PHENOMENOLOGICAL SURVEY OF M–THEORY

ALON E. FARAGGI

Theoretical Physics Department, University of Oxford, Oxford OX1 3NP, UK
E-mail: faraggi@thphys.ox.ac.uk

The Standard Model data suggests that the quantum gravity vacuum should accommodate two pivotal ingredients. The existence of three chiral generations and their embedding in chiral 16 SO(10) representations. The $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifolds are examples of perturbative heterotic string vacua that yield these properties. The exploration of these models in the nonperturbative framework of M–theory is discussed. A common prediction of these constructions is the existence of super–heavy meta–stable states due to the Wilson–line breaking of the GUT symmetries. Cosmic ray experiments in the forthcoming years offer an exciting experimental window to the phenomenology of such states.

1 Introduction

Over the past few years substantial progress has been achieved in the basic understanding of string theory. The picture which emerged, and which is depicted qualitatively in figure 1, is that the different string theories in ten dimensions are limits of a more fundamental theory, traditionally dubbed M–theory. The question remains how to connect these advances, and string theory in general, to experimental data. I think that it will be universally agreed that this is a vital question that string theory faces, and opinions may differ on what is the most suitable methodology to advance this issue.

In this respect I think that there are some prevailing misconceptions. The first question that should be posed is why one should be interested in string theory in the first place. Physics is first and foremost an experimental science, and after all the recent celebrated advances in string theory have to do with unification of theories in ten or eleven dimensions, and what has this to do with experimental physics? Nevertheless, given the present experimental data, the exploration of string theory is well motivated.

One misconception is the reference to string theory as the “theory of everything”. I think that it is besides the point. The primary questions in my view are those of relevance and utility. Namely, the experimental data that we observe reveal patterns of gauge charges and mass spectra. The main issue is whether the structures that appear in string theory are relevant and can be utilized toward understanding the physical origin of the observed patterns. In this respect another misconception is the view that the sole merit of string phenomenology is to find the one true vacuum that corresponds to the observed world. While this is a well posed goal that is pursued with vigor

---

*INVITED TALK PRESENTED AT SUGRA 20, BOSTON MA, MARCH 17–21 2003.
†OUTP-03-17P
and intent, I think that it is again besides the point. In the first place, as is especially evident following the string duality developments, none of the string limits can fully characterize the true vacuum. The true vacuum should have some nonperturbative realization. The perturbative string theory limits and the eleven dimensional classical limit are effective limits that can at best probe some of the properties of the true vacuum. In this view it may well be that some limits may highlight some properties of the vacuum, whereas other limits may be more instrumental to extract different properties. A good example of this is the dilaton stabilization problem. As is well known, in the perturbative heterotic–string limit the dilaton, whose VEV governs the string gauge and gravitational couplings, has a run–away potential and cannot be stabilized at a finite value. However, we should regard the heterotic limit as the zero coupling expansion of the more basic theory. With our present understanding of string theories in the context of their M–theory embedding it is clear that we should not in fact expect the dilaton to be stabilized in the heterotic limit. In order to stabilize the dilaton we have to move away from the zero coupling expansion, or to move away from the perturbative heterotic–string limit. The existence of the classical eleven dimensional limit in which the dilaton is interpreted as the moduli of the eleventh dimension lends credence to this general expectation. Thus, the issue of dilaton stabilization may be more accessible, even if not yet fully resolved, in other limits of the underlying theory, rather than in the perturbative heterotic string limit. The problem of supersymmetry breaking may be similar.

The primary questions in respect to string and M–theories are therefore those of relevance and utility. The relevance follows from the basic structure
of the Standard Model of particle physics. The Standard PArticle Model (SPAM) matter sector is composed of three chiral generations, charged under the three group factors $SU(3) \times SU(2) \times U(1)_Y$. The most remarkable aspect of the SPAM, and for me its essence, is the GUT embedding of the Standard Model representations. Most striking is the embedding in $SO(10)$ in which each generation fits into a single 16 spinorial representation of $SO(10)$. The GUT embedding of the Standard Model spectrum is depicted in figure 2. If we regard the gauge charges of the Standard Model states as experimental observables, as they were in the process of its experimental discovery, then 54 parameters are needed at the level of the Standard Model to account for the matter charges. The embedding in $SO(10)$ reduces this number of parameters to three, which is the number of chiral 16 $SO(10)$ representations needed to accommodate the Standard Model states.

