NEW TRENDS IN PARTICLE THEORY *

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Abstract

I discuss some new trends in Particle Theory beyond the Standard Model. Some topics which are briefly covered include electroweak baryogenesis, gauge versus non-gauge discrete symmetries, the strong-CP problem, dynamical symmetry breaking scenarios, supersymmetric grand unification, new aspects in low-energy supersymmetry and superstring phenomenology.

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1 Introduction

The official title of this talk is probably too ambitious. I guess nobody knows what are the future trends of Particle Theory. Instead of that I will just review some work done in the general field of Physics Beyond the Standard Model in the last couple of years or so. I will skip here the standard praising of the Standard Model (SM) and the immediate list of reasons why there should be something else Beyond the Standard Model (BSM). Instead of that let me display for you the present situation of the Stock Market of BSM ideas:

- Susy phenomenology ↑
- Weak-scale baryogenesis ↑
- Astroparticle (solar ν’s, COBE…) ↑
- String phenomenology ↑
- $W_L - W_L$ scattering at LHC/SSC ↑
- Constraints on BSM from EW-loops ↑
- Non-commutative geometry models ↑
- Technicolor and ETC ↑
- Axions; Global symmetries ↓
- B-violation at high energies in SM ↓
- $t - \bar{t}$ condensates ↓
- Wormholes solving the Cosm.C. problem ↓
- Non-SUSY SU(5) ↓
- 17 KeV neutrino ↓↓

As in real life, this is a very speculative market and the declared tendency of each particular topic is not directly related to the intrinsic value of each idea and is probably even less directly connected to reality. Of course, there are many interesting topics which are not in the list and even those in the list cannot be thoroughly reviewed. Fortunately, some of the most interesting topics have been discussed by other speakers in this conference: Astroparticle physics (Silk), neutrino physics (Spiro, Petcov…), constraints on BSM physics from electroweak loops (Altarelli…), extra Z-bosons (Cvetic, Taxil) etc. Concerning neutrino physics I just would like to make a very trivial comment: we should not overemphasize neutrino mass estimates based on the see-saw mechanism(s). These type of estimates are purely qualitative and should only be taken as such. Let us now turn to review a few topics which have received much attention in the last couple of years or so.
2 B and L-number violation in the SM at high temperatures

The SM has baryon number $B$ and the three lepton numbers $L_i$ as accidental global $U(1)$ symmetries. They are good classical symmetries of the *minimal* SM but are violated by quantum effects associated to the electroweak interactions. Indeed, the currents associated to $B$ and $L_i$ in the SM have mixed anomalies with the $SU(2)_W$ gauge bosons, and hence these symmetries are not respected by quantum mechanical effects. It was 't Hooft [1] who first realized that $SU(2)_W$ non-perturbative (instanton) effects can give rise to $B$ and $L_i$-violating effective interactions. Numerically, the rate of these interactions is of order $|T|^2 \simeq e^{-S} \simeq e^{-4\pi/\alpha_W} \simeq 10^{-80}$, and hence negligible for all practical purposes. However, it turns out that at *high temperatures* the $B$-violating interactions are no longer suppressed [2]. There are classical Higgs and $SU(2)_W$ gauge-field configurations ("sphalerons" [3]) which interpolate between vacua with different $B$-(and $L_i$-) number. Thus at high temperatures ($T \geq M_W/\alpha^2$) $B, L_i$-violating interactions are unsuppressed. The above fact has two important (cosmological) consequences:

i) **Weak-scale baryogenesis**

It could well be that these electroweak effects could be the source of $B$-violation required to generate the primordial baryon/antibaryon asymmetry of the universe [4]. However, the other two ingredients [5] required for this generation, CP-violation and departure from thermal equilibrium, do not seem so easy to get within the minimal SM. Although the SM automatically has a source of CP-violation from the KM-phase, this turns out to be too small (for recent controversy about this point see ref.[6]). This can easily be cured by a modest extension of the SM including extra scalars (or by going to the SUSY-SM which has additional sources of phases). The third point, breakdown of thermal equilibrium, is the toughest to get. Here the nature of the electroweak phase transition is of the outmost importance. It turns out that, contrary to what one would naively expect, the character of this very fundamental phase-transition is poorly known. If the phase transition is first order, baryogenesis is feasible. Estimations suggest that one should not expect a first order phase transition in the minimal SM for Higgs masses $m_H \geq 60$ GeV. Taking into account the LEP data, this leaves little space for a first order phase transition to develop. There are, however many theoretical uncertainties involved (see Fodor’s contribution to these proceedings) and one cannot rule out completely this possibility. On the other hand, extending the SM by adding a non-minimal Higgs sector makes life easier concerning baryogenesis.
Erasure of a primordial baryon asymmetry

It could well be that the low-temperature baryogenesis scenario outlined above does not work, in which case one could look back to the more traditional high-temperature baryogenesis scenarios which were so popular in the eighties [7]. Those schemes take the $B$ and CP-violation from explicit couplings present in most Grand Unified Theories (GUTs) like $SU(5)$ or $SO(10)$. Departure from thermal equilibrium is in this case very easy to obtain from late decay of superheavy particles (e.g. coloured scalars) generically present in GUTs. These scenarios are not free of problems either: inflation may completely dilute any primordial baryon asymmetry created in this way unless the reheating takes place at temperatures $\leq 10^{13}$ GeV. In the last few years it has also been realized that the high temperature electroweak effects discussed above may also erase any primordial $B$ asymmetry, since those effects are in thermal equilibrium below temperatures $T \leq 10^{12}$ GeV or so. There is a loop-hole though in this argumentation. If a net $B-L$ density is generated at a primordial stage, electroweak effects will be unable to erase it. This is because the combination $B-L$ has no mixed anomaly with $SU(2)_W$, and hence this symmetry is respected by all standard model interactions [2]. Thus the idea is to generate a primordial $B-L$ density in some GUT scenario like $SO(10)$ (minimal $SU(5)$ does not work because it has an exact global $B-L$ symmetry). High temperature electroweak effects may partially convert a baryon asymmetry into a lepton asymmetry (or viceversa) but are unable to erase the net $B-L$ [8].

If one takes this option of a primordial $B-L$ generation, one may still be in trouble if there are additional explicit interactions in the effective Lagrangian violating $B$ and/or $L$. Even if these additional interactions are very tiny or even suppressed by inverse powers of large masses (like GUT scale or Planck mass), they may be sufficiently effective in erasing the primordial asymmetry. If one insists in preserving this asymmetry one can obtain constraints on $B$ and/or $L$-violating terms like e.g., Majorana masses for the left-handed neutrinos. One finds for any of the three neutrinos the bound [9]

$$m_{\nu_i} \leq 10^{-3} \text{ eV} \quad (1)$$

if the primordial $B$-asymmetry is to be preserved. Other stringent limits are also found for other $B$ and/or $L$-violating terms [10] present in different models (e.g. $R_p$-violating operators in the SUSY-SM). If the above limits are strict, this would mean bad news for $\nu$-oscillation experiments (or high-temperature baryogenesis!).

