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More about soft terms and FCNC in realistic string constructions

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Abstract: In realistic four-dimensional string constructions the presence of anomalous U(1)'s is generic. In addition, the associated Fayet-Iliopoulos contribution to the D-term can break the extra gauge symmetries. As a consequence, physical particles can appear combined with other states. We show that even if a three-generation standard-like model has originally flavour-independent soft scalar masses, the particle mixing contribution may generate non-universality among them. Thus FCNC effects which were apparently absent reappear. We also discuss the size of these contributions in an explicit model, and how they can be suppressed.

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Since the late eighties a number of interesting four-dimensional string vacua with particle content not far from that of the supersymmetric (SUSY) standard model have been found (see e.g. the discussion in [1] for heterotic string models and [2] for intersecting brane models, and references therein). Let us recall that the gauge groups obtained after compactification are generically larger than the standard model gauge group, containing extra U(1) symmetries, $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)^n$ [3]. Actually, at least one of these U(1)'s is usually anomalous[1]. For example, it was found in [5] that only 192 different three-generation models containing the $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)^n$ gauge group can be constructed within the Z_3 heterotic orbifold with two Wilson lines. The matter content of 175 of them was analyzed in detail and only 7 of them turn out not to have an anomalous U(1) associated. Although the anomaly is cancelled by the four-dimensional Green-Schwarz (GS) mechanism, it generates a Fayet-Iliopoulos (FI) contribution to the D-term [6]. This effect is crucial for model building [7] since some scalars C's, in particular $\mathrm{SU}(3) \times \mathrm{SU}(2)$ singlets, acquire large vacuum expectation values (VEVs) in order to cancel the FI contribution, breaking the extra gauge symmetries, and allowing the construction of realistic standard-like models [8 − 10][2]. In particular, many particles, which we will refer to as ξ, acquire a high mass because of the generation of effective mass terms. These come for example from operators of the type $C\xi\xi$. In this way vector-like triplets and doublets and also singlets become heavy and disappear from the low-energy spectrum. The remarkable point is that the standard model matter remain massless, surviving through certain combinations with other states.

The FI breaking also has an important implication in the flavour structure, that is, derivation of realistic quark/lepton mass matrices in string models. For example, stringy 3-point couplings in heterotic Z_3 orbifold models do not lead to completely realistic quark/lepton mass matrices before the FI breaking[3]. However, after the FI breaking, the particle mixing appears through Yukawa couplings including scalars C's with non-vanishing VEVs, and that can induce non-trivial quark/lepton mass matrices for light modes [14]. Detailed study on realizing quark/lepton mass matrices has been carried out in [15], showing the possibility for deriving realistic quark/lepton mass matrices in string models.

Generally speaking, a flavour structure leading to realistic fermion mass matrices may affect somehow their superpartner sector, that is, soft SUSY breaking sfermion masses and the so-called A-terms. Thus, it is quite important to study sfermion masses and A-terms in each flavour structure, of course including the above-mentioned particle mixing scenario through the FI breaking. Indeed, there are strong experimental constraints on flavour changing neutral currents (FCNCs). That requires degeneracy at least between the

[1]In [4], some conditions for the absence of the anomalous U(1) in heterotic string models are discussed, and classifications of models with anomalous U(1) are also attempted.

[2]In brane models where several anomalous $\mathrm{U}(1)'s$ are usually present the FI terms do not necessarily trigger further gauge symmetry breaking since they may vanish at the orbifold singularity. However, the possibility of using non-vanishing FI terms for model building is also present and has been discussed in the literature (see e.g. [11]).

[3]Stringy 3-point couplings in non-prime order orbifold models have possibilities for realizing non-trivial quark/lepton mass matrices [12, 13]. Higher dimensional operators may also be important to derive realistic quark/lepton mass matrices.

first and the second generations of sfermion masses, unless they are sufficiently large like $O(10)$ TeV. For example, the FCNC constraints on the Kaon system require

$$Re\left((m^d_{12})^2_{LL,RR}/m^2_{ave}\right) \leq 4\times 10^{-2}\frac{m_{ave}(GeV)}{500}, \quad (1)$$

$$Im\left((m^d_{12})^2_{LL,RR}/m^2_{ave}\right) \leq 3\times 10^{-3}\frac{m_{ave}(GeV)}{500}, \quad (2)$$

where m_{ave} denotes the average squark mass, and $(m^d_{12})^2_{LL,RR}$ denotes the (1,2) entry of left-handed and right-handed down sector squark mass squared matrices in the super-CKM basis, respectively. The latter $(m^d_{12})^2_{LL,RR}$ is given by multiplying the corresponding mass squared difference between the first and second squarks by diagonalizing matrices $(V^d_{L,R})_{12}$ of the left-handed and right-handed quarks. Thus, the mass difference between the first and second squarks is severