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Beautiful Mirrors and Precision Electroweak Data

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Abstract

The Standard Model (SM) with a light Higgs boson provides a very good description of the precision electroweak observable data coming from the LEP, SLD and Tevatron experiments. Most of the observables, with the notable exception of the forward-backward asymmetry of the bottom quark, point towards a Higgs mass far below its current experimental bound. The disagreement, within the SM, between the values for the weak mixing angle as obtained from the measurement of the leptonic and hadronic asymmetries at lepton colliders, may be taken to indicate new physics contributions to the precision electroweak observables. In this article we investigate the possibility that the inclusion of additional bottom-like quarks could help resolve this discrepancy. Two inequivalent assignments for these new quarks are analysed. The resultant fits to the electroweak data show a significant improvement when compared to that obtained in the SM. While in one of the examples analyzed, the exotic quarks are predicted to be light, with masses below 300 GeV, and the Higgs tends to be heavy, in the second one the Higgs is predicted to be light, with a mass below 250 GeV, while the quarks tend to be heavy, with masses of about 800 GeV. The collider signatures associated with the new exotic quarks, as well as the question of unification of couplings within these models and a possible cosmological implication of the new physical degrees of freedom at the weak scale are also discussed.

1 Introduction

The electroweak precision tests, driven primarily by the experiments at LEP, the Tevatron and the SLC, have, in recent years, held much of the attention of the field. Taken in conjunction with the measurement of the top mass and certain other low energy measurements, these experiments have vindicated the Standard Model (SM) to an unprecedented degree of accuracy [1]. While startling deviations from the SM expectations have occasionally appeared, only to disappear later as the precision increased, the results of the precision tests have been remarkably steady over the last five years. Yet, certain discrepancies persist. It is thus contingent upon us to examine their significance and especially to ascertain whether they could be pointers to new physics at the weak scale.

In this article, we shall concentrate upon the most obvious of such a possible deviation [2], namely the forward-backward asymmetry (A^b_{FB}) of the b-quark, the measurement of which shows a 2.9σ deviation from the value predicted by the best fit to the precision electroweak observables within the SM [1,3]. One might, of course, argue that this discrepancy is but a result of experimental inaccuracies and/or just a large statistical fluctuation. This viewpoint is supported, to some extent, by the observation that the corresponding SLD measurement of the b-asymmetry factor A_b using the LR polarized b asymmetry is in much better agreement with the SM [1]. It has also been argued that any correction to the $\bar{b}bZ$ vertex, large enough to 'explain' A^b_{FB} would have shown up in the very accurate measurement of R_b, the branching fraction of the Z into b's. However, we shall demonstrate that this need not be so. But more importantly, given the remarkable consistency amongst the four LEP experiments as regards A^b_{FB}, it is perhaps worthwhile to take this deviation from the SM seriously and to speculate on possible explanations thereof.

Let us begin by reviewing the relevant data at the Z-peak. We parametrize the effective $Zb\bar{b}$ interaction by

$$\mathcal{L}_{Zb\bar{b}} = \frac{-e}{s_W c_W} Z_\mu \bar{b} \gamma^\mu \left[\bar{g}^b_L P_L + \bar{g}^b_R P_R \right] b \tag{1}$$

where $s_W \equiv \sin\theta_W$, $c_W \equiv \cos\theta_W$ and $P_{L,R}$ are the chiral projection operators. An analogous definition holds for the other fermions. Within the SM, the tree-level values of the chiral couplings $g^f_{L,R}$ are determined by gauge invariance. The weak radiative corrections to the same are well-documented and are insignificant for all but the b-quark. Clearly then,

$$R_b \equiv \frac{\Gamma(Z \to b\bar{b})}{\Gamma(Z \to \text{hadrons})} \simeq \frac{(\bar{g}^b_L)^2 + (\bar{g}^b_R)^2}{\sum_q \left[(\bar{g}^q_L)^2 + (\bar{g}^q_R)^2\right]} \tag{2}$$

where the sum is to be done over all the light quarks. The forward-backward asymmetry at LEP, on the other hand, is given by

$$A^b_{FB}|_{\sqrt{s}\simeq m_Z} = \frac{3}{4} A_\ell \, A_b \tag{3}$$

with

$$A_b \simeq \frac{(\bar{g}^b_L)^2 - (\bar{g}^b_R)^2}{(\bar{g}^b_L)^2 + (\bar{g}^b_R)^2}$$

$$A_\ell \simeq \frac{(g_L^\ell)^2 - (g_R^\ell)^2}{(g_L^\ell)^2 + (g_R^\ell)^2}. \tag{4}$$

Small corrections also accrue to the above observable from a non-zero b-quark and c-quark masses as well as QCD, electroweak and electromagnetic vertex corrections [4–6]. Whereas the observed values are

$$R_b(obs) = 0.21646 \pm 0.00065 , \quad A_{FB}^b(obs) = 0.0990 \pm 0.0017 , \tag{5}$$

the SM expectations for a top quark mass of 174.3 GeV and a Higgs mass close to its present experimental bound, are $R_b(SM) \simeq 0.2157$ and $A_{FB}^b(SM) \simeq 0.1036$. Thus, while the observed value for R_b is consistent with the SM, that for A_{FB}^b shows, as emphasized before, a relatively large deviation from the predicted value. This relatively large discrepancy may be reduced by choosing larger Higgs masses, although only at the cost of worsening the agreement between theory and experiment for other observables, most notably the lepton asymmetries.

It has been noted, for example in Ref. [7], that the overall consistency of the SM with the data improves if we dismiss altogether the measurement of the forward-backward asymmetry. Such an act of exclusion leads to a preference for new physics scenarios that produce a negative shift in the oblique electroweak parameter S [8], an example being provided by supersymmetric theories with light sleptons [7]. We, instead, choose to consider all experimental data on equal footing.

In this article, we investigate a possible way of resolving the disagreement between the hadronic and leptonic asymmetries through the introduction of new quark degrees of freedom at the weak scale thereby inducing non-trivial mixings with the third generation of quarks. In section 2, we examine the experimental status in order to determine the necessary modifications in the couplings of the right- and left-handed bottom quarks. As the required modification in the right-handed sector turns out to be too large to be obtainable via radiative corrections, we investigate, in section 3, the possibility that tree-level mixing of the bottom quark with exotic quarks might be responsible for the observed deviations. All possible assignments for such quarks are examined for their effects on the precision electroweak observables and the two simplest choices identified. The fits to the data for the two cases are presented in sections 4 and 5 respectively. Other phenomenological consequences, including the question of unification, will be investigated in sections 6 and 7. We reserve section 8 for our conclusions.