More recently it has been realized that there are different effects which may somewhat relax these limits:

- In the presence of supersymmetry, due to the existence of a new approximate global $U(1)$ symmetry (R-symmetry) beyond $B$ and $L$, electroweak effects are again unable to erase a primordial $B$-asymmetry [11]. Since SUSY is not an exact symmetry, its only effect is to postpone the electroweak erasing down to temperatures $T \leq 10^7$ GeV or so (instead of $T \leq 10^{12}$ GeV). This is sufficient
to relax e.g. the $\nu$-mass bounds to $m_\nu \leq 10$ eV. This allows for more brilliant prospects for neutrino oscillations.

ii) Fermion mass effects on the baryon number densities may also be important. It has recently been pointed out that those may also avoid the erasure of a primordial asymmetry (see Dreiner’s contribution to these proceedings).

iii) The smallness of the electron Yukawa coupling makes that the $e_R$ gets very late into thermal equilibrium. This effect turns out to relax also the above type of bounds drastically [80].

Both in the high- and low-temperature scenarios for baryogenesis, it is clear that one cannot neglect the $B/L$-violating electroweak effects. Imposing that a given BSM scheme is consistent with either one or the other of these scenarios may be a extremely effective constraint on physics beyond the standard model. For example, minimal $SU(5)$ cannot yield any baryon asymmetry since it will necessarily be diluted by electroweak effects. Other mechanisms to generate baryon asymmetry may be at work, one of the most interesting ones being the one in ref.[12].

A final speculative remark is in order. I mentioned above that $B - L$ is an anomaly-free global symmetry. This is only partially correct: it does not have mixed anomalies with SM interactions but it does have mixed gravitational anomalies. Thus one expect the existence of gravitational effects violating $B - L$. One could speculate [13] on the possibility that at extremely high temperatures, just below the Planck mass, these gravitational effects could generate a $B - L$ primordial asymmetry. The usual electroweak effects would just redistribute the relative amounts of $B$ and $L$. This scenario would have the beauty that the very existence of an asymmetry would be a direct consequence of the anomaly structure of the SM. But, although attractive in principle, it is not obvious how to make such an scenario to work out in detail.

3 The Crisis of Global Symmetries

We mentioned above the four global $U(1)$ symmetries of the SM ($B$ and $L_i$). They are accidental symmetries of the theory and are a mere consequence of $SU(3) \times SU(2) \times U(1)$ gauge invariance, Lorentz invariance and renormalizability. In going to physics beyond the standard model one normally needs to impose new global symmetries to achieve different phenomenological goals. Examples of those are: 1) The Peccei-Quinn $U(1)_{PQ}$ symmetry, introduced to solve the strong-CP problem; 2) The discrete global symmetries introduced in multi-Higgs models in order to avoid flavour-changing neutral currents (FCNC); 3) Discrete (or continuous) "horizontal" symmetries introduced in order to get appropriate "textures" for the fermion mass matrices; 4) The usual $Z_2$ R-parity (or other "generalized matter parities") imposed in the minimal supersymmetric standard model (MSSM) in order to avoid fast proton decay.

All these global symmetries are really imposed by hand and, unlike gauge
symmetries, they are not motivated by any fundamental symmetry principle. In the last few years it has been realized that a number of gravitational effects badly violate global symmetries: terms which are forbidden from a Lagrangian by imposing a global symmetry (continuous or discrete) are regenerated by gravitational dynamics (wormholes, blackholes [14]). Thus global symmetries are not effective in fulfilling their expected duties! This is what I call in this section The Crisis of Global Symmetries.

A good example of the effect of this ”crisis” is the difficulties which are expected for a Peccei-Quinn type of solution for the strong-CP problem. This solution [15] requires the existence of a global $U(1)_{PQ}$ symmetry which has mixed anomalies with QCD. This symmetry is spontaneously broken, giving rise to the corresponding (pseudo-)Goldstone boson, the axion. Due to the anomaly, the axion field $a$ behaves as an ”effective $\theta$ parameter” which couples to gluons as $a/f_a G^\mu\nu \tilde{G}_{\mu\nu}$. Non-perturbative QCD effects generate a scalar potential for $a$ of the form

$$V_{\text{QCD}}(a) = (m_a^{QCD})^2 f_a^2 (1 - \cos \bar{a})$$

where $\bar{a} = a + \theta$. This potential is minimized for, $<\bar{a}> = 0$, yielding an elegant solution to the strong CP-problem (the problem of the unexpected smallness of the QCD $\theta$ parameter). Let us now assume that there are new operators of dimension $D = 4 + n$ which explicitly violate the $U(1)_{PQ}$ global symmetry. As I discussed above, gravitational effects are expected to generate such terms. In this case there will be additional contributions to the axion scalar potential. A simple estimation leads to a contribution [16]

$$V_{\text{grav}}(a) = (m_a^{\text{grav}})^2 f_a^2 (1 - \cos (n\bar{a} + \delta))$$

where $\delta$ is a number of order one. This potential is no longer minimized for $<\bar{a}> = 0$ and hence, in order not to spoil the solution provided by eq.[2], $m_a^{\text{grav}}$ has to be very small. In order to keep a $\theta$ parameter $\leq 10^{-10}$, as required by experiment, one needs to have

$$(m_a^{\text{grav}})^2 \leq 10^{-10} (m_a^{QCD})^2$$

A naive estimation tells us that in order for this contribution to be so supressed, the operator of lowest dimension $D$ violating $U(1)_{PQ}$ needs to have $D \geq 12$!! Thus an appropriate (gauge) symmetry should thus guarantee the absence of operators with $D \leq 12$ for the Peccei-Quinn mechanism to work. This looks like a bit hard to get.

There are in the literature other proposed non-axionic solutions to the strong-CP problem:

i) Assume that the laws of nature are CP conserving ($\theta_{\text{bare}} = 0$) and that this symmetry is spontaneously broken [17]. Once this symmetry breaking occurs, a non-vanishing (calculable) $\theta$ will appear and , hence one has to cook carefully the model in order to get small loop corrections to $\theta$. Although this is certainly a possibility, the actual realizations of this general idea are quite contrived.
ii) If the "current" u-quark mass vanishes, there is a global chiral $U(1)$ symmetry which allows to "rotate away" the $\theta$ angle. This possibility [18] is very neat and simple but hard to reconcile with the standard lore of chiral Lagrangian estimations of light-quark masses. Given the simplicity of this possibility I think it should however be seriously reconsidered.

Notice that these two alternatives also need the existence of global symmetries (CP in the first case, a chiral $U(1)$ in the second). Thus these two alternatives are also jeopardized by generic gravitational effects!

The case of the proposed solutions to the strong-CP problem is just an example. There are other BSM schemes which make use of discrete symmetries in a fundamental way and are also in trouble. A second prominent example is the R-parity discrete symmetry of the minimal supersymmetric standard model (MSSM). Unlike what happens in the simple standard model, in its SUSY version $B$ and $L_i$ are *not* automatic accidental symmetries of the theory. In particular, the most general Lagrangian consistent with SUSY, $SU(3) \times SU(2) \times U(1)$ and Lorentz invariance allows for Yukawa couplings which violate $B$ and/or $L_i$ symmetries. This would be phenomenologically catastrophic and hence one imposes by hand some *global* discrete symmetry which forbids the dangerous couplings. The simplest example of such a symmetry is "R-parity", a $Z_2$ symmetry under which the usual SM particles are even and all their SUSY-partners are odd. This symmetry forbids all $B$ and $L_i$-violating couplings. In view of our previous discussion, a global symmetry is not enough protection and the MSSM would be in trouble.

### 4 A New Guiding Principle?

There is a way out to the above "global symmetry crisis". If the symmetries which are phenomenologically required are *gauge symmetries*, they will be immune to the problematic gravitational effects. It is well known that those effects cannot e.g. violate charge conservation which is a symmetry associated to a gauge theory (QED). The intuitive reason for this difference between global and local is simple. The most important characteristic of *local* symmetries is that they play an important role in fixing what are the *actual physical degrees of freedom* of a theory. Local symmetries not only tell us what terms are allowed or forbidden in the effective Lagrangian (this is also done by the global symmetries) they also allow us to get rid of spurious non-physical states in the theory. Gravitational (or any other) perturbative or non-perturbative effects may generate terms violating a global symmetry but they can never modify the number of physical degrees of freedom of a theory and, hence, they cannot violate a gauge symmetry. This suggest to impose the following physical principle:

"*All symmetries (even discrete ones) should be gauge symmetries (unless they are accidental)*"

By accidental symmetries I mean symmetries like baryon and lepton numbers,
which are a mere consequence of renormalizability and gauge invariance. This could be a good possibility to obtain approximate Peccei-Quinn $U(1)_{PQ}$ symmetries in specific models [19]. Notice that Peccei-Quinn symmetries are anomalous and, hence, cannot be gauged in a straightforward way. It is interesting to remark, though that in string models PQ symmetries may be gauged under certain conditions [20].

The above principle tells us that even discrete symmetries should be gauged. Many particle theorists are not familiar with the concept of discrete gauge symmetry. The most intuitive way to generate a discrete gauge symmetry (DGS) is to start with a standard gauged $U(1)$ theory and break that symmetry spontaneously through the vev of a scalar field [21]. If the scalar field has charge $Nq$ and the rest of the particle spectrum has charges $q_i = M_i q$, with not all $M_i$ equal to a multiple of $N$, one can check that there is an unbroken $Z_N$ subgroup of the original $U(1)$ which remains unbroken. Locally, there is no way to distinguish a gauged from a global discrete symmetry but gauge symmetries give rise to some non-local interaction effects of the Bohm-Aharanov type which are not present in the global symmetry case [22]. These effects may have only, at most, cosmological relevance. However, there are other two important practical differences compared to the global case. One of them we already mention, DGS are immune to destabilizing gravitational effects. The other is that, just like it happens with gauged $U(1)$ symmetries, discrete symmetries should be anomaly-free [23]. The discrete charges of chiral fermions should obey certain restrictive discrete anomaly cancellation conditions. These conditions look very much like discretized versions of usual anomaly cancellation conditions. For example, the discrete $Z_N$ anomaly cancellation conditions look like [23]

$$\sum_i (q_i) = 0 \mod N \tag{5}$$

where $\exp(i 2\pi q_i / N)$ is the $Z_N$ charge of each of the $SU(M)$ fermion $M$-plets.

The discrete anomaly cancellation conditions should be obeyed by any DGS and this may lead to interesting phenomenological implications. For example the discrete $Z_N$ symmetries guaranteeing sufficient proton stability in the SUSY standard model ("generalized matter parities") were classified in ref.[23]. It was found that only four of them (a $Z_2$ and three $Z_3$s) are discrete anomaly-free with the particle content of the MSSM. The usual R-parity is one of them. Of course, this anomaly-freedom criterion may be applied to many other BSM schemes like, for example, discrete symmetries giving rise to appropriate quark-mass matrix textures.

It is also worth remarking that string theory does not have much sympathy for global symmetries either. In fact, there is a general theorem in strings which states that any exact continuous (e.g., $U(1)$) symmetry cannot be just global, it has to be a gauge symmetry [24]. An equivalent theorem for discrete (e.g. $Z_N$) symmetries has not been proved. However, it has been shown in explicit four-dimensional strings that many of the $Z_N$ symmetries have a gauge origin,
and that could well be the case for all discrete symmetries in string models. In particular, it has been shown that CP itself may be understood as a discrete gauge symmetry [25]. It has also been shown in plenty of four-dimensional string examples that the discrete anomaly cancellation conditions mentioned above are indeed satisfied [26]. The origin of this cancellation is still unclear but, like in the continuous gauge case, it is probably a consequence of the important "modular invariance" property of string theories.

5 The Naturality Problem: the Strongly Interacting Approach

By "naturality problem" I mean the problem of the instability of the Higgs sector of the standard model under quantum corrections. This problem goes under different names in the literature: gauge hierarchy problem, the mass problem etc. This is a problem which has concerned many particle theorists since more than fifteen years ago. Although a minority of physicists still maintain that this is not a problem since one can always renormalize the scalar mass to the value we wish, the immense majority think that there is indeed a problem since using physical (cut-off) regulators we need to make ridiculous fine-tunings to maintain the Higgs scalar sufficiently light to really induce $SU(2) \times U(1)$ breaking. As is well known, there are still essentially two schools of thought concerning this problem: i) The strongly interaction scenarios and ii) the supersymmetry approach. Let me say a few words about the first of these and postpone the second to the next section.

The strongly interacting schemes assume that, at energies not much above the weak scale, there are new strongly interacting phenomena which will allow us eventually to understand the origin of the symmetry-breaking (Higgs) sector of the standard model. In these schemes the Higgs field (and sometimes also the quarks and leptons) are composite particles. The main problem of this approach a priori is that very little is known about the non-perturbative physics of chiral gauge theories. Hence what one normally does is to imitate the physics of the only relatively well understood non-perturbative gauge theory, QCD, which is not chiral. This is what is done in the simplest and most attractive scenario of this type, Technicolor.

In Technicolor [27] one assumes that there are new QCD-like strong interactions at a scale of order a TeV with gauge group $G_T$. The theory contains fermions coupling to $G_T$ called technifermions $\Psi_i$. Instead of a Higgs vev, one assumes that there are non-vanishing vevs for technifermion bilinears, $\langle \bar{\Psi}_R \Psi_L \rangle \neq 0$, breaking $SU(2) \times U(1)$ spontaneously. All this works very nicely and the $W$ and $Z^0$ bosons indeed get a mass in the usual way. However, and this is the key problem, at this level the quarks and leptons remain massless. In order to provide masses to fermions, the best idea available is the introduction of extra new gauge interactions, Extended Technicolor (ETC) [28], whose crucial property is that they
connect the usual quarks and leptons with the technifermions. The latter are massive, they get a dynamical mass due to technicolor interactions. Then one can draw one-loop graphs in which a quark (lepton) splits into a techniquark and an ETC boson which then recombine again into a quark (lepton). These loops provide masses for the quarks and leptons of order

\[ m_{q,l} \simeq \frac{g^2_{ETC}}{(4\pi^2)} \frac{\langle \bar{\Psi}_R \Psi_L \rangle}{M^2_{ETC}} \]  

where \( g_{ETC} \) is the gauge coupling and \( M_{ETC} \) the mass scale of the new gauge ETC interactions. The idea of ETC is very nice in principle but very problematic in practice. The above formula has to provide masses for all quarks and leptons. It is very hard for such a one-loop effect to generate masses big enough to account for the masses of the third generation of quarks and leptons (particularly the t-quark). To increase the above contributions one has to either increase the value of the condensate \( \langle \bar{\Psi}_R \Psi_L \rangle \) or to decrease the ETC gauge boson masses \( M_{ETC} \). To increase the condensate without increasing at the same time the W and Z\(^0\) masses is not easy. The second possibility is also very problematic: the ETC gauge bosons necessarily change flavour leading to enormous FCNC unless \( M_{ETC} \geq 100 \) TeV or so. There is an extra problem for the ETC theories. Usually these theories not only produce dynamically the Goldstone bosons required for the Ws and Z\(^0\) to get massive, they also give rise to a plethora of other composite scalars (pseudo-Goldstone bosons), some of which should have already been seen at present accelerators like LEP. All these (and other) problems lead to a decline in the popularity of the Technicolor and ETC ideas during the years 1982-1988 or so.

There has been in the last few years a certain "discrete revival" of Technicolor and of the strongly-interacting Higgs sector schemes in general [29]. Concerning ETC, work has been done in trying to avoid the FCNC problem of these theories. In this connection, two main lines have been explored. The idea in both schemes is enhancing the value of the condensate \( \langle \bar{\Psi}_R \Psi_L \rangle \) without increasing the values of the masses of the W and Z\(^0\), which are fixed by the "technipion" decay constant \( F_{T\pi} \). One of the ideas to achieve this goes under the name of walking technicolor [30]. It was pointed out that, if the \( \beta \)-function of the Technicolor interactions is small, there is an enhancement of the ratio \( \frac{\langle \bar{\Psi}_R \Psi_L \rangle}{F^2_{T\pi}} \). It turns out however that the achieved enhancement is not enough to account for the mass of the third generation fermions. The second main idea put forward to increase this ratio is the use of ultraviolet fixed-point models [31]. The idea is that, if Technicolor interactions have an ultraviolet fixed point (a zero of the \( \beta(\alpha) \)-function) at a finite value of \( \alpha \), the \( \bar{\Psi}_R \Psi_L \) bilinear has large "anomalous dimensions", i.e. it gets a large enhancement factor. The problem is that it is not clear whether examples of non-Abelian theories of the above characteristics exist. It has been recently pointed out that non-Abelian gauge theories with a large number of fermions could have this property. Assuming that a Technicolor theory with the required properties exist, one can do some interesting model-building [32].
In spite of the efforts, no completely compelling ETC model with phenomenological promise has been constructed up to now. Furthermore, as I mentioned above, very special properties have to be assumed for a Technicolor theory to have any chance of surviving. On the other hand, these difficulties may well be due to our lack of understanding of non-perturbative gauge theories, and not really to the idea itself.

An alternative to Technicolor which has received attention in the recent past is based on the assumption of top-antitop quark condensation, as a substitute for techniquark condensation [33]. This is suggested by the fact that the top-quark is extraordinarily heavy compared to the other quarks. It is assumed that a bilinear condensate \( \langle \bar{t}_R t_L \rangle \neq 0 \) forms due to some unknown strong interactions giving rise to \( SU(2) \times U(1) \) breaking. In the original formulation of this idea, these unknown interactions were described by a Nambu-Jona Lasinio type of model which lead to some interesting predictions like \( m_{\text{Higgs}} = 2m_{\text{top}} \) and some more problematic results like \( m_{\text{top}} \geq 210 \text{ GeV} \). This latter result gives rise to large loop contributions to the \( \rho \)-parameter which are several standard deviations away from the electroweak data. There are also extra theoretical concerns: nothing is known about the origin of the masses of the rest of the quarks and leptons nor about the origin of the top condensation. It seems to me that the general idea of the t-quark condensates is much better than the actual implementations done up to date.

Due to our lack of knowledge of the precise strongly interacting dynamics which could be waiting for us above the weak scale, perhaps the wiser approach is trying to parametrize the process of \( SU(2) \times U(1) \) breaking in the most general possible way. In this connection, the most promising way seems to use the effective (chiral)-Lagrangian approach which is used so successfully in describing the chiral symmetry-breaking dynamics of QCD [34]. The effective chiral Lagrangian describing the \( SU(2) \times U(1) \) symmetry-breaking contains a definite set of operators [35] whose coefficients should be determined experimentally. Each specific model (e.g., ETC models, minimal SM, etc) corresponds to definite numbers for those coefficients. One of the most important experimental tests of a strongly-interacting Higgs sector would be the study at LHC/SSC energies of longitudinal W-boson scattering, \( W_L W_L \rightarrow V_L V_L \), where \( V_L = W_L, Z_L \). Since the longitudinal degrees of freedom of the massive gauge bosons correspond to the Goldstone boson, in a theory with a strongly interacting Higgs sector the cross section for \( W_L W_L \) scattering should be large. Scattering of \( W_L \)s should be accessible from \( W_L \) bremsstrahlung in \( p-p \) collisions at LHC/SSC. Detailed computations [36] show, however, that the rate for these reactions to be above background would normally require the existence of some resonance (e.g., a techni-\( \rho \)) in the \( W_L W_L \) channel.
6 The Naturality Problem: the Supersymmetric Option

In the last twelve years Supersymmetry (SUSY) has emerged as a serious alternative to avoid the naturality problem [37]. In this approach there are no new strong interactions and the Higgs sector is weakly interacting. This symmetry introduces a supersymmetric partner for each particle with opposite statistics and spin differing by 1/2 unit. SUSY thus transforms fermions into bosons and vice versa. The building blocks of a renormalizable SUSY field theory are the "chiral multiplets" and the "vector multiplets". A chiral multiplet \((\psi; \phi)\) contains a complex scalar \(\phi\) and a Weyl spinor \(\psi\) whereas a vector multiplet \((A^\mu; \lambda)\) contains a gauge boson \(A^\mu\) and its "gaugino" \(\lambda\) which is a Weyl spinor. The usual quark(\(q\)), leptons (\(l\)) and Higgs(\(H, \bar{H}\)) fields fit into chiral multiplets along with their SUSY-partners, the squarks(\(\tilde{q}\)) , sleptons(\(\tilde{l}\)) and Higgsinos (\(\tilde{H}, \tilde{\bar{H}}\)). The \(SU(3) \times SU(2) \times U(1)\) gauge bosons fit into vector multiplets along with their gauginos (gluinos, winos and bino). With these building blocks one easily builds a SUSY version of the SM. Everything works as in the usual field theory of the SM but with additional interactions involving the SUSY-partners. The number of coupling constants is the same as in the non-SUSY SM. But now, due to the presence of the additional partners and couplings, the Higgs mass parameters are stable under radiative corrections (they are not renormalized), providing a solution to the naturality problem.

