M_m$ is the 11-dimensional $M$-theory mass scale, $M_\rho$ is the mass scale associated to the 1-dimensional interval and the other parameters are like in type-I theory. If we have that $M_W = M_m^2 / M_{Planck}$ then we obtain the following relations:

$$\frac{M_W}{M_{Planck}} = \left(\frac{M_c}{M_m}\right)^6 \left(\frac{M_\rho}{M_m}\right)^{\alpha_9} = (\sqrt{2}\pi)^{2/3} \left(\frac{M_c}{M_m}\right)^6$$

(1.7)

Therefore, from the first relation it seems that we can also obtain the hierarchy by a small hierarchy between the compactification scales and the string scale. However, the second relation tells us that if the compactification scale $M_c$ is very small compared with the $M$-theory scale, then the gauge coupling would be far too small to agree with experiment. We can still use a standard value for $\alpha_9$ and obtain the hierarchy between the electroweak and Planck scales, but only in terms of a similarly large hierarchy between the $M$-theory scale and the length of the 11-dimensional interval. The ultimate origin of this hierarchy then remains unexplained. On the other hand, contrary to the case with weak-scale $M$-theory, it is in principle possible to have an intermediate scale $M$-theory without having an unacceptably large value for the interval length, $\rho$, since the observed hierarchy implies in this case $\rho \sim 10^{-12}$ meters.

2 Phenomenological Issues

A number of questions and phenomenological implications appear in this scheme:

i) Gauge coupling unification

If the string scale is of order $10^{11}$ GeV, one would also expect that gauge couplings of the SM should also join at this scale. We now argue that this kind of unification does not need fast power-like running, as would be mandatory for a weak-scale-string scenario. Indeed, if there are further particles charged only under $SU(2) \times U(1)$ with masses of order $M_W$ in the massless spectrum beyond those of the MSSM, the $SU(2) \times U(1)$ couplings would grow faster and so intersect the $SU(3)$ coupling precociously.

To see this, recall the one-loop formulae for the running of the SM gauge couplings:

$$\sin^2 \theta_W(M_Z) = \frac{3}{8} \left(1 + \frac{5\alpha(M_Z)}{6\pi} \left(b_2 - \frac{3}{8} b_1 \right) \log\left(\frac{M_Z}{M_{Z_0}}\right)\right)$$

(2.1)
\[
\frac{1}{\alpha_3(M_Z)} = \frac{3}{8} \frac{1}{\alpha(M_Z)} - \frac{1}{2\pi} \left( b_1 + b_2 - \frac{8}{3} b_3 \right) \log \left( \frac{M_s}{M_Z} \right)
\]

(2.2)

where \(b_i\) are the SM beta-function coefficients. In the MSSM one has \(b_1 = 11, b_2 = 1, b_3 = -3\) yielding unification at \(2 \times 10^{16}\) GeV. In the present case one can check that increasing \(b_2\) by three units makes \(\alpha_2\) cross \(\alpha_3\) at around \(M_s = 10^{11}\) GeV as required. In order to obtain consistent values with both \(\alpha_3\) and \(\sin^2 \theta_W\) one has further to increase the value \(b_1\) by around eleven units. An example of additional particles which can produce these beta functions is given by supplementing the three SM quark-lepton generations with a collection of new chiral fields transforming like e.g.

\[
4[(1, 1, 1) + (1, 1, -1) + (1, 2, 1/2) + (1, 2, -1/2)]
\]

under \(SU(3) \times SU(2) \times U(1)\). These are just four standard sets of of vector-like leptons. A pair of doublets in this case should be identified with the SM Higgs. In this case one finds \(\alpha_3(M_Z) = 0.12\) and \(\sin^2 \theta_W(M_Z) = 0.23\) for \(M_s = 3 \times 10^8 M_W\). Of course this particularly simple choice is not unique, and other combinations could be possible. As discussed below, such an extension of the MSSM might be interesting if one wants to gauge a symmetry like lepton number to insure proton stability.

Another alternative is to consider extensions to the SM gauge group like \(SU(4) \times SU(2)_L \times SU(2)_R\). (Such a group often appears in orientifold constructions.) This kind of gauge group can also lead to precocious unification. Clearly, a more systematic study of the unification possibilities at the geometrical scale \(M_s \propto 10^{11}\) GeV would certainly be interesting.

We conclude that gauge coupling unification can be easily accommodated in this scheme, although it is not, properly speaking, predicted.

ii) Proton stability

It is well known that generic versions of the SUSY SM contain \(D = 4\) operators which violate baryon- and/or lepton number unless some symmetry (like e.g., the standard discrete R-parity) is imposed on the model. In addition, even if those \(D = 4\) terms are forbidden, if the string and unification scales are of order of the intermediate scale in general there will appear dimension 5 and dimension 6 (and even higher) operators which will violate baryon and lepton numbers.
Thus there must exist a symmetry forbidding these unwanted effects also. The problem of dimension five operators (although in a less severe form) is in fact already present in the scheme with $M_s = 10^{16}$ GeV and is much more problematic in the schemes with $M_s = 1$ TeV. Thus this is really a problem common to all schemes with $M_s < M_{Planck}$. A natural solution is the presence of a continuous gauged $U(1)$ symmetry. In this context a symmetry like $U(1)_{B-L}$ which appears in left-right symmetric models may forbid R-parity violating operators of dimension 4 but again is not going to forbid some dangerous dimension 5 operators.

Another natural alternative is to consider a pseudoanomalous $U(1)_X$ whose triangle anomalies are cancelled via the 4-dimensional Green-Schwarz mechanism. The simplest flavour independent such a $U(1)_X$ one can think of is one which assigns charge=1 for all SM quarks and leptons, charge=-2 for the Higgs doublets $H, \bar{H}$. This simple symmetry naturally forbids baryon and lepton violating dimension-4 and 5 operators. Alas, in the case of perturbative heterotic vacua this mechanism is very restrictive. Indeed for the mechanism to be possible the mixed anomalies $A_i$ of this $U(1)_X$ with the SM gauge groups $G_i$, i=1,2,3 have to satisfy $A_1 : A_2 : A_3 = 5/3 : 1 : 1$ \[14\]. In the case of the MSSM one finds instead:

$$A_1 : A_2 : A_3 = 6 : 4 : 6 \quad (2.4)$$

If e.g., one adds the contribution of the extra states needed for gauge coupling unification which we suggested in the previous paragraph on has more freedom.

But in fact everything is simpler in the class of Type I $D = 4$ strings that we are considering. It has been recently realized \[15\] that in the context of Type IIB $D = 4$ orientifolds there is a generalized Green-Schwarz mechanism at work in which the mixed anomalies are not constrained to be in definite ratios. This is because of the generic presence of several twisted RR axionic fields which couple differently to different group factors. Thus an anomalous $U(1)_X$ like the one proposed above can be perfectly consistent within the context of Type I theories without the need of choosing particular charges for the extra particles present to get gauge coupling unification.

A $U(1)_X$ symmetry with quark and lepton charges as considered here was first considered by Weinberg \[26\] in the early days of SUSY model-building in order to avoid proton decay via dimension-4 and -5 operators. However he introduced extra exotic particles to get cancellation of $U(1)$ anomalies. This is not required in the
context of string theory. In any event the presence of such an anomalous $U(1)$ is enough to forbid $D = 4, 5$ operators. This is true even if the $U(1)$ becomes massive by swallowing RR axionic field, because then the symmetry would survive perturbatively as a global $U(1)$ symmetry.

As explained above, a simple anomalous $U(1)$ like the one discussed above only supresses dimension 4 and 5 baryon/lepton number violating operators, but not dimension 6 which need also to be supressed if the string scale is as low as $10^{11}$ GeV. However, the relevant operators will always involve quarks and lepton of the first two generations and hence one expects further supression factors. A classical example for this are the dimension 6 operators coming from the exchange of colour-triplet Higgs fields in the minimal SUSY-$SU(5)$ operator. From the proton-stability point of view the Higgs triplet can be as light as $10^{11}$ GeV because the triplet couples with very supressed Yukawa couplings to the first two generations.

Let us finally remark that anomalous $U(1)$s with similar general characteristics to the one we are proposing here do appear in specific $D = 4, N = 4$ Type IIB orientifolds (see ref. [6, 15]). In particular, $SU(n)$ groups come along with $U(1)$ factors with the same couplings in $U(n)$ gauge groups. Thus, e.g., if we have a set of 3-branes with the SM non-Abelian gauge group $SU(3) \times SU(2)$ one indeed expects to have $U(3) \times U(2)$ factors. One linear combination of the two $U(1)$s would be the hypercharge whereas the orthogonal one would typically be an anomalous $U(1)$.

Another alternative is the presence of anomaly-free $U(1)$s gauging either lepton or baryon numbers. As explained in [17] this is possible if one adds vector-like sets of leptons to the MSSM. We have seen that extra vector-like sets of leptons are wellcome in order to get gauge coupling unification, thus this is an alternative which could be at work.

iii) Soft terms and radiative electroweak symmetry breaking

In a simple scheme in which we assume that the SM and the hidden sector live in two separated sets of p-branes, only closed-string fields are able to transmit SUSY-breaking from the hidden to the visible sector. Thus it is natural to assume that the mediators of SUSY-breaking will be the dilaton and moduli fields. Specifically, in the context of Type IIB $D = 4, N = 1$ orientifolds the complex dilaton $S$ and untwisted moduli fields $T_i$ are able to couple to both sectors of p-branes. This is not in general
the case for twisted closed-string moduli, which couple only to sets of p-branes with positions in transverse space on top of the given orbifold singularities. In any event, it would make sense from the visible sector point of view to parametrize SUSY breaking in terms of the vacuum expectation values of the auxiliary fields of dilaton and untwisted moduli, $F_S, F_{T_i}$. This is the spirit previously applied to heterotic compactifications in refs. [18, 19, 20] and generalized to Type I type of vacua in ref. [21]. Thus possibilities like dilaton/modulus dominated boundary conditions may be particularly relevant in the present scheme. However one does not need to assume that $F_S$ and $F_{T_i}$ are the only fields contributing to the vacuum energy since in the hidden sector there may be other fields (e.g. the twisted moduli in Type IIB orientifolds) who contribute to SUSY-breaking but which do not couple directly to the visible p-branes.

The existence of fields which can contribute to SUSY-breaking but which do not couple to the visible p-brane sector somewhat changes the results for soft terms from dilaton/moduli dominance as computed in heterotic models. As explained in [22, 21], the structure and couplings of the massless p-brane sector in Abelian Type IIB $D = 4$ orientifolds is quite analogous to the untwisted sector of Abelian heterotic orbifolds and so is the Kahler potential (modulo some redefinitions of the $S$ and $T_i$ chiral fields) and renormalizable superpotential. Thus one obtains similar soft terms results as those found for the untwisted sector of heterotic orbifolds. The case of the existence of some field $\phi$ contributing to SUSY-breaking but not coupling to the visible world was in fact considered in chapter 8 of [19]. One finds for the case of an overall modulus field $T$ the result for the soft scalar and gaugino masses (assuming a vanishing cosmological constant):

$$m_0^2 = m_3^2/2(1 - \cos^2\theta \cos^2\phi) \ ; \ M_{1/2}^2 = 3m_3^2/2 \sin^2\theta$$

where $tg\theta = |F_S|//(\sqrt{3}|F_T|)$ and $sin\theta_\phi = |F_M|//(\sqrt{3}m_3^2)$. Thus in the absence of a SUSY-breaking contribution from $\phi$ one has $\cos\theta_\phi = 1$ and one goes back to the usual dilaton/modulus dominated limit with $M_{1/2}^2 = 3m_0^2$. On the other hand if there are fields $\phi$ contributing to SUSY-breaking but not coupling to the visible p-branes (like e.g. the twisted fields we mentioned above) one sees that smaller gaugino masses with $M_{1/2} \leq \sqrt{3}m_0$ are now possible. Notice that in the above expressions one has $m_3^2 \propto \alpha_3/2(M_c/M_s)^6$.