2 Bottom Quark Couplings Confront Data

Let us assume a purely phenomenological stance and attempt to determine $\bar{g}_{L,R}^b$ from the data. Even in the limit of infinite precision, the ellipse and the hyperbola representing the solution spaces for eqs.(2, 3) intersect at *four* points with the coordinates given by

$$(\bar{g}_L^b, \bar{g}_R^b) \approx (\pm 0.992 g_L^b(SM), \pm 1.26 g_R^b(SM)) , \tag{6}$$

where we indicate on the right the approximate values of the left- and right- handed couplings necessary to fit the bottom-quark production data at the Z-peak[1]. Clearly, no experiment performed at the Z-peak can reduce the degeneracy any further.

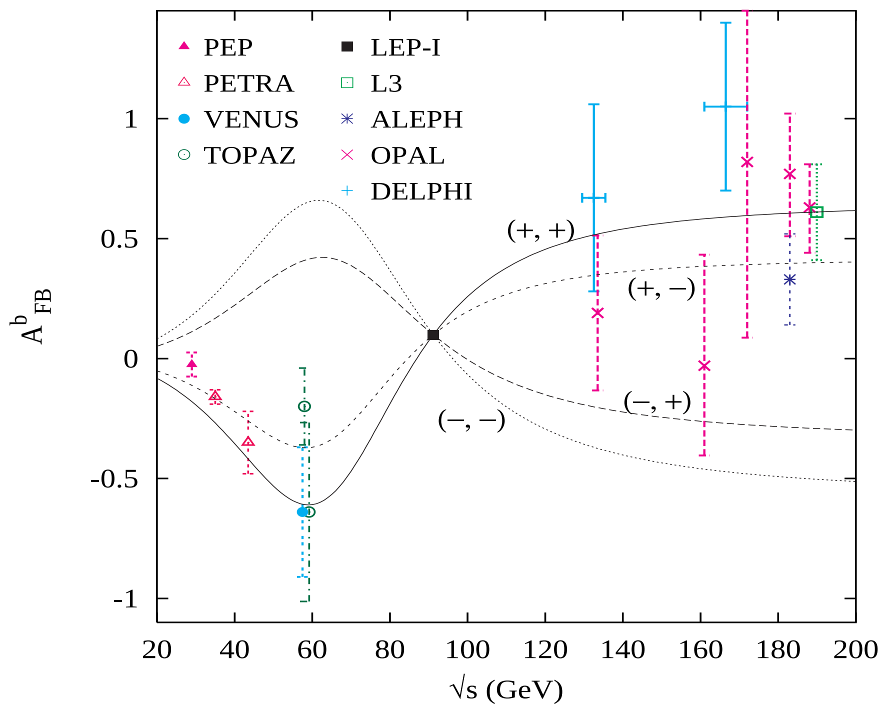

Figure 1: *The forward-backward asymmetry for the b-quark as a function of $\sqrt{s}$ for the four solutions of eq.(6). The signs in the parentheses refer to those for $(\bar{g}^b_L, \bar{g}^b_R)$ in the same order as in eq.(6) with $(+,+)$ being SM-like. The experimental data correspond to the measurements reported in Refs. [10–20].*

Off the Z-peak though, the photon-mediated diagram becomes important thereby affecting the forward-backward asymmetry of the bottom-quark. Such data, thus, could discriminate amongst the four solutions described above. The asymmetry is easy to calculate and in Fig. 1, we plot the same as a function of the center of mass energy of the e^+e^- system for each of the solutions[2] in eq.(6). It is quite apparent that the two solutions with $\bar{g}^b_L \approx -g^b_L(SM)$ can be summarily discarded. Interestingly enough, the data does not readily discriminate between the two remaining solutions. This, though, is not unexpected as $|g^b_R| \ll |g^b_L|$ within the SM. A similar analysis can be performed for R_b as well, but the off-peak measurements of this variable are not accurate enough to permit a similar level of discrimination.

[1] A similar analysis, although restricted to modifying the magnitude but not the sign of the couplings, was performed in Ref. [9]

[2] Had we instead held the magnitudes of the couplings to their SM values, the resulting curves would have been barely distinguishable from those in Fig. 1.

It is quite interesting to note that the agreement with the next best measurement of A^b_{FB}, viz. that at PETRA (35 GeV) is much better for the $(+,-)$ choice than for the SM (or the 'SM-like' solution). This observation can be quantified by performing a χ^2 test including all the data shown in Fig.1. It can easily be ascertained that the χ^2 is indeed significantly improved if the sign of $\bar{g}^b_R$ were to be reversed. Whether this information actually calls for a such a reversal is, of course, open to interpretation.

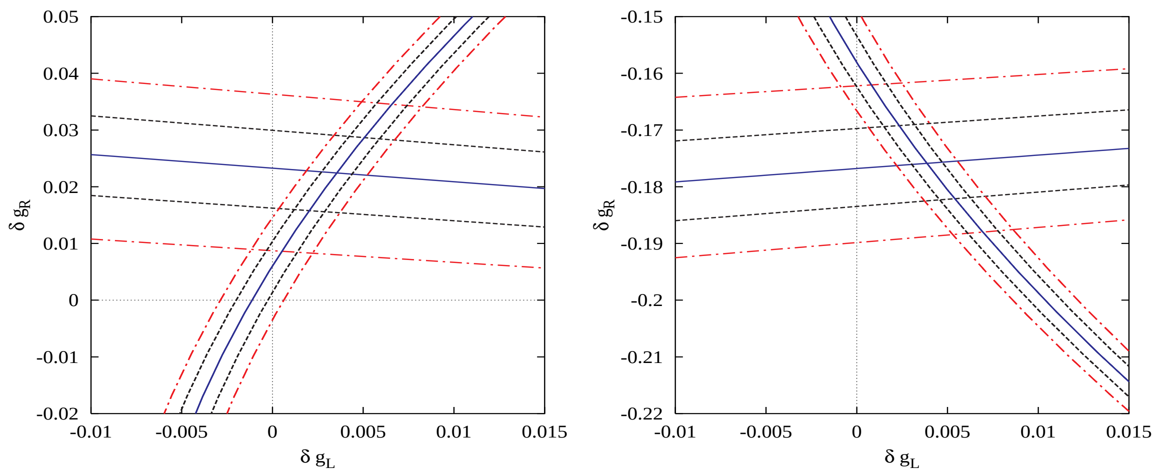

Figure 2: *The regions in the $Z\bar{b}b$ coupling parameter space that are favoured by the observed values of A^b_{FB} (flatter curves) and R_b (steeper curves). For each set, the innermost curve leads to the experimental central value while the sidebands correspond to the 1σ and 2σ error bars. The Standard Model point is at the origin.*

Any resolution of the A^b_{FB} anomaly through a modification of the $Zb\bar{b}$ couplings must then lie within one of two disjoint regions of the parameter space, regions that we exhibit in Fig. 2. What immediately catches the eye is that the required shifts in the coupling satisfy $|\delta g_R| \gg |\delta g_L|$, a condition that would prove crucial at a later stage of our analysis. At this point, it is perhaps worthwhile to note that the two other (ruled out) branches of the solution space would have required a very large $|\delta g_L|$, a shift that is very hard to obtain in any reasonable model.

3 Beautiful Mirrors

We now turn to the question of whether the required $\delta g_{R,L}$ could arise naturally as consequences of ordinary-exotic quark mixing. To keep the discussion simple, yet without losing track of any subtle effects, let us, for now, confine ourselves to just one additional set of quarks. Any extension of the model would not change the qualitative aspects of our analysis. We shall also, for the time being, neglect any mixing with quarks of the first two generations[3]. At this stage, we do not make any further assumptions about the quantum numbers of these new quarks. Working in the basis (b'_1, b'_2), where the primes indicate weak-interaction eigenstates and b'_2 refers to the exotic b-quark, the mass matrix can be parametrized as

$$\mathcal{L}_{m_b} = -\sum_{ij} \bar{b}'_{iL} M_{ij} b'_{jR} + \text{h.c.}, \quad M \equiv \begin{pmatrix} M_{11} & M_{12} \\ M_{21} & M_{22} \end{pmatrix} \tag{7}$$

where the subscripts L, R refer to the quark chirality, and the (in general, complex) elements M_{ij} represent either a bare mass term or one derived from the Higgs mechanism. It is a straightforward task, then, to obtain the mixing matrices for the left- and right-handed quarks (as well as the mass eigenvalues) by diagonalizing the matrices $MM^\dagger$ and $M^\dagger M$ respectively. Ordinary-exotic mixings would generically introduce additional parameters in the charged current structure and, more importantly for us, in the neutral current sector as well. Any such deviation from the SM structure depends crucially on the isospins of the exotics, and for the new left- and right-handed b' fields, we denote these values by $t_{3L(R)}$. Let us concentrate on the neutral currents in the b-sector, more specifically on the part independent of the charge generator[4] Expressed in terms of the physical states (mass eigenstates), these can be parametrized as

$$\begin{aligned} J^3_\mu(b) &= \frac{e}{s_W c_W} \sum_{ij} \bar{b}_i \gamma_\mu (L_{ij} P_L + R_{ij} P_R) b_j \ , \\ L &\equiv \begin{pmatrix} t_{3L} s_L^2 - \frac{1}{2} c_L^2 & -(t_{3L} + \frac{1}{2}) s_L c_L \\ -(t_{3L} + \frac{1}{2}) s_L c_L & t_{3L} c_L^2 - \frac{1}{2} s_L^2 \end{pmatrix} \\ R &\equiv \begin{pmatrix} t_{3R} s_R^2 & -t_{3R} s_R c_R \\ -t_{3R} s_R c_R & t_{3R} c_R^2 \end{pmatrix} \end{aligned} \tag{8}$$

where $s_{L,R} \equiv \sin\theta_{L,R}$ etc parametrize the left- and right-handed mixing matrices in the b-sector and s_W (c_W) denote the sine (cosine) of the weak mixing angle. As expected, we would have flavour-changing neutral currents if either $t_{3L} \neq -1/2$ or $t_{3R} \neq 0$. The presence of any such coupling would play a crucial role in the discovery of such an exotic and we shall return to this later. Since the shifts in $g^b_{L,R}$ are given by

$$\delta g^b_L = \left(t_{3L} + \frac{1}{2} \right) s_L^2 \ , \qquad \delta g^b_R = t_{3R} s_R^2 \ , \tag{9}$$

it follows that the right handed component of the exotic cannot be a $SU(2)_L$ singlet.

In principle, we could allow each of b'_L and b'_R to lie in any (and inequivalent) representation of $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$. However, the requirement of anomaly cancellation

[3] A negligibly small mixing with the first and second generations may be enforced, for instance, by additional gauge interactions, such as Top-color [21–25], Top-flavor [26–29] or Bottom-color [30].

[4] Any mixing which respects $U(1)_{\text{em}}$ must be between objects of the same charge Q and thus there are no mixing effects proportional to Q.

indicates that a vector-like coupling for the exotics is the most economic choice. In addition, the introduction of vector-like fermions, unlike the one of their chiral counterparts, do not lead to large contributions to the oblique electroweak parameter S [8], thereby preserving the agreement with precision data. We shall, therefore, limit ourselves to only such quarks. We refer to these exotic quarks as (beautiful) mirrors in the sense that they occur in vector-like pairs and they have the same electric and color charges as the ordinary bottom- (or beauty-) quark.

As we have already mentioned, nonzero quark mixing requires that there be mass terms connecting the ordinary b-quark to its exotic counterpart. Demanding that the only scalars in the theory be the $SU(2)$ Higgs boson doublets restricts the choice of the exotics to a $SU(2)$ singlet and two varieties each of $SU(2)$ doublets and triplets. The phenomenological requirement of $t_{3R} \neq 0$ eliminates the singlet and one of the triplets as the possible source for the large modification of the right-handed bottom quark coupling. The choice then devolves to one of $\Psi_{L,R} = (3,2,1/6), (3,2,-5/6)$ and $(3,3,2/3)$. The phenomenological consequences of the first and last possibilities are similar and hence we shall concentrate on studying the first two cases.