Of course, SUSY cannot be an exact symmetry of nature and has to be broken in some way. It is the process of SUSY breaking which introduces additional parameters in the SUSY-SM. One may introduce terms in the Lagrangian which explicitly break SUSY but, in order not to get the stability of the scalar masses spoiled, these additional terms have to be of some restricted type (soft SUSY-breaking terms). This restricted type of terms are precisely the type of terms one obtains if SUSY (or better, its gauge version, Supergravity) is spontaneously broken in a "hidden sector " of the theory, but I will not elaborate on this point here. Let me just say that in this scheme the mass scale of the usual SUSY partners is fixed by the mass of the "gravitino", the SUSY partner of the graviton. The only phenomenologically important point is that certain SUSY-breaking soft terms appear now in the Lagrangian, including scalar masses, gaugino masses and some extra scalar interactions. In the simplest model, the Minimal Supersymmetric Standard Model (MSSM), there are only four SUSY-breaking soft terms:
i) Universal gaugino masses $M_{1/2}$  
ii) Universal scalar masses $M_0$  
iii) Trilinear scalar couplings proportional to $hM_0A$  
iV) Mixed Higgs mass term of the type $B\mu H\bar{H} + h.c.$.

Here $h$ is the corresponding Yukawa coupling and $\mu$ is a possible SUSY-preserving Higgs supermultiplet mass term which may in general be present in the original lagrangian. Thus in the MSSM the list of new SUSY parameters is:

$$M_{1/2}, M_0, A, B, \mu$$

All these couplings are universal (e.g., all scalar masses of all different squark, slepton and Higgs fields are equal) at a large mass scale (the grand unification or the Planck mass scales).

One of the most attractive features of the SUSY versions of the SM is that $SU(2) \times U(1)$ symmetry-breaking appears as a direct consequence of SUSY-breaking. One can see that, once the above soft terms are generated, loop effects generate a scalar potential for the Higgs fields which automatically induces $SU(2) \times U(1)$-breaking. At the unification scale all soft scalar masses are equal to $M_0$ but the low energy evolution computed through the renormalization group equations is different: the squark mass $^2$ increases at low energies, the sleptons and $\tilde{H}$ masses vary very little whereas the mass $^2$ of the other Higgs $\tilde{H}$ becomes negative! This is precisely what we need in order to spontaneously break $SU(2) \times U(1)$ and at the same time avoid color or charge symmetry breaking. It is important to remark that the above behaviour is very generic and is essentially a consequence of the multiplet structure and quantum numbers of the SUSY-SM. One point is however important for the mechanism to work out correctly: the $\tilde{H}$ Higgs field gets a negative mass $^2$ only if its coupling to the top-quark (the t-quark Yukawa coupling) is sufficiently large. Numerically, one essentially needs to have $m_t \geq 60$ GeV if the mechanism is to work in a natural way. When this mechanism was proposed these values for $m_t$ seemed fantastically large but nowadays we know that the top-quark mass is in fact very large and the radiative $SU(2) \times U(1)$ breaking mechanism is in fact a very natural one. Needless to say, in order to obtain the minimum of the Higgs potential at the correct scale (i.e. in order to reproduce the measured values of the $W$ and $Z^0$ masses), the full parameter space $M_{1/2}, M_0, A, B, \mu, h_{top}$ is constrained, but the correct numbers are obtained for very wide ranges of the above parameters.

In the last three years or so there has been a certain increase in the number works in the field of supersymmetry, particularly on the minimal supersymmetric standard model. Much of this sociological effect was motivated by the famous joining of the three gauge coupling constants at a single unification scale which, with the advent of the LEP data, became more striking [38]. Many of the SUSY topics studied in the early eighties were reconsidered and analized in more detail. Some of the SUSY topics recently reconsidered are the following:

i) Unification of gauge coupling constants
Indeed, when one runs up in energies [39] the three gauge coupling constants \(g_1, g_2, g_3\) using the renormalization group, the low energy data is consistent with unification at a single point \(M_X \simeq 10^{16}\) GeV [40]. In the case of \(g_1\), unification takes place if the standard GUT boundary condition \(g_1^2 = \frac{3}{5} g_{SU}^2\) holds. An equivalent way of stating the same fact is that, if there is unification of the gauge coupling constants into a standard GUT (\(SU(5), SO(10),\) etc), one can compute one of the gauge couplings in terms of the others. At the one-loop level and ignoring threshold corrections one gets the well known formulae

\[
\sin^2 \theta_W(M_Z) = \frac{3}{8} (1 + \frac{5\alpha(M_Z)}{6\pi} (b_2 - \frac{3}{5} b_1) \log(\frac{M_X}{M_Z}))
\]

\[
\frac{1}{\alpha_s(M_Z)} = \frac{3}{8} (\frac{1}{\alpha(M_Z)} - \frac{1}{2\pi} (b_1 + b_2 - \frac{8}{3} b_3) \log(\frac{M_X}{M_Z}))
\]

where in the MSSM one has \(b_1 = 11, b_2 = 1\) and \(b_3 = -3\). Since the experimental errors for \(\sin^2 \theta_W\) and \(\alpha\) are the smallest, what makes sense is to compute \(\alpha_s\) in terms of the other two. A number of refinements have been introduced for this computation in the last three years including [41]: i) Effect of superheavy GUT-thresholds. This can only be done in a specific GUT model like minimal SUSY \(SU(5)\). ii) Effect of the low energy sparticle thresholds. iii) Effect of two loop corrections involving the t-quark. iv) Possible corrections to the GUT values of the coupling constants due to non-renormalizable effects coming (presumably) from gravity. Including all these possible sources of uncertainties one finds the following result [41]

\[
\alpha_s(M_Z) = 0.125 \pm 0.01 = 0.125 \pm 0.001 \pm 0.005 \pm 0.0005 \pm 0.005 \pm 0.002 \pm 0.006
\]

The first two errors in the above expression come from the errors in the original input parameters \(\alpha(M_Z)\) and \(\sin^2 \theta_W(M_Z)\). The other four additional errors come from each of the four effects i)-iv) listed above. The result obtained for \(\alpha_s(M_Z)\) is in very good agreement with the experimental average \(\alpha_s(M_Z) = 0.12 \pm 0.01\). It is important to remark that in the case of the non-SUSY unification one obtains the prediction \(\alpha_s(M_Z) = 0.075\), which is clearly ruled out.