The GUT embedding of the Standard Model is also supported by the logarithmic running of the Standard Model parameters. Quantum field theories, in general, and their specific realization in the form of the Standard Model, predict that the gauge and matter couplings evolve logarithmically with the energy scale. This evolution has been confirmed experimentally in the energy range accessible to collider experiments. It is also consistent qualitatively with the big desert scenario, suggested by grand unified theories, in the gauge and matter sectors of the Standard Model. The scalar sector of the Standard Model is not protected against radiative corrections from higher scales, and therefore the Standard Model needs to be augmented with an additional sector that protects the scalar states from higher scale corrections. It is natural to demand that the new sector preserves the successes of the Standard Model and its GUT embedding. Such a new sector is provided by
supersymmetry. While the jury is still out on the validity of supersymmetry and its contemporary interpretation, it seems to me as the best of all evils.

While the Standard Model, GUTs, supersymmetry and point quantum field theories have enjoyed impressive successes, they still fall short of providing a comprehensive framework for the fundamental forces and matter that are observed experimentally. In the first place gravity is not yet included as a quantum theory. It is gravely unsatisfactory to have two fundamental theories, gravity and quantum mechanics, that are incompatible. More concretely, from the point of view of the Standard Model data, GUTs and supersymmetry do not explain the origin of many of the parameters that we observe in the Standard Model. Specifically, the existence of flavor with its intricate mass and mixing spectrum is unaccounted for. It is therefore plausible that the origin of these additional structures must be sought in a theory that unifies gravity and the gauge interactions.

Superstring theories are unique in the sense that they provide exactly that. Namely, string theories give rise to precisely the structures that are observed in the Standard Model, like matter and gauge spectrum, and they provide consistent framework for perturbative quantum gravity. Hence the utility of string theory. In regard to its relevance, the jury is of course still out on that, but we note that string theory gives rise to additional sectors, that may be precisely what is needed to understand the detailed spectrum of the Standard Model. These new sectors include the requirement of compactified manifolds that may account for the existence of flavor. Thus, string theories gives rise to the structures that may eventually prove relevant for the understanding of the Standard Model data. Hence its relevance.

Getting back to the qualitative picture of figure 1, the question is: how should we utilize the new understanding of string theories to advance its phenomenological studies. As discussed above, none of the string limits should be regarded as fundamental and therefore each limit can at best reveal some properties of the true nonperturbative vacuum. We should also consider the possibility that in the end the true fundamental vacuum may be probed only by its perturbative limits, and that the underlying nonperturbative theory be defined only for conceptual consistency. The new understanding of string theory suggests the following approach. Suppose that we are able to identify in some limit a class of vacua that appear viable from a phenomenological perspective. Vacua here refers to specific classes of compactified manifolds. Further insight into the properties of the phenomenological vacua may therefore be gleaned by compactifying the other string limits on the same class of compactified manifolds.

As reasoned above, the essential property of the Standard Particle Model is its embedding in $SO(10)$ GUT. From this perspective, the primary guides in the search for phenomenological string vacua should be the existence of three chiral generations and their $SO(10)$ embedding. The class of string vacua that we seek are those for which there exist a limit that preserves the $SO(10)$
embedding of the Standard Model spectrum. The only perturbative string limit which enables the $SO(10)$ embedding of the Standard Model spectrum is the heterotic $E_8 \times E_8$ string. The reason being that this is the only limit that produces the spinorial 16 representation in the perturbative massless spectrum. In this respect it is likely that other M–theory limits provide more useful means to study other properties of the fundamental vacuum, such as dilaton and moduli stabilization.

2 Perturbative phenomenology

The study of phenomenological string vacua proceeds with the compactification of the heterotic string from ten to four dimensions. A class string compactifications that preserve the $SO(10)$ embedding of the Standard Model spectrum are those that are based on the $Z_2 \times Z_2$ orbifold and have been extensively studied by utilizing the so–called free fermionic formulation. The structure of these models have been amply reviewed in the past. The models are constructed by specifying a set of boundary condition basis vectors and the one-loop GSO projection coefficients. The first five basis vectors of the realistic free fermionic models consist of the NAHE set. The gauge group after the NAHE set is $SO(10) \times E_8 \times SO(6)^3$ with $N = 1$ space–time supersymmetry, and 48 spinorial 16 of $SO(10)$, sixteen from each sector $b_1$, $b_2$ and $b_3$, which are the three twisted sectors of the corresponding $Z_2 \times Z_2$ orbifold. The $Z_2 \times Z_2$ orbifold is special precisely because of the existence of three twisted sectors, that naturally yields three generation models, one from each of the twisted sectors. The construction proceeds by adding to the NAHE set three additional boundary condition basis vectors which break $SO(10)$ to one of its subgroups. At the same time the number of generations is reduced to three generations. One spinorial of $SO(10)$, decomposed under the final $SO(10)$ unbroken subgroup, is obtained from each of the twisted $b_1$, $b_2$ and $b_3$. Consequently the weak hypercharge, which arises as the usual combination $U(1) = 1/2U(1)_{B-L} + U(1)_{T_{3R}}$, has the standard $SO(10)$ embedding. The models contain several electroweak Higgs multiplets and couplings that may yield qualitatively viable fermion mass texture.