4 Scenario with Standard Mirror Quark Doublets

Our first scenario relies on the introduction of the fermion doublets

$$\Psi^T_{L,R} = (\chi, \omega) \equiv (3,2,1/6), \tag{10}$$

which are a mirror copy of the standard quark doublets of the Standard Model. The most general Yukawa and mass term in the Lagrangian is then

$$\mathcal{L} \supset -\left(y_1 \overline{Q'_L} + y_2 \overline{\Psi_L}\right) b'_R \phi - \left(x_1 \overline{Q'_L} + x_2 \overline{\Psi'_L}\right) t'_R \tilde{\phi} - M_1 \overline{\Psi'_L} \Psi'_R + h.c., \tag{11}$$

where the primes, once again, denote weak eigenstates. Note that, on account of Ψ'_L and Q'_L having the same quantum numbers, a mass term of the form $\overline{Q'_L} \Psi_R$ can be trivially rotated away. In the basis (b', ω'), we then have a mass matrix of the form

$$M_b = \begin{pmatrix} Y_1 & 0 \\ Y_2 & M_1 \end{pmatrix} , \quad Y_i \equiv y_i \langle \phi \rangle \tag{12}$$

and an analogous one for the top. For the sake of simplicity, we shall assume that the mass matrices are real. In the phenomenologically interesting regime of $Y_1 \ll Y_2 < M_1$, we have, for the eigenvalues of the mass eigenstates b, ω and the mixing angles

$$\begin{aligned} m_b &\approx Y_1 \left(1 + \frac{Y_2^2}{M_1^2}\right)^{-1/2} , & m_\omega &\approx (M_1^2 + Y_2^2)^{1/2} , \\ \tan\theta^b_R &\approx \frac{-Y_2}{M_1} & \tan\theta^b_L &\approx \frac{-Y_1 Y_2}{M_1^2 + Y_2^2} , \end{aligned} \tag{13}$$

with analogous expressions for the top-sector. A few points are to be noted:

- Since both ω'_L and χ'_L have the same quantum numbers as their ordinary counterparts, neutral currents in these sectors remain unmodified.

- As $\delta g^b_R < 0$ (while $g^b_R(SM) > 0$), a small δg^b_R would only worsen the fit. Rather, we must demand a large negative correction that would take us to the second allowed region in the parameter space (see Fig.2). For example, a 1σ agreement for each of A^b_{FB} and R_b is obtained for

$$\delta g^b_R = \frac{-s^2_R}{2} \approx -0.165 \quad \Longrightarrow \quad Y_2 \approx 0.7\, M_1. \tag{14}$$

- The top-sector mass matrix is as in eq.(12) but with $y_i \to x_i$. Since the y's and x's are independent, one could, in principle, set $x_2 = 0$ (this, for example, could be ensured by imposing a discrete symmetry). In such a case, the top sector sees no additional mixing and x_1 is the usual top Yukawa coupling.

- Since no exotic quark has yet been seen at the Tevatron collider, $M_1 \gtrsim 200$ GeV.

- In general, due to the large mixing in the bottom sector and the fact that the right-handed mirror quarks carry non-trivial weak charges, a potentially large correction to the precision electroweak parameters will accrue.

- A right-handed W-t-b interaction is induced with strength proportional to $s^b_R s^t_R$. Measurement of $b \to s\gamma$ requires $s^b_R s^t_R < 0.02$ [31], leading us to consider the case of negligible χ-t mixing.

In order to address the question of how well does this scenario fit the data, we have computed the corrections to the S, T and U parameters, with respect to a reference Higgs mass value of 115 GeV. While the corrections to U are small, the corrections to T and S are large and increase with the overall scale of quark masses. For instance, for $M = 200$, 225, 250 GeV the corrections to the T parameter are $\Delta T \simeq 0.35$, 0.42, 0.54 respectively, while the correction to the parameter S is somewhat insensitive to the masses and measures $\Delta S \simeq 0.1$.

The large corrections to the T parameter, together with the relatively large corrections to the right-handed bottom couplings, tend to increase the hadronic width and the total width of the Z to unacceptable levels. This problem can be ameliorated by including the mixing of the bottom quark with a quark $\xi_{R,L}$ carrying the quantum numbers of the right-handed bottom quark and its mirror partner. In the basis (b', ω', ξ'), the simplest modification to the mass matrix that fulfills this requirement is given by:

$$M_b = \begin{pmatrix} Y_1 & 0 & Y_3 \\ Y_2 & M_1 & 0 \\ 0 & 0 & M_2 \end{pmatrix} \quad , \quad Y_i \equiv y_i \langle \phi \rangle \tag{15}$$

Note that, as happens with $(M_b)_{12}$, the element $(M_b)_{31}$ could also be trivially rotated away. The inclusion of small, but non-zero, values of the elements $(M_b)_{23}$ and $(M_b)_{32}$ only serves to complicate matters without modifying the main phenomenological consequences of this model.

Ignoring small terms proportional to the bottom quark mass, the left-handed mixing angle is now given by

$$s_L \simeq \frac{Y_3}{\sqrt{Y_3^2 + M_2^2}} \tag{16}$$

The main effect of the mixing with these weak singlet quarks is to reduce the left-handed coupling of the bottom quark and thus the partial width of the Z into b's and hence into hadrons as such. The scenario described above can thus clearly improve the agreement with A_{FB}^b. For small values of s_L, as demanded by experimental results, the oblique corrections to the precision electroweak observables are still dominated by the large mixing of the bottom quark with the weak mirror quark doublet. The presence of the new quarks will, of course, induce additional radiative corrections to the b-quark couplings. It is easy to see though that, given the above-mentioned mass and mixing angle pattern, these corrections are tiny compared to those induced at the tree level, and hence could be safely neglected. Once again, non-observation of an exotic quark at the Tevatron implies $M_2 \gtrsim 200$ GeV.