ii) \(m_b/m_\tau\) and the mass of the top quark

In many grand unified theories like \(SU(5)\) or \(SO(10)\) the Yukawa couplings of the b-quark and the \(\tau\)-lepton are equal at the unification scale \(M_X\) [42]. At low energies these couplings get renormalized differently and in the SUSY case [43] one then gets a predicted ratio of the type [44]:

\[
\frac{m_b(M_Z)}{m_\tau(M_Z)} = \left(\frac{\alpha_s(M_Z)}{\alpha_s(M_X)}\right)^{8/9} \left(\frac{\alpha_1(M_Z)}{\alpha_1(M_X)}\right)^{10/99} \left(1 + \frac{6}{(4\pi)^2} h_{top}^2 F(M_Z)\right)^{-1/12}
\]

where \(F(M_Z) \simeq 290\) and \(h_{top}\) is the t-quark Yukawa coupling. The first two factors come from renormalization due to gluon and hypercharge boson exchange, whereas the third factor comes from loops involving a virtual top quark. Of
course, in order to compare with the physical masses of $b$ and $\tau$, one has to run $m_b$ down to half the upsilon mass. In the early eighties the preferred values for the top quark mass were relatively small and hence, the last factor in (10) was usually ignored. Now we know that $m_{\text{top}}$ is large and hence that term cannot in general be neglected. Notice that, if we knew $m_b$, $m_\tau$, $\alpha_s$ and $\alpha_1$ with very good precision, we should be able to extract what is the value of the top Yukawa coupling [44]. With $h_{\text{top}}$ so determined, one gets a relationship between $m_{\text{top}}$ and $tg\beta$ since both are related by

$$m_{\text{top}} = h_{\text{top}}(\frac{\sqrt{2}M_W}{g_2})sin\beta.$$  

(\beta is defined by $tg\beta = <\bar{H}> / <H>$). There have been a number of recent analysis of the $m_{\text{top}}$ dependence of the $m_b/m_\tau$ ratio [45]-[47]. The results are very sensitive to the value taken for $\alpha_s$ and also on the value of $m_b$. If one takes for $\alpha_s$ the value given in eq.[9], impose the "experimental" condition $0.85m_b^0(5GeV) \leq 4.45GeV$ [41] and takes the LEP constraints on the top mass $120GeV \leq m_{\text{top}} \leq 160GeV$, one finds that the region $3 \leq tg\beta \leq 40$ would be forbidden. However, this bound dissapears for values of $\alpha_s$ slightly smaller than 0.120. Furthermore, slight corrections to the GUT identity $h_b = h_\tau$ which may come from a variety of sources close to the GUT scale make also the bound on $tg\beta$ to disappear (see S. Pokorski contribution to the parallel session).

One interesting point recently analized is whether a GUT boundary condition [48] equating all third generation Yukawa couplings $h_{\text{top}} = h_b = h_\tau$ is consistent both with data and the theoretical constraints. This type of relationship appears in simple SO(10) models and has the virtue of reducing the number of free parameters in explicit models of fermion masses. The answer [41],[49] seems to be yes, however one needs to have $m_{\text{top}} \simeq 180 - 190$ GeV, the fixed point value. Furthermore, very large values for $tg\beta \simeq 50$ are required, which is very unnatural to get in a radiative $SU(2) \times U(1)$ breaking model [50].

iii) Ansätze for quark and lepton mass matrices

Another topic which has received new attention is the construction of predictive ansätze for fermion mass matrices within the context of SUSY-GUTs. Some simple ansätze with "texture zeros" lead to attractive predictions for masses and mixing angles like e.g., $|V_{us}| = \sqrt{m_d/m_s}$ or $|V_{cb}| = \sqrt{m_c/m_t}$. The subject has been nicely reviewed in the paralell sessions by S.Raby and G.Ross.

iv) SUSY-Higgs masses

This topic is of very direct phenomenological relevance [51]. The issue which has been reconsidered in this case is the validity of the theoretical upper bounds on the lightest SUSY neutral Higgs scalar. At the tree level in the MSSM there is always a neutral scalar which is necessarily lighter than the $Z^0$ mass. On the other hand, for a heavy top quark (which is the experimental case) the loop
corrections to the masses of the scalars give large contributions of order
\[ \delta m^2 \simeq \frac{3}{8\pi^2} \frac{g^2 m_{\text{top}}^4}{\sin^2 \beta M_W^2} \log(1 + \frac{m_{\tilde{q}}^2}{m_{\text{top}}^2}). \] (12)

One then finds that the lightest Higgs scalar has to be lighter than something like 130 GeV for values of \( m_{\text{top}} \) of order 180 GeV and squark masses around one TeV. Thus, unfortunately, there is no guarantee that LEP-II will be enough to check the Higgs sector of the MSSM. Another important issue about the SUSY-Higgs sector is whether the combined data from both LEP-II and LHC will be enough to probe it. If one draws a plot of one the neutral Higgs mass (e.g., that of the pseudo-scalar \( A \)) versus \( \tan \beta \), one finds a certain window for \( 5 \leq \tan \beta \leq 20 \) and \( 100 \text{ GeV} \leq m_A \leq 200 \text{ GeV} \) in which the rate (and/or signature) for Higgs-particle production is too small to be detectable. It is an important challenge to look for interesting signatures to close this window and some ideas on how to close it have already been put forward [52].

v) Other SUSY-topics

Many other areas of SUSY standard model phenomenology have been recently reanalyzed. Amongst those the following: a) SUSY proton decay. The decay rate coming from dimension five operators in minimal SUSY-SU(5) is close to the experimental limits on nucleon instability [53]. Some authors even claim one can already constraint the masses of SUSY-particles (or \( \tan \beta \)) using those limits (see talks in the parallel sessions by Arnowitt and Nath); b) Upper bounds on the masses of SUSY-particles from naturallity arguments (see talks by Arnowitt, Nath and Ross); c) Constraints on MSSM parameters in order to get appropriate amount of dark matter in the form of lightest stable neutralinos (see talks by Roszkowsky and Ross); d) SUSY contributions to the decay \( b \rightarrow s\gamma \) etc.

It is certainly intriguing how the MSSM has passed a number of important tests in the last decade. I find particularly significant the joining of the three coupling constants at a single point and also that within the MSSM a heavy top quark leads in a natural way to the spontaneous breakdown of \( SU(2) \times U(1) \). The other merits of SUSY which are occasionally mentioned (like e.g., correct \( m_b/m_\tau \) por large \( m_{\text{top}} \); correct amount of dark matter predicted; not too much proton decay etc.) are more model-dependent. One has to remark also that essentially all these interesting points are present not only in the MSSM but also in simple extensions like models with R-parity violation and models with an extra singlet Higgs scalar (sometimes called the NMSSM =next to MSSM).

7 Challenges for Supersymmetric Unification

Not everything is nice and simple within the realm of the SUSY standard model. There are a good number of issues which still need to be understood. I will briefly mention here the four problems which look more relevant to me. The first three
are the following: a) In the SUSY standard model the soft terms which break SUSY are in general complex. This gives rise to new sources of CP-violation beyond the KM-phase. In particular, one finds that, unless the complex phase appearing in the soft terms are small (smaller than $10^{-2} - 10^{-3}$), one gets large contributions to the electric dipole moment of the neutron (EDMN) two or three orders of magnitude above experimental limits; b) In SUSY versions of the SM there are also new sources of flavour-changing neutral currents (FCNC). These are due to the exchange of sparticles in box diagrams and appear if e.g. the squarks are not degenerate in mass; c) In the models in which supersymmetry-breaking takes place in a hidden sector of the theory, one can have cosmological problems with gravitino and other singlet fields in charge of SUSY-breaking which can spoil standard nucleosynthesis if some stringent constraints on the masses of those particles are not obeyed [54].