Unlike the situation in the weak-scale string scenarios, in the present scheme, once
SUSY breaking of order $\alpha_3/\sqrt{2}(M_3^2/M_2^2)$ is transmitted to the visible 3-brane sector, radiative $SU(2) \times U(1)$ breaking occurs in the usual way, since there is plenty of room for the evolution of the soft masses from $10^{11}$ GeV to the weak scale. Since, nevertheless, the space for running is substantially smaller than in the usual MSSM scheme, for a given set of soft terms, radiative breaking in the present scheme will in general require higher top-quark Yukawa couplings. This is because larger top-quark Yukawa couplings make the Higgs field squared-mass run faster towards smaller values.

iv) Invisible axions and the Strong CP problem

One of the bonuses of the present scheme is that it seems to provide a general scenario for the solution of the strong CP problem. It is well known that astrophysical constraints coming from the stability of red-giants impose the lower bound on the decay constant of an axion $M_{PQ} > 10^{10}$ GeV. In fact there are also cosmological bounds which give an upper bound of the same order of magnitude, fixing the Peccei-Quinn scale around the intermediate scale. In any event, it is well known that in Type I string theory there may be a plethora of RR scalars with axion-like couplings (like the twisted RR fields mentioned in the previous subsection). The typical scale of the corresponding decay constants is expected to be of order $M_s = 10^{11}$ GeV, consistent with the astrophysical bounds. Thus the RR scalars of Type I theory provide naturally with the required invisible axions at the appropriate mass scale.

String axions appearing in perturbative heterotic string theory have been studied in the past as possible candidates for invisible axions. In particular the model-independent ImS axion as well as the imaginary partners of $T$ moduli which can also get axionic couplings at one loop. However in the perturbative heterotic case the natural Peccei-Quinn scale is the string scale or slightly below and thus it is too large. Also it is difficult to avoid that these would be axions would get a mass upon SUSY-breaking from space-time or world-sheet instantons.

In the present scheme the string scale coincides with the Peccei-Quinn scale, so this part of the problem is solved. On the other hand in e.g., Type IIB orientifolds not only the untwisted moduli $S, T_i$ but also twisted Ramond-Ramond (RR) fields have axionic couplings. Although $ImS$ and or $ImT_i$ might get too-large masses from interactions with the $N = 0, 3$-brane sector, that is not going to be in general the case for some twisted RR-fields $ImM_a$. Indeed, these fields couple only to 3-branes situated
on the corresponding orbifold singularity (whose blowing-up mode is related to \(ReM_a\)). Thus \(ImM_a\) axionic fields coupling to SM gauge groups will not couple to hidden sector groups. This will guarantee that only QCD instantons will determine the potential of these axions. This will then lead to an automatic solution of the strong CP problem.

v) Neutrino masses

Intermediate scales of order \(10^{10-13}\) GeV are popular in the literature in the context of see-saw models of neutrino masses. Integrating out any heavy Majorana neutrinos, whose masses we suppose are of order \(M_s\), generates the effective interaction \(LLHH\) in the low-energy superpotential, where \(L\) and \(H\) represent SM lepton and higgs doublets. The coefficient of such an operator is of order \(a/M_s\), where \(a\) contains products of neutrino Yukawa couplings. The Majorana masses which follow for the light neutrino species are then \(m_\nu \sim a\langle H\rangle^2/M_s \sim a(10\text{ eV})\).

For reasonable values for \(a\), masses this size can lie below the experimental limit (from \(\beta\beta\) decay) of \(\sim 1\) eV for the electron neutrino, and can easily lie in a range which is consistent with atmospheric, solar and/or LSND results. (Accomodating all three typically would require light sterile neutrinos, which are also possible in the scenario we are considering.)