The parameters S, T and U are in one to one correspondence with the variations of the parameters ϵ_3, ϵ_1 and ϵ_2, introduced in Ref. [32] with respect to a given value of the top quark mass and the Higgs mass. The relation between these parameters is given by:

$$\begin{aligned} \Delta\epsilon_1 &= \alpha T, \qquad \Delta\epsilon_3 = \frac{\alpha S}{4 s_W^2}, \\ \Delta\epsilon_2 &= -\frac{\alpha U}{4 s_W^2}, \end{aligned} \tag{17}$$

and, within the SM, for a top quark mass $m_t = 174.3$ GeV and a Higgs mass $m_H = 115$ GeV, $\epsilon_1 = 5.6\ 10^{-3}$, $\epsilon_2 = -7.4\ 10^{-3}$ and $\epsilon_3 = 5.4\ 10^{-3}$ [7]. The dependence of the ϵ_i parameters on the Higgs and top quark masses may be found, for instance, in Ref. [33]. The dependence of the most important observables on the parameters ϵ_i are given by

$$\begin{aligned} \Gamma_Z &\simeq 2.489\ (1 + 1.35\ \epsilon_1 - 0.46\ \epsilon_3 + ...)\ \text{GeV} \\ \sin^2\theta_l^{\text{eff}} &\simeq 0.2310\ (1 - 1.88\ \epsilon_3 + 1.45\ \epsilon_1) \\ \frac{m_W^2}{m_Z^2} &\simeq 0.7689\ (1 + 1.43\ \epsilon_1 - \epsilon_2 - 0.86\ \epsilon_3)\ , \end{aligned} \tag{18}$$

where the dots reflect the contributions associated with the variations of the b-coupling due to radiative corrections and the mixing with the b quarks. There is, in addition, a dependence on the precise value of $\alpha(M_Z)$ and $\alpha_s(M_Z)$. In our computations we have used the central values for the hadronic contribution to $\alpha(M_Z)$ namely $\Delta\alpha_{\text{had}}^{(5)} = 0.02761$ [34], while the strong gauge coupling was allowed to float around the central value of 0.118 [35].

Variations in the parameter T larger than 0.3 tend to induce a large positive correction to the total Z width, and are therefore disfavored by the data. Consequently, this model leads to a better fit to the data whenever the new quarks are relatively light and the Higgs is heavy. In general, a heavy Higgs leads to a negative contribution to T and a positive contribution to S, leading to a better agreement with the total width of the Z. However, the same effects worsen the agreement of the data with the leptonic asymmetries, which mainly depend on $\sin^2\theta_l^{\rm eff}$. Therefore, the model leads to a correlation of the quark and Higgs masses: The heavier the new quarks, the heavier the Higgs needs to be. The value of α_s also plays an important role in this process, since it may lower the total hadronic width without modifying the leptonic asymmetries.

We have made a fit to the data within this model. Including the values of Y_2, M_1, α_s, m_t, m_H and s_L as variables in the fit (the fit is quite insensitive to the scale M_2, provided it remains below a a few TeV), we obtain that the best fit to the data is obtained for the mirror quark mass parameter M_1 close to the present experimental bound on this quantity, while the preferred values of the Higgs mass are about 300 GeV. Raising the quark bound to 250 GeV leads to an optimal value of the Higgs mass of about 850 GeV. The best fit gives $\alpha_s \simeq 0.116$ and $m_t \simeq 173$ GeV. For the exotic sector, the corresponding values are

$$Y_2 \simeq 0.71\ M_1, \qquad M_1 \simeq 200\,\mathrm{GeV} \tag{19}$$

while

$$s_L^2 \simeq 0.008. \tag{20}$$

The best fit to the ratio Y_2/M_1 and to s_L^2 are virtually independent of M_1 for 200 GeV $\lesssim M_1 \lesssim$ 250 GeV.

In Fig. 3 we show the 1- and 2-σ regions in the m_H–m_χ parameter space determined by the best fit to the data. As emphasized above, the model leads to a preference for light quarks, with masses below or about 250 GeV, within the reach of the Tevatron collider (see section 6), while the preferred values of the Higgs mass are much larger than in the Standard Model, a feature that appears in many models [36].

For the parameters providing the best fit, all measured precision electroweak observables, including the lepton and hadron asymmetries and the Z widths are within 2σ of the predictions of this model, and, in particular, the bottom asymmetries are within 1σ of the predicted values. Similar results are obtained for slightly larger values of M_1, although the model is clearly disfavored for quark masses $M_1 > 250$ GeV. While the agreement between theory and experiment for the bottom-quark asymmetries is remarkably better than in the Standard Model, the lepton asymmetries remain essentially the same as in the SM. This is due to the tension between these observables and the total Z-width within this model and is reflected in the fact that the fit produces

$$\sin^2\theta_l^{\rm eff} \simeq 0.2315\ . \tag{21}$$

As a result, the left-right lepton asymmetry is about 1.9 standard deviations away from the value measured at SLD thus marginally worsening the discrepancy obtained within the Standard Model. The W mass is, instead, in excellent agreement with the predictions of this model.

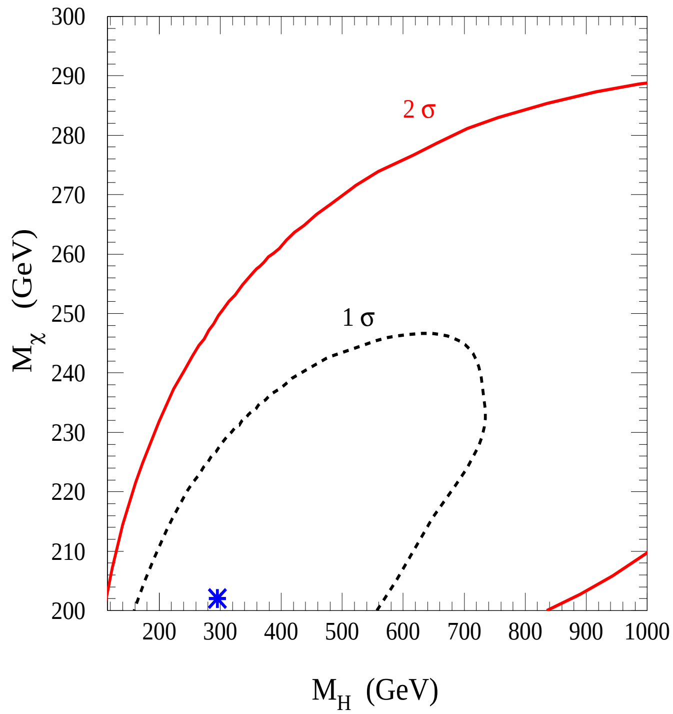

Figure 3: *Region in the m_H–m_χ parameter space (in the model with standard mirror quark doublets) that is consistent with the best fit point (marked) at the 68% C.L. and 99.5% C.L. respectively.*

5 Top-less Mirror Quark Doublets

Let us now analyze the case in which the mirror quarks belong to a doublet in which there is no quark with the same charge as the top quark, viz.,

$$\Psi^T_{L,R} = (\omega, \chi) \equiv (3, 2, -5/6). \tag{22}$$

This model has some advantages with respect to the model analyzed above. First of all, since the weak partner of the ω has charge $-4/3$, there is no mixing involving the top quark. Second, the model allows for a modification of the right-handed bottom couplings with moderate mixing angles. In the basis (b', ω'), the mass matrix reads

$$M_b = \begin{pmatrix} Y_1 & 0 \\ Y_R & M_1 \end{pmatrix} \ , \quad Y_i \equiv y_i \langle \phi \rangle \ , \tag{23}$$

where the zero entry is now enforced by gauge invariance. The right and left-handed mixing angles have similar expressions to the ones found in the above model. However, the mixing of the right-handed bottom leads to a positive shift of g^b_R,

$$\delta g^b_R \simeq \frac{1}{2} s^2_R \tag{24}$$

and therefore small values of s_R can lead to a relevant shift of A_{FB}^b in the direction required by experiment.