I must say that the above three problems are interesting constraints on explicit models of SUSY-breaking but do not seem to me difficult to overcome. Indeed, there are different SUSY-breaking models in the literature in which the above points do not cause any trouble. On the other hand, in my opinion, the fourth problem that I want to mention is really serious. This is the famous doublet-triplet splitting problem which has been with us already for more than fifteen years. Let me remark from the outset that this is not really a problem of SUSY, but a problem of GUTs. Let us briefly recall what it is. Consider the simplest case of $SU(5)$. The Weinberg-Salam doublet $H_2$ which breaks the electroweak symmetry is contained in a five-plet of $SU(5)$ along with a triplet of coloured scalars $H_3$. The latter have to be superheavy, otherwise they would mediate very fast proton decay through dimension six operators. Thus we have to arrange the parameters of our $SU(5)$ Lagrangian in such a way that $H_2$ remains light (to be available for electro-weak symmetry breaking) but $H_3$ is superheavy (to avoid fast proton decay). These doublet-triplet splitting requires a fine-tuning of the Lagrangian parameters to one part in $10^{14}$!! Supersymmetry does not solve this problem, it just guarantees that, provided the fine-tuning is done, radiative corrections will not spoil it. Although some ideas in order to obtain the doublet triplet splitting without fine-tuning have been suggested (sliding singlet mechanism [55], missing partner mechanism [56], Higgses as Goldstone bosons [57] etc..) all of them either do not work (in the case of the first of them) or are really cumbersome and ad-hoc.

In my opinion this doublet-triplet splitting problem of GUTs is sufficiently important to consider seriously the possibility of giving up on GUTs (but not on SUSY!!). Is the GUT idea really needed? Supersymmetry is certainly enough to solve the hierarchy problem but we do not want to get rid of the nice properties of GUTs. Amongst those one of the nicest is charge quantization, i.e. the fact that $Q_e = 3Q_d$, which is automatic in GUTs like $SU(5)$ and $SO(10)$. In fact this charge quantization property does not require any GUT. It is well known [58] that it may be equally obtained if one imposes anomaly cancellation on one family of quarks and leptons. Thus we really do not need unification to get that one. The second very nice property of standard GUTs is the prediction of gauge coupling
unification which occurs precisely for the GUT scale value $\sin^2\theta_W = 3/8$. This nice prediction is not easy to obtain in a non-unified theory. One interesting alternative is to consider string theories which have the remarkable property of gauge coupling unification even in the absence of a unification group. This brings us to our last and most speculative subject.

8 Superstring Phenomenology

The most outstanding virtue of four-dimensional (4-D) supersymmetric heterotic strings [59] is that they are finite theories of quantum gravity, a property which is not true of any 4-D field-theory. Furthermore, they allow for chiral gauge interactions like the ones of the SM. Thus 4-D strings are the only known candidates for unified theory of all interactions. These are theories of closed strings which contain in their spectrum an infinity of particles (string excitations), all of them but a few with masses of order of the Planck mass. One identifies the massless states as candidates to describe the observed world. There is a unique type of string interaction (the merging of two closed strings into a single closed string) and this leads to identities between the gauge coupling constants and the Newton coupling constant, as we will mention below. An important property of string models is that the coupling constants are not constants but fields whose vacuum expectation values determine the physical values.

There are at present explicit examples of 4-D strings with three quark-lepton generations and a gauge group containing the $SU(3) \times SU(2) \times U(1)$ interactions of the standard model. All these models typically have extra particles (e.g., heavy leptons and/or quarks, extra gauge bosons). Although there is not at present a specific model which exactly mimics absolutely all the desired properties of, e.g., the SUSY-SM, they get tantalizingly close. Furthermore, these string models have some attractive generic features (e.g., hierarchies of Yukawa couplings; existence of a multiplicity of generations; constraints on the possible particle representations; possibility of gauging some anomalous $U(1)$s etc.) which open new avenues for model-building of unified theories. The construction of 4-D strings which resemble at low energies the SM and the study of possible generic features of unified models based on 4-D strings goes under the name of string phenomenology [60]. This is a vast field and I cannot discuss all the aspects of it. I will concentrate on two topics which have recently received some attention: SUSY-breaking soft terms and gauge coupling unification in 4-D strings.

One of the interesting features of string models is that they have natural candidates to constitute the "hidden sector" breaking supersymmetry in $N = 1$ supersymmetric models. In particular there are a couple of singlet scalar fields $S$ and $T$ which couple to usual matter only through non-renormalizable terms suppressed by inverse Planck mass powers [61]. The vev of these fields have a clear physical interpretation: $\langle ReS \rangle = 1/g^2$ determines the size of the string (i.e. gauge) coupling constant $g^2$ and $ReS$ is called the dilaton. Concerning
the other one has \(< ReT >= R^2\), where \(R\) is the overall size of the six extra compactified dimensions present in compactified string models. The field \(T\) is called the modulus. Of course it would be very interesting to compute from first principles those two vevs which would tell us what should be the size of the gauge couplings \(g^2\) and what is the size of the compactification radius \(R^2\). The latter should be of order one in Planck mass units. To find those vevs we would need to know the scalar potential \(V(S,T)\) and then look for its minima. Unfortunately this potential vanishes order by order in perturbation theory and hence \(ReS, ReT\) are undetermined at this level. Thus the expectation is that non-perturbative effects will raise that degeneracy and create a non-vanishing potential for \(S\) and \(T\). Indeed in simple models of "gaugino condensation" \([62]\) such non-vanishing potentials are in fact generated. In these models it is assumed that the gaugino fields \(\lambda\) corresponding to extra ("hidden sector") gauge factors normally present in string models condense (i.e. \(< \lambda\lambda >\neq 0\)) and this gives rise to SUSY-breaking and a non-vanishing scalar potential \(V(S,T)\) \([63],[64]\). For particular classes of hidden-sector gauge groups and particle content, the values obtained for \(< ReS >= 1/g^2\) are consistent with the extrapolation of the low energy gauge coupling constants \([65]\). Within this type of models one can also obtain results for the values of the SUSY-breaking soft terms \(M_{1/2}, M_0, A, B\) that we discussed above \([66]\).