In addition to the standard GUT spectrum, the string models also contain exotic states which arise from the basis vectors that break the $SO(10)$ symmetry. These states carry either fractional $U(1)_Y$ or $U(1)_{Z'}$ charge. Such states are generic in superstring models and impose severe constraints on their validity. In some cases the exotic fractionally charged states cannot decouple from the massless spectrum, and their presence invalidates otherwise viable model. In the NAHE based models the fractionally charged states always appear in vector–like representations, and, in general, mass terms are generated from renormalizable or nonrenormalizable terms in the superpotential. The analysis of ref. demonstrated the existence of free fermionic models
with solely the MSSM spectrum in the low energy effective field theory of the Standard Model charged matter. In general, unlike the “standard” spectrum, the “exotic” spectrum is highly model dependent.

The free fermionic string models provide the arena for studying many of issues that pertain to the phenomenology of the Standard Model and Unification. Many of these issues have been the subject of past studies, that include among others: top quark mass prediction, several years prior to the actual observation by the CDF/D0 collaborations; generations mass hierarchy; CKM mixing; superstring see–saw mechanism; Gauge coupling unification; Proton stability; and supersymmetry breaking and squark degeneracy.

3 \( Z_2 \times Z_2 \) orbifold correspondence

The key property of the fermionic models that is exploited in trying to elevate the analysis of these models to the nonperturbative domain of M–theory is the correspondence with the \( Z_2 \times Z_2 \) orbifold compactification. The correspondence of the NAHE-based free fermionic models with the orbifold construction is illustrated by extending the NAHE set, \( \{1, S, b_1, b_2, b_3\} \), by one additional basis vector \( \xi_1 \). With a suitable choice of the GSO projection coefficients the model possesses an \( SO(4)^3 \times E_6 \times U(1)^2 \times E_8 \) gauge group and \( N = 1 \) space-time supersymmetry. The matter fields include 24 generations in the 27 representation of \( E_6 \), eight from each of the sectors \( b_1 \oplus b_1 + b_2 + b_3 + \xi_1 \) and \( b_2 \oplus b_2 + b_3 + \xi_1 \). Three additional 27 and 27 pairs are obtained from the Neveu-Schwarz sector.

To construct the model in the orbifold formulation one starts with the compactification on a torus with nontrivial background fields. The subset of basis vectors, \( \{1, S, \xi_1, \xi_2\} \), generates a toroidally-compactified model with \( N = 4 \) space-time supersymmetry and \( SO(12) \times E_8 \times E_8 \) gauge group. The same model is obtained in the geometric (bosonic) language by tuning the background fields to the values corresponding to the \( SO(12) \) lattice. The metric of the six-dimensional compactified manifold is then the Cartan matrix of \( SO(12) \), while the antisymmetric tensor is given by \( b_{ij} = g_{ij} \) for \( i > j \). When all the radii of the six-dimensional compactified manifold are fixed at \( R_I = \sqrt{2} \), it is seen that the left- and right-moving momenta reproduce the massless root vectors in the lattice of \( SO(12) \). Adding the two basis vectors \( b_1 \) and \( b_2 \) corresponds to the \( Z_2 \times Z_2 \) orbifold model with standard embedding. Starting from the Narain model with \( SO(12) \times E_8 \times E_8 \) symmetry, and applying the \( Z_2 \times Z_2 \) twist on the internal coordinates, reproduces the spectrum of the free-fermion model with the basis \( \{1, S, \xi_1, \xi_2, b_1, b_2\} \). The Euler characteristic of this model is 48 with \( (h_{11}, h_{21}) = (27, 3) \), and it is denoted as \( X_2 \). The four dimensional gauge symmetry at this stage can be either \( E_6 \times U(1)^2 \times SO(4)^3 \times E_8 \), or \( SO(10) \times U(1)^3 \times SO(4)^3 \times SO(16) \).
depending on the choice of GSO phase \( c(\xi_1, \xi_2) = \pm 1 \).