As in the previously analyzed model, the values of R_b and of the hadronic width can be improved by allowing an additional mixing with a quark ξ with the same quantum numbers as the right-handed bottom quark and its mirror partner. In the basis (b', ω', ξ'), we assume a mass matrix quite similar to the one in the last section:

$$M_b = \begin{pmatrix} Y_1 & 0 & Y_L \\ Y_R & M_1 & 0 \\ 0 & 0 & M_2 \end{pmatrix} \quad , \quad Y_i \equiv y_i \langle \phi \rangle \tag{25}$$

As with the previous model, the matrix element $(M_b)_{31}$ can be trivially rotated away, while the inclusion of small (compared to M_i) but non-vanishing matrix elements $(M_b)_{23}$ and $(M_b)_{32}$ do not change the main results of our analysis.

Ignoring small effects induced by the bottom mass, the left-handed and right-handed mixing angles are given by

$$s_L \simeq \frac{Y_L}{\sqrt{Y_L^2 + M_2^2}}, \qquad s_R \simeq \frac{Y_R}{\sqrt{Y_R^2 + M_1^2}}. \tag{26}$$

The main difference between this model and the one analyzed before lies in the smallness of all the mixing angles. Since all Yukawa couplings are small compared to the explicitly gauge invariant masses, the corrections to the oblique parameters S, T and U are small. The corrections to T become relevant only for quark masses above 500 GeV, while the corrections to S and U remain small even for masses in the multi-TeV range. Since the expected bottom-quark asymmetry has now migrated much closer to the measured value, the data now prefers non-negligible values of the T parameter so as to permit a better agreement with the lepton asymmetries and the W mass. This can only be achieved by pushing the quark masses up, while keeping the Higgs mass close to its experimental lower bound.

We have analyzed all the precision observables in the context of this model, as described by the parameters m_H, m_t, α_s, M_1, Y_R, M_2 and Y_L. The best fit to the data is obtained for a Higgs mass close to the present experimental bound and mirror quark doublets with mass of about

$$M_1 \simeq 825\,\mathrm{GeV} \quad \text{and} \quad Y_R \simeq 160\,\mathrm{GeV}, \tag{27}$$

implying that $s_R^2 \simeq 0.036$. The best fit value of M_2 is close to its experimental bound, $M_2 \simeq 200$ GeV, while $Y_L \simeq 15$ GeV, leading to $s_L^2 \simeq 0.006$. Similar to the previously analyzed scenario, changing M_2 while keeping the ratio of Y_L/M_2 does not alter the fit in any significant way. The best fit values of α_s and m_t are $\alpha_s \simeq 0.116$ and $m_t \simeq 176$ GeV.

In Fig. 4 we show the 1- and 2-σ regions in the m_H–M_1 parameter space obtained by the best fit to the data. As emphasized above, the Higgs tends to be light, in the region most accessible to the Tevatron collider and a $\sqrt{s} = 500$ GeV linear collider. The quarks, instead, tend to be heavy with masses of about 1 TeV. For M_1 above a few TeV though,

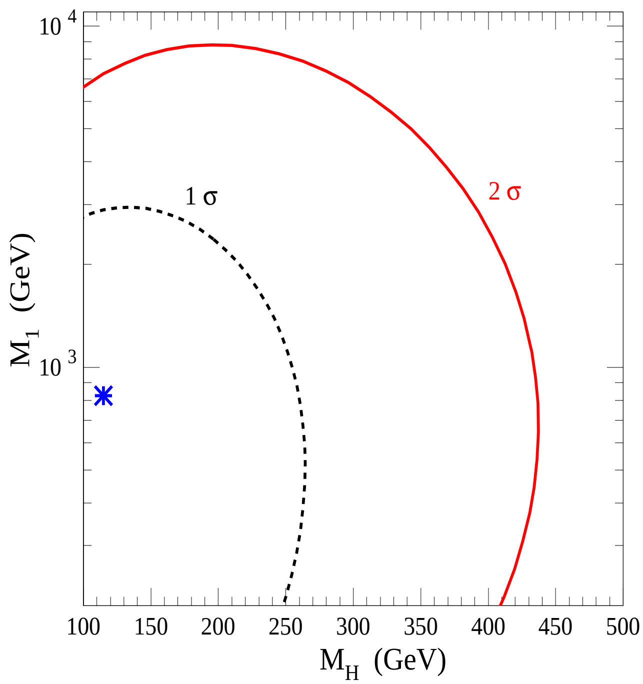

Figure 4: *Region in the m_H–m_χ parameter space (in the model with top-less mirror quark doublets) that is consistent with the best fit point (marked) at the 68% C.L. and 99.5% C.L. respectively.*

the Yukawa coupling Y_R needed to improve the fit to the data becomes large, spoiling the perturbative consistency of the theory at low energy scales.

This model provides a surprisingly good agreement with the experimental data. For the parameters providing the best fit to the data, the left-right lepton asymmetry measured at SLD is 1.2 standard deviations from the theoretically predicted value, while almost all other measured observables are within 1 σ of the predictions of this model. The only exceptions are the charm forward-backward asymmetry and the total hadronic cross section measured at LEP, which stay within 2 σ of the measured values. As already pointed out, the fitted value

$$\sin^2 \theta_l^{\rm eff} \simeq 0.2313 \tag{28}$$

exhibits a much better agreement with the leptonic asymmetries than in the model with Standard Mirror Quarks.

6 Implications at Present and Future Colliders

Although the two models presented above share many features, there are subtle differences as far as the collider signatures are concerned. We shall examine, in some detail, the scenario containing a top-like quark and then point out the differences with the second model.