There is some more general information about soft terms which can be obtained without resorting to specific gaugino condensation models. One can find simple expressions for the SUSY-breaking soft terms if one assumes that the two fields \(S,T\) play the leading role in providing the source for SUSY-breaking \([67]-[70]\). One does not need to know the precise way in which SUSY-breaking takes place to get those results. In this case the "goldstino" \(\tilde{\eta}\) (the Goldstone particle of SUSY-breaking) is predominantly a linear combination \(\tilde{\eta} = \sin\theta \tilde{S} + \cos\theta \tilde{T}\) of the fermionic partners of \(S\) and \(T\). For \(\sin\theta = 1\) the dilaton \(S\) is the leading source of SUSY-breaking whereas for \(\cos\theta = 1\) it is the modulus \(T\) which is dominant. Now, depending on the particular string model and the value of \(\sin\theta\) one gets different results for the soft terms \([70]\). The effective low energy Lagrangian of the model is determined by the standard \(N = 1\) supergravity functions which are the Kahler potential \(K(S,T,C_i)\) and the gauge kinetic function \(f^a, a\) being a gauge group label. The tree level form of \(f^a\) is universal for any string model, \(f^a = k^a S\), where the \(k^a\) are numerical constants to be discussed below. The kahler potential \(K\) is a model-dependent function. In a simple (but large) class of models (symmetric orbifolds) this function takes the general form \([71]\):

\[
K(S,T,C_i) = -\log(S + S^*) - 3\log(T + T^*) + +(T + T^*)^{n_i}C_iC_i^* \quad (13)
\]

where the \(C_i\) correspond to the matter fields, like (s)quarks, (s)leptons etc. The \(n_i\) are particle-dependent integers whose most common values are -1,-2,-3 for this class of models \([68]\). One can plugg this expression in the general form of the scalar potential and assume that \(S\) and \(T\) give the dominant source of SUSY-breaking. If one farther imposes that the cosmological constant should vanish
one finds expressions for the soft terms of the following type [70]:

\[ m_i^2 = n_{3/2}(1 + n_i \cos^2 \theta) \]
\[ M_{1/2} = \sqrt{3}m_{3/2}\sin \theta \]
\[ A_{ijk}^0 = -\sqrt{3}m_{3/2}(\sin \theta + (3 + n_i + n_j + n_k) \sqrt{3} \cos \theta) \]  

(14)

where the indices \( ijk \) refer to the three particles coupled through the \( A \)-term and the superindex 0 indicates that this is the \( A \) parameter of a large, non-suppressed Yukawa coupling (like the one expected for the top). In some classes of orbifolds (e.g. \( Z_2 \times Z_2 \)) and for the large \( T \) limit of Calabi Yau compactifications one has \( n_i = -1 \) and then one gets the even simpler boundary condition [69],[70]:

\[ M_{1/2} = -A = \sqrt{3} m_i \]  

(15)

Notice that in fact those boundary conditions also apply in the limit \( \sin \theta = 1 \) (dilaton dominance limit) for any \( n_i \) choice, i.e., in a model-independent way [69], since the dilaton dependence of the effective low energy Lagrangian is indeed model-independent. Although the above expressions for the soft terms involve some assumptions (most notably that \( S,T \) dominate SUSY-breaking and the imposition that the cosmological constant vanishes) it is remarkable how far one can go in obtaining equations based on string physics which may be experimentally tested. Indeed, expressions like the ones above lead to certain constraints on the low energy supersymmetric mass spectrum when run down in energies according to the renormalization group equations. Such type of analysis has been recently done in ref.[72],[70].

Another interesting topic in string phenomenology is the issue of gauge coupling unification. Suppose we had a 4-D string with gauge group including the one of the standard model. Then one finds [73]

\[ g_3^2k_3 = g_2^2k_2 = g_1^2k_1 = \frac{4\pi M_{string}^2}{M_{Planck}^2} \]  

(16)

Thus even in the absence of a GUT group there is unification of coupling constants. The \( k_a \) corresponding to the non-Abelian gauge factors are positive integers called the Kac-Moody levels. In practically all models considered up to now one has \( k_2 = k_3 = 1 \). In fact this is always the case for string models obtained upon direct compactification of the \( E_8 \times E_8 \) heterotic string. The corresponding constant \( k_1 \) for abelian factors like the hypercharge is a model-dependent rational number. Consistency with the low energy spectrum of the standard model requires \( k_1 \geq 1 \) [74] but otherwise there is no other model independent constraint on it. This leads to the model-independent constraint \( \sin^2 \theta_W(M_{string}) = k_2/(k_1 + k_2) = 1/(1 + k_1) \leq 1/2 \). The standard GUT value for \( k_1 \) is 5/3, leading to \( \sin^2 \theta_W(M_{GUT}) = 3/8 \), but in the non-unified string case \( k_1 \) needs not be equal to the GUT result [75]. Since the couplings are unified at
the string scale $M_{\text{string}} \simeq 4 \times 10^{17}$ GeV [76] and not at the previously mentioned SUSY-GUT scale $M_X$, instead of eqs.[7,8] one gets:

$$sin^2 \theta_W(M_Z) = \frac{1}{1 + k_1 \alpha(M_Z)} (1 + \frac{k_1 \alpha(M_Z)}{2\pi} (b_2 - \frac{b_1}{k_1}) \log(\frac{M_{\text{string}}}{M_Z}))$$

(17)

$$\frac{1}{\alpha_s(M_Z)} = \frac{1}{1 + k_1} \left( \frac{1}{\alpha(M_Z)} - \frac{1}{2\pi} (b_1 + b_2 - (1 + k_1)b_3) \log(\frac{M_{\text{string}}}{M_Z}) \right)$$

(18)

Let us assume for a moment that a) $k_1 = 5/3$ as happens e.g. in $E_6$ type of string compactifications; b) that the only particles charged under $SU(3) \times SU(2) \times U(1)$ are the ones of the minimal supersymmetric standard model and c) neglect possible string threshold effects. Then one finds [77],[78] $sin^2 \theta_W(M_Z) = 0.218$ and $\alpha_s(M_Z) = 0.20$, far away from the experimental results (but still better than non-SUSY $SU(5)$). Thus one has to give up at least one of the above three assumptions a) to c). One can give up assumption b) and consider further charged particles apart from those of the MSSM. If the extra introduced particles are not to be exotic this will require the introduction of some intermediate mass scale(s) between $M_{\text{Planck}}$ and the weak scale [78]. This opens a Pandora’s box of possibilities. Concerning assumption c), indeed string threshold effects can be relatively large, since infinite towers of particles may give contributions in the vicinity of $M_{\text{string}}$. String threshold effects have been computed in some class of orbifold models [79]. Phenomenological analysis shows [77],[68] that the string threshold corrections may be large as long as $ReT \simeq 4 - 20$. These corrections go in the good direction only for very restricted possible models. Finally, one can give up assumption a) and consider values of $k_1$ different from 5/3 [75]. The best agreement is found for $4/3 \leq k_1 \leq 3/2$. For $k_1 = 1.444$ one finds $sin^2 \theta_W(M_Z) = 0.233$ and $\alpha_s(M_Z) = 0.14$ ; for $k_1 = 1.466$ one gets $sin^2 \theta_W(M_Z) = 0.235$ and $\alpha_s(M_Z) = 0.137$. Given the inherent uncertainties coming from string threshold effects these results can be considered as succesfull, and show that claiming for the necessity of intermediate scales in direct string unification may be premature.

On the other hand specific string models with $k_1$ in the interesting range still have to be found. In any case it is clear that the constraint of gauge coupling unification in string models may be very important in selecting possible string unification schemes.

9 Conclusion

The general field of Particle Theory beyond the Standard Model is quite speculative and, at the moment, it is quite difficult to favour on a firm basis any of the avenues discussed above better than any other. It may well be that the directions of the arrows showing the tendency of different topics listed in the introduction flip in some cases upside down. Our main hope relies on the planned new experiments and it is also our hope that some of the ”trends” briefly described above will be promoted to real physics!
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