The \( Z_2 \times Z_2 \) orbifold of the \( SO(12) \) lattice, which is realized at the free fermionic point in the moduli space, differs from the \( Z_2 \times Z_2 \) orbifold on \( T_2^1 \times T_2^2 \times T_2^3 \), which gives \( (h_{11}, h_{21}) = (51, 3) \). In \cite{ref12} it was shown that the two models may be connected by adding a freely acting twist or shift. Denoting the three complex coordinates of the \( T_2^1 \times T_2^2 \times T_2^3 \) tori by \( z_1, z_2 \) and \( z_3 \). Acting with the \( \{\alpha, \beta\} = Z_2 \times Z_2 \) twists on this space produces a model with 48 twisted fixed points, 16 from each of the twisted sectors \( \alpha, \beta \) and \( \alpha \beta \). The resulting manifold has \( (h_{11}, h_{21}) = (51, 3) \), and is denoted as \( X_1 \). The additional freely acting shift \( \gamma : (z_1, z_2, z_3) \rightarrow (z_1 + 1/2, z_2 + 1/2, z_3 + 1/2) \) produces again fixed tori from the three twisted sectors \( \alpha, \beta \) and \( \alpha \beta \) and does not produce any additional fixed tori. Under the action of the \( \gamma \)-shift, the fixed tori from each twisted sector are paired. Therefore, \( \gamma \) reduces the total number of fixed tori from the twisted sectors by a factor of 2, yielding \( (h_{11}, h_{21}) = (27, 3) \). This model therefore reproduces the data of the \( Z_2 \times Z_2 \) orbifold at the free-fermion point in the Narain moduli space. The precise form of the shift that reproduces the \( SO(12) \) lattice, and hence the \( Z_2 \times Z_2 \) at the free fermionic point is discussed in ref. \cite{ref12}. However, all the models that are obtained from \( X_1 \) by a freely acting \( Z_2 \)-shift have \( (h_{11}, h_{21}) = (27, 3) \) and hence are connected by continuous extrapolations.

The connection between \( X_1 \) and \( X_2 \) by a freely acting shift has profound consequences. The result of adding the freely acting shift \( \gamma \) is that the new manifold \( X_2 \), while still admitting three twisted sectors, is not simply connected and hence allows the breaking of the \( SO(10) \) symmetry by utilizing the Hosotani–Wilson symmetry breaking mechanism \cite{ref13}. Thus, we can regard the utility of the free fermionic machinery as singling out a specific class of compactified manifolds. In this context the freely acting shift has the crucial function of connecting between the simply connected covering manifold to the non-simply connected manifold. Precisely such a construction has been utilized in \cite{ref14} to construct non-perturbative vacua of heterotic M-theory.

### 4 M–embeddings

The profound new understanding of string theory that emerged over the past few years means that we can use any of the perturbative string limits, as well as eleven dimensional supergravity to probe the properties of the fundamental M–theory vacuum. The pivotal property that this vacuum should preserve is the \( SO(10) \) embedding of the Standard Model spectrum. This inference follows from the fact that also in the strong coupling limit heterotic M–theory produces discrete matter and gauge representations. Additionally, the underlying compactification should allow for the breaking of the \( SO(10) \) gauge symmetry. In string theory the prevalent method to break the \( SO(10) \) gauge group is by utilizing Wilson line symmetry breaking. Compactification
of M–theory on manifolds with $SU(5)$ GUT gauge group that can broken to the Standard Model gauge group were discussed in [14]. In [15], the analysis was extended to $SO(10)$ GUT gauge group that can be broken to $SU(5) \times U(1)$. This work was reviewed in [3] and here I discuss relevant points for further explorations of the phenomenological free fermionic models.

The key to the construction of ref. [14] is the utilization of elliptically fibered Calabi–Yau threefolds. These manifolds are represented as a two dimensional complex base manifold and a one dimensional complex fiber with a section. On these manifolds the equation for the fiber is given in the Weierstrass form

$$y^2 = x^3 + f(z_1, z_2)x + g(z_1, z_2) = (x - e_1)(x - e_2)(x - e_3).$$

Here $f$ and $g$ are polynomials of degrees 8 and 12, respectively and are functions of the base coordinates; $e_1$, $e_2$ and $e_3$ are the three roots of the cubic equation. Whenever two of the roots coincide the fiber degenerates into a sphere. Thus, there is a locus of singular fibers on the base manifold. These singularities are resolved by splitting the fiber into two spherical classes $F$ and $F - \mathbb{N}$. One being the original fiber minus the singular locus, and the second being the resolving sphere.