Any new quark [37], with a mass below or about 250 GeV, as preferred in the model with standard quark doublets, should be observable at the next run of the Tevatron collider. The main decay of the χ-quark will be similar to that of the top quark, namely,

$$\chi \to b + W^+. \tag{29}$$

Although the ω-quark tends to be somewhat heavier than the χ-quark, yet it should be possible to pair-produce it at the Tevatron collider, especially if M_1 turns out to be in the (lower) range preferred by the electroweak fit. On account of the phase space restrictions, it would decay mainly through the flavor violating channel, *viz.*,

$$\omega \to b + Z \tag{30}$$

a mode that has already been looked for at the Run I of the Tevatron with a resultant lower limit of about 200 GeV on m_ω [38]. This very same flavor-changing neutral current interaction, would, at a next generation linear collider, lead to a rather significant cross section for the process

$$e^+ + e^- \to \bar{b} + \omega, \tag{31}$$

as well as the conjugate state, thereby leading to rather striking signatures.

There is no prediction for the singlet quark mass within this models. If they were light, in the range accessible to the Tevatron collider, they will mainly decay via its flavor violating couplings

$$\xi \to b + Z\ , \tag{32}$$

or, in the event of a non-zero $(M_b)_{23}$, also into an ω-quark:

$$\xi \to \omega + Z. \tag{33}$$

Due to its larger center of mass energy, the LHC, should be able to test this model for even larger values of the ω- and χ- quark masses.

Finally, we turn to the Higgs. Note that the Yukawa coupling matrix is proportional to that in eq.(15) with the gauge invariant masses $M_{1,2}$ switched off. This immediately implies that the Yukawa interactions are not diagonal in the mass-basis. This is quite crucial especially since the coupling y_2 is quite large (actually, of the same order as the top Yukawa in the SM). While the full expressions for the resultant Yukawas are cumbersome, they simplify considerably in the limit of a vanishing y_3 to

$$\mathcal{L}_{bw\phi^0} \sim y_2 \left(c_R\, \bar{\omega}_L b_R - s_R\, \bar{\omega}_L \omega_R\right) \phi^0 + h.c. + \mathcal{O}(y_1) \tag{34}$$

Thus, for $m_{\phi^0} > m_\omega + m_b$, a condition satisfied over almost the entirety of the preferred parameter space (see Fig.3), the Higgs is afforded an additional decay mode. And, if the Higgs is more than twice as heavy as the ω (again, true for a very large part of the 1σ preferred area), a further decay channel would open up, although with a branching fraction smaller than that for the flavour-changing mode. While this results in a severe depletion of the 'gold-plated' modes ($\phi^0 \to ZZ \to 4l$), presumably these non-canonical channels would lend themselves easily to discovery, especially with the b- (and lepton-) richness of the final state.

In the top-less mirror quark model, instead, the quark doublets are predicted to be heavier. Only the LHC (or a second generation linear collider) would have enough center of mass energy to produce the ω and χ in this case. The quark χ (with a $-4/3$ charge) will decay into

$$\chi \to b + W^- \ , \tag{35}$$

leading to a top-like signature with a wrong sign W. Similarly, due to phase space restrictions,

$$\omega \to Z + b \tag{36}$$

Since the flavour changing coupling is smaller in this case as compared to the previous model, non-diagonal production at a linear collider would be somewhat suppressed. Also, with the Higgs preferring to be light, and with its couplings to the new states not being as large as in the other model, Higgs phenomenology remains largely unchanged from the SM.

As happens in the model with standard mirror quarks, there is no prediction for the singlet mirror quark masses. If they are available at the colliders, they will decay via its flavor violating coupling:

$$\xi \to b + Z \ . \tag{37}$$

And since ξ prefers to be lighter than ω, a non-zero $(M_b)_{23}$ would, once again open an additional decay channel, viz.,

$$\omega \to \xi + Z \ . \tag{38}$$

7 Unification

The question of unification within these models is an interesting one. It might be argued that this is not a pertinent one, since we have not detailed any mechanism to keep the Higgs boson and vector quarks naturally light. The hierarchy of masses may arise through some hitherto undiscovered mechanism, which must be related to the one leading to the breakdown of the electroweak symmetry, and will probably require extra gauge symmetries [21–30, 39–41]. The presence of any such extra symmetry would certainly alter unification in a highly model-dependent way. It is still possible though to discuss the issue of unification without delving into the details of the implementation of such a symmetry, as long as one assumes that the low-energy spectrum is determined, besides additional complete SU(5) multiplets [41]. This is the approach that we adopt in this section.

We shall proceed with a one-loop analysis, taking into account that the possible two-loop effects are of the order of the small threshold effects at the Grand Unification scale, and hence become strictly relevant only for the construction of a complete grand unified model, an objective which is beyond the scope of this article.

In the model with standard mirror quarks the beta-function coefficients of the three gauge couplings are given by

$$
\begin{aligned}
b_3 &= -11 + \frac{4}{3}n_g + 2 \\
b_2 &= -\frac{22}{3} + \frac{4}{3}n_g + \frac{n_H}{6} + 2 \\
b_1 &= \frac{4}{3}n_g + \frac{n_H}{5} + \frac{2}{5}
\end{aligned}
\tag{39}
$$

where the last term in each line relates to the extra contribution induced by the presence of the mirror doublet and singlet bottom quarks, n_g is the number of generations and n_H is the number of Higgs doublets. To be consistent with the electroweak fits described earlier, we assume that only one Higgs doublet plays a relevant role in electroweak symmetry breaking.

This model predicts a shift of the hypercharge beta-function that is somewhat smaller than the equal shifts of the beta-functions of the weak and strong gauge couplings. As is well-known, within the SM, the strong and weak gauge couplings meet at approximately 10^{17} GeV (for $n_H = 1$). The hypercharge coupling, however, crosses the other two at a much lower scale. This big 'discrepancy' can be reduced by postulating a large number of Higgs doublets, but only at the cost of bringing down the unification scale to $\sim 10^{12}$ GeV, a value palpably inconsistent with proton stability. The introduction of the new quarks and the consequent shift in the beta-functions works to reduce this very same difference in the scales at which the couplings meet. The improvement is quite significant. For $n_H = 1, 2$ and 3, and 'unification' scales of approximately 5×10^{16} GeV, 2×10^{16} GeV and 10^{16} GeV, the couplings differ from the average "unification" value by less than 3, 1.5 and 1 percent respectively[5]. These corrections are small, and considering the large scales involved, as emphasized above, could easily be accommodated (even for $n_H = 1$) by threshold effects due to the presence of heavy particles (with GUT scale masses) or Planck scale suppressed operators.