Thus in the present scheme the right-handed neutrinos could be just massive string states with masses of order \(M_s\), making current neutrino experiments windows onto string physics!

vi) Wimpzillas and ultra-high energy cosmic rays

Intermediate scales have also recently appeared in the context of supermassive, stable, particles with masses of order \(10^{12} - 10^{16}\) GeV (the so-called ‘Wimpzillas’ [25]) which could constitute an interesting candidate for cold dark matter. Particles associated to the string scale could provide candidates for such states. Furthermore, decaying particles associated to the same string scale could be candidate sources to generate ultra-high energy cosmic rays [16] (see [27] and references therein) with energies \(\propto 10^9 - 10^{10}\) GeV found experimentally.

vii) Cosmology

The cosmological implications of the intermediate scale may be many. Although usually inflation is considered much closer to the Planck scale, there is no impediment for it to occur at an intermediate scale. The amount of fine tuning needed to obtain the
expected number of $e$-foldings and the spectrum of perturbations is of course model-dependent. For a single inflaton field this will imply an inflaton mass of the order of $M_s^2/M_{Planck}$, i.e. the electroweak mass, which is quite reasonable, compared with a $10^{-13}$ GeV inflaton in the TeV string scenario [10]. Furthermore with an intermediate scale string theory it is straightforward to obtain large enough reheating temperatures in order to produce the standard model fields after inflation [10]. Being sufficiently higher than the electroweak scale, there should not be any problem to generate the baryon asymmetry. A detailed study of all these issues may be interesting.

viii) Generating $M_c$ and the breathing mode

Considerable effort [4] has been devoted to understanding the hierarchy $M_c/M_s$ in models with a weak-scale string scale. This proves to be reasonably difficult to do, due to the partially-conflicting constraints that: (i) $M_c$ be small enough to give the correct value for $M_{Planck}$, and yet that the breathing mode, $T$, of the compact dimensions not be so light as to be ruled out by constraints on long-range, gravitational-strength scalar interactions.

Clearly a hierarchy like $M_c/M_s \sim 10^{-2}$ is much easier to generate, to the extent that it does not require much explanation at all. Furthermore the constraints on the mass of the breathing mode in our scenario are also much weaker. This is because the mass of any such mode can be at most as large as $M_c$. Since for energies larger than $M_c$ the breathing mode is a component of the extra-dimensional metric, and so is required to be massless by general covariance. Generically the breathing mode can be much lighter than this upper bound. Direct experimental limits on ‘fifth forces’ preclude this mass being smaller than an inverse millimetre, which is easy to satisfy when $M_c \sim 10^9$ GeV, but more difficult when $M_c \ll M_s \sim 1$ TeV.

3 Conclusions

In summary, we have proposed as a natural scheme one in which the string scale $M_s$ is identified with the intermediate scale $M_s = \sqrt{M_W M_{Planck}}$. An important property of this scenario is that the hierarchy $M_W/M_{Planck}$ appears as a consequence of the compactification scale being just a couple of orders of magnitude below the string scale, due to the large power appearing in $M_W/M_{Planck} = \frac{1}{2} \alpha_3^2 (M_c/M_s)^6$. This natural
generation of the huge hierarchy is a nice feature compared to the standard assumption that both string and Planck scales are close to one another. Essentially what we assume is the existence of a perturbative Type I string vacuum with an $N = 1$ supersymmetric 3-brane world (including the SM) and a separate 3-brane world with no supersymmetry.

We remark that the weak and Planck scales in this scheme are related by a ‘dual-looking’ formula: $M_W = 1/\alpha' M_{\text{Planck}}$. It would be interesting to see whether this is a reflection of some underlying duality between the physics of these two scales.

Our scheme can easily accomodate logarithmic gauge-coupling unification at the price of supplementing the particle content of the MSSM just above the weak scale. This is not as natural as in a conventional grand-unified scenario, but is certainly simpler than gauge-coupling unification in the 1 TeV scenarios, which requires power-like running of coupling constants.

Another attractive aspect of the present scheme is that the string scale turns out to coincide with the Peccei-Quinn scale that is required by astrophysical constraints for invisible axion models which solve the strong CP problem. Since string theory has abundant axion-like fields (like e.g., the twisted RR fields in Type I orientifold models), and some of these may couple to the SM but not to the hidden sector, this problem could be naturally solved.

String theory at an intermediate scale could also be relevant for other phenomenological issues like neutrino masses and cosmology. In any event, it seems that if indeed the string scale is of order the intermediate scale $\sim 10^{11}$ GeV, string theory could be much more amenable to experimental test than previously thought.

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References

[1] E. Witten, Nucl. Phys. B471 (1996) 135, [hep-th/9602070].

[2] J.D. Lykken, [hep-th/9603133].