A nonperturbative vacuum state of the heterotic M–GUT–theory on the observable sector is specified by a set of M–theory 5–branes wrapping a holomorphic 2–cycle on the 3–fold. The 5–branes are described by a 4–form cohomology class $[W]$ satisfying the anomaly–cancellation condition. This class is Poincaré–dual to an effective cohomology class in $H_2(X, \mathbb{Z})$ that can be
written as
\[ W = c_2(TX) - c_2(V_1) - c_2(V_2) = \sigma_*(w) + c(F - N) + dN, \]
where \( c_2(TX), c_2(V_1) \) and \( c_2(V_2) \) are the second Chern classes of the tangent bundle and the two gauge bundles on the fixed planes; \( c, d \) are positive definite integers, \( \omega \) is a class in \( B \), and \( \sigma_*(\omega) \) is its pushforward to \( X \) under \( \sigma \).

The key to the M–theory embedding of the free fermionic models is their correspondence with the \( Z_2 \times Z_2 \) orbifold. The starting point toward this end is the \( X_1 \) embedding manifold with \((h_{11}, h_{21}) = (51, 3)\). The manifold is then rendered non–simply connected by the freely acting involution and the methodology of ref. 14,15 can be adopted to construct viable M–theory vacua. The difference however is that now the fiber is more singular than the ones previously considered. The fiber of \( X_1 \) in Weierstrass form is given by
\[ y^2 = x^3 + f_8(w, \tilde{w})xz^4 + g_{12}(w, \tilde{w})z^6, \]
where
\[ f_8 = \eta - 3h^2, \quad \text{and} \quad g_{12} = h(\eta - 2h^2), \]
\[ h = K \prod_{i,j=1}^4 (w - w_i)(\tilde{w} - \tilde{w}_j) \]
and
\[ \eta = C \prod_{i,j=1}^4 (w - w_i)^2(\tilde{w} - \tilde{w}_j)^2. \]
Taking \( w \to w_i \) (or \( \tilde{w} \to \tilde{w}_i \)) we have a \( D_4 \) singular fiber. These \( D_4 \) singularities intersect in 16 points, \((w_i, \tilde{w}_j), i, j = 1, \ldots, 4\), in the base. The resolution of the singular fiber in this case is more involved than the simpler ones previously considered. It is expected that the richer structure of fiber classes will yield a richer class of M–theory vacua with the possibility of new features appearing.

Figure 3 illustrates qualitatively the approach to the phenomenological application of M–theory advocated in this paper. In this view the different perturbative M–theory limits are used to probe the properties of a specific class of compactifications. In this respect one may regard the free fermionic models as illustrative examples. Namely, in the heterotic limit this formulation highlighted the particular class of models that are connected to the \( Z_2 \times Z_2 \) orbifold. In order to utilize the M–theory advances to phenomenological purposes, our task then is to now explore the compactification of the other perturbative string limits on the same class of spaces, with the aim of gaining further insight into their properties. In this spirit compactifications of type I string theory on the \( Z_2 \times Z_2 \) orbifold that are connected to the free fermionic models have been explored16.
5 Conclusions

The Standard Particle Model data suggests two pivotal ingredients that should be accommodated in the vacuum of the fundamental quantum gravity theory. The existence of three chiral generations and their embedding in $SO(10)$ representations. String vacua based on $Z_2 \times Z_2$ orbifold compactifications that admit these requirements have been constructed. In these construction the free fermionic point in the moduli space plays an important role, and may be singled out due to the maximally enhanced symmetries generated at this point and its relation to the self–dual point under the T–duality. To go beyond the perturbative analysis the advances in M–theory has to be employed, that perhaps will tell us what is special about the $Z_2$ orbifold? In the heterotic limits of M–theory, the prevailing method to break the GUT gauge group is the Hosotani–Wilson symmetry breaking mechanism. A fascinating aspect of this symmetry breaking mechanism on topologically non–trivial manifolds is that it gives rise to ultra–massive meta–stable states that provide different candidates for explaining the cosmic ray events beyond the GSK cutoff. Developing the experimental and phenomenological tools to decipher these events will be the subject of intense activity in forthcoming years.