Observe that, since the model lacks supersymmetry, dangerous dimension five operators are absent from a potential grand unified scenario. Moreover, the unification scale is sufficiently large to avoid the constraints coming from proton decay induced via dimension 6 operators. However, for the Higgs masses and Yukawa couplings associated with the best fit to the precision electroweak data, the model tends to induce a Landau pole in the Higgs quartic couplings below the GUT scale, particularly for relatively heavy quarks, $M_1 \simeq 250$ GeV. The most obvious means of avoiding this problem is to give up the property of perturbative unification. An alternate way would be to effect a suitable modification such as the one that we discuss shortly.

In the model with the non-standard mirror doublet quarks, instead, the unification relations are not improved with respect to the Standard Model case, since the shift in the

[5] Two-loop effects will produce small modifications to these numbers.

beta function of the hypercharge gauge coupling is larger than in the ones associated with the strong and weak coupling beta-functions. However, this model presents an interesting property: the mirror doublet and singlet quarks introduced in this model are contained in the adjoint (**24**), and in the $\mathbf{5} + \bar{\mathbf{5}}$ representations of $SU(5)$. Indeed, the weak doublet and singlet quarks in this model have precisely the same quantum numbers as the 24-plet partners of the standard model gauginos and of the color-triplet Higgsinos of the minimal supersymmetric standard model (MSSM), respectively[6].

On a more speculative note, if one were willing to accept the presence of a complete **24** of fermions at the weak scale, together with the standard mirror doublet and singlet quarks[7], one could obtain excellent unification relations without the need of supersymmetry, together with an excellent fit to the precision electroweak observables. Within this assumption, the top-less mirror quark doublets will be the ones leading to a relevant mixing with the bottom quark, while the standard mirror quark doublets should have only small mixing with the three generation of quarks. In this case, the Higgs tends to be light. For a top-less doublet lighter than approximately 1 TeV, the Yukawa couplings are relatively weak and the Landau pole problem is avoided. This model may lead, instead, to a conflict with the Higgs potential stability [42]. Whether this is a real physical problem can only be answered by studying the possibility of ours being a metastable vacuum. In the SM with a similarly light Higgs, $m_H \simeq 115$ GeV, the requirement of strict stability would suggest the presence of new physics far below the Planck scale. The requirement of being in a metastable vacuum with lifetime longer than the age of the Universe, instead, allows the Standard Model description to be valid up to scales close to the Planck scale [43, 44]. We postpone for a future study a more detailed analysis of these questions.

Observe that the above-mentioned possibility leads to the potential presence of fields with the quantum numbers of the MSSM gauginos at low energies. In the absence of a symmetry like R-Parity, the fields with the quantum number of the Wino and of the Bino will mix with the leptons and neutrinos, respectively, and hence their couplings to the leptons and the Higgs bosons should be very small. If one assumes the presence of a second Higgs doublet, with no relevant role in the electroweak symmetry breaking mechanism, a coupling of order one of this field with the Bino-like field and, for instance, the third generation leptons may induce the proper annihilation rate to make the Bino-like field a good dark matter candidate[8]. Alternatively, the Bino may play the role of a sterile neutrino. Without the addition of new fields, the gluino-like particle tends to be very long lived or even stable. We also reserve for a separate study the analysis of the cosmological and phenomenological consequences of such a scenario.

[6] In the minimal supersymmetric $SU(5)$ model, the colored-triplet Higgsinos are assumed to acquire GUT scale masses , creating the so-called doublet triplet splitting problem. A similar hierarchy problem would exist in the model described in this section.

[7] All these fields are contained in the adjoint of E_6.

[8] Discrete symmetries may need to be imposed in order to avoid dangerous lepton flavor violating processes.

8 Conclusions

The Standard Model with a light Higgs boson is in very good agreement with the precision electroweak observables measured at the Tevatron, SLD and LEP colliders. Although there is no clear indication of the need for new physics in the electroweak precision measurement data, the prediction for the effective leptonic weak mixing angle extracted from the hadronic and leptonic observables are several standard deviations away from each other. In this article we have analyzed a possible way of fixing this discrepancy by introducing mirror quarks with quantum numbers similar to those of the left-handed and right-handed bottom quarks.

While the two models analyzed in our article lead to an improvement of the general fit to the precision electroweak data, they present qualitatively different characteristic that make themselves easily distinguishable from the experimental point of view. In the model with standard mirror quark doublets, only negative shifts to the right-handed bottom coupling may be obtained by means of the mixing with the doublet and singlet quarks. The very smallness of this coupling within the SM, however, allows us not only to change its magnitude but reverse its sign as well. Apart from improving the agreement to the precision electroweak data to a great extent, this also leads to interesting predictions for the bottom quark asymmetries away from the Z peak. Moreover, the best fit to the data within this model is obtained for quarks light enough to be accessible at the Tevatron collider, and relatively heavy Higgs bosons. Finally, the unification relations are significantly improved with respect to the Standard Model case, and the potential unification scale is sufficiently large in order to avoid proton decay via dimension six operators. However, perturbative unification within this simple extension of the Standard Model is not possible, due to the presence of a Landau pole in the Higgs quartic couplings at scales below the potential GUT scale.

On the other hand, the model with non-standard mirror quark doublets leads to mild modifications of the left- and right-handed bottom quark couplings induced via small mixings of the mirror quarks with the standard ones. Besides this, the model prefers relatively light Higgs bosons, possibly in the range testable at the next run of the Tevatron collider, while the mirror quarks tend to be heavy, only accessible at the LHC. An interesting property of this model is that the quantum numbers of the exotic quarks are precisely the ones of the heavy coloured gauginos and Higgsinos within minimal supersymmetric SU(5) scenarios.

One can contemplate the possibility of taking both standard and exotic mirror quark doublets and the mirror down quark singlets, and including particles with the quantum numbers of the standard gauginos in the minimal supersymmetric Standard Model and, eventually, an additional Higgs doublet. This extension of the Standard Model allows a remarkable improvement in the fit to the precision electroweak observable data, leads to the possibility of achieving consistency with the unification of gauge couplings and has all the ingredients necessary to lead to an explanation of the dark matter content of the Universe.

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