[3] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, [hep-ph/9803315];
    I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, [hep-ph/9804398];
    I. Antoniadis, S. Dimopoulos, A. Pomarol and M. Quiros, [hep-ph/9810419].

[4] N. Arkani-Hamed, S. Dimopoulos and J. March-Russell, [hep-th/9809124].

[5] K. Dienes, E. Dudas and T. Gherghetta, [hep-ph/9803466]; [hep-ph/9806292]; [hep-ph/9807522];
    K. Dienes, E. Dudas, T. Gherghetta and A. Riotto, [hep-ph/9809406].

[6] G. Shiu and S.H. Tye, [hep-th/9805157].

[7] C. Bachas, [hep-ph/9807415].

[8] Z. Kakushadze and S.H. Tye, [hep-th/9809147].

[9] K. Benakli, [hep-ph/9809582].

[10] K. Benakli and S. Davidson, [hep-ph/9810280];
    D.H. Lyth, [hep-ph/9810320].

[11] L. Randall and R. Sundrum, [hep-th/9810155].

[12] T. Banks and M. Dine, Nucl. Phys. B479 (1996) 173;
    E. Dudas and C. Grojean, [hep-th/9704177];
    I Antoniadis and M. Quirós, [hep-th/9707208];
    Z. Lalak and S. Thomas, [hep-th/9707223];
    E. Dudas, [hep-th/9709043];
    H.P. Nilles, M. Olechowski and M. Yamaguchi, [hep-th/9707143];
    T.Li, J.Lopez and D. Nanopoulos, [hep-ph/9702237];
    K. Choi, [hep-th/9706171];
    K. Choi, H.B. Kim and C. Muñoz, [hep-th/9711158].
A. Lukas, B. Ovrut and D. Waldram, hep-th/9711197 ; hep-th/9710208 . E. Mirabelli and M. Peskin, hep-th/97065002 .

[13] E. Cáceres, V. Kaplunovsky and M. Mandelberg, Nucl. Phys. B493 (1997) 73; J. Ellis, A. Faraggi and D.V. Nanopoulos, Phys. Lett. B419 (1998) 123, hep-th/9709049.

[14] L.E. Ibáñez, Phys. Lett. B303 (1993) 55

[15] L. E. Ibáñez, R. Rabadán and A. Uranga, hep-th/9808139.

[16] S. Weinberg, Phys. Rev. D26 (1982) 287.

[17] L.E. Ibáñez and G.G. Ross, Nucl. Phys. B368 (1992) 3.

[18] M. Cvetic, A. Font, L.E. Ibáñez, D. Lüst and F. Quevedo, Nucl. Phys. B36 (1999) 194;
L.E. Ibáñez and D. Lüst, Nucl. Phys. B382 (1992) 305;
V. Kaplunovsky and J. Louis, Phys. Lett. B306 (1993) 269.

[19] A. Brignole, L.E. Ibáñez and C. Muñoz, Nucl. Phys. B422 (1994) 125.

[20] A. Brignole, L. E. Ibáñez, C. Muñoz and C. Scheich, Z.f.Phys.C 74, (1997) 157, hep-ph/9508258.

[21] L.E. Ibáñez, C. Muñoz and S. Rigolin, to appear.

[22] G. Aldazabal, A. Font, L.E. Ibáñez and G. Viñoler, hep-th/9804026.

[23] D. Dicus et al., Phys. Rev. D22 (1980) 839;
M. Fukugita, S. Watamura and M. Yoshimura, Phys. Rev. D26 (1982) 1840.

[24] E. Witten, Phys. Lett. B149 (1984) 351;
K. Choi and J.E. Kim, Phys. Lett. B165 (1985) 71;
T. Banks and M. Dine, hep-th/9605136 ; hep-th/9609040;
K. Choi, hep-th/9706171.

[25] See e.g. E. Kolb, D. Chung and A. Riotto, hep-ph/9810361.
[26] S. Weinberg, private communication (1994).

[27] V.A. Kuzmin and V.A. Rubakov, Phys.Atom.Nucl. 61(1998)1028;
    M. Birkel and S. Sarkar, hep-ph/9804285.