Acknowledgments

I would like to thank Carlo Angelantonj, Dave Clements, Alessandro Caffarela, Claudio Coriano, Emilian Dudas, Richard Garavuso Jose Isidro, Marco Macone, Sander Nooij and Michael Plümacher for collaboration and discussions. Work supported in part by the Royal Society and PPARC.

References

1. K. Kawai et al, Nucl. Phys. B 288, 1 (1987); I. Antoniadis et al, Nucl. Phys. B 289, 87 (1987).
2. I. Antoniadis et al, Phys. Lett. B 231, 65 (1989); A.E. Faraggi et al, Nucl. Phys. B 335, 347 (1990); I. Antoniadis et al, Phys. Lett. B 245, 161 (1990); A.E. Faraggi, Phys. Lett. B 278, 131 (1992); Nucl. Phys. B 387, 239 (1992); Phys. Rev. D 46, 3204 (1992).
3. A.E. Faraggi, hep-ph/9707311, hep-th/9910042, hep-th/0208125.
4. A.E. Faraggi and D.V. Nanopoulos, Phys. Rev. D 48, 3288 (1993); A.E. Faraggi Int. J. Mod. Phys. A 14, 1663 (1999); hep-th/9511093.
5. J.L. Lopez and D.V. Nanopoulos, Phys. Lett. B 251, 73 (1990); I. Antoniadis et al, Phys. Lett. B 278, 257 (1992); Phys. Lett. B 279, 281 (1992); A.E. Faraggi, Nucl. Phys. B 403, 101 (1993); Nucl. Phys. B 407, 57 (1993); Phys. Lett. B 329, 208 (1994); A.E. Faraggi and E. Halyo, Phys. Lett. B 307, 305 (1993); Phys. Lett. B 307, 311 (1993); Nucl. Phys. B 416, 63 (1994); A.E. Faraggi and J.C. Pati, Phys. Lett.
B 400, 314 (1997); J. Ellis et al, Phys. Lett. B 425, 86 (1998); A.E. Faraggi and C. Coriano, hep-ph/0306186.

6. S. Chang et al, Phys. Lett. B 397, 76 (1997); Nucl. Phys. B 477, 65 (1996); C. Coriano et al, Nucl. Phys. B 614, 233 (2001).

7. S. Chaudhuri et al, Nucl. Phys. B 469, 357 (1996); G.B. Cleaver et al, Nucl. Phys. B 525, 3 (1998); Phys. Rev. D 59, 055005 (1999).

8. G.B. Cleaver et al, Phys. Lett. B 455, 135 (1999); Int. J. Mod. Phys. A 16, 425 (2001); Nucl. Phys. B 593, 471 (2001); Mod. Phys. Lett. A 15, 1191 (2000); Int. J. Mod. Phys. A 16, 3565 (2001); Nucl. Phys. B 620, 259 (2002).

9. A.E. Faraggi, Phys. Lett. B 274, 47 (1992); Phys. Lett. B 377, 43 (1995); Nucl. Phys. B 487, 55 (1996).

10. I. Antoniadis et al, Phys. Lett. B 302, 202 (1993); K.R. Dienes and A.E. Faraggi, Phys. Rev. Lett. 75, 2646 (1995); Nucl. Phys. B 457, 409 (1995). A.E. Faraggi, Nucl. Phys. B 428, 111 (1994); Phys. Lett. B 499, 147 (2001); Phys. Lett. B 520, 337 (2001); J.C. Pati, Phys. Lett. B 388, 532 (1996); J. Ellis et al, Phys. Lett. B 419, 123 (1998); I. Antoniadis et al, Phys. Lett. B 241, 24 (1990); A.E. Faraggi and J.C. Pati, Nucl. Phys. B 526, 21 (1998); A.E. Faraggi and O. Vives, Nucl. Phys. B 641, 93 (2002).

11. Y. Hosotani, Phys. Lett. B 129, 1983 (1993).

14. R. Donagi et al, Adv. Theor. Math. Phys. 5, 93 (2002).

15. A.E. Faraggi et al, Nucl. Phys. B 641, 111 (2002); hep-th/0209245. A.E. Faraggi and R. Garavuso, Nucl. Phys. B 659, 224 (2003).

17. C. Coriano and A.E. Faraggi, Phys. Rev. D 65, 075001 (2002); S. Sarkar and R. Toldra, Nucl. Phys. B 621, 495 (2002); C. Barbot and M. Drees, Phys. Lett. B 533, 107 (2002); A. Cafarella et al, hep-ph/0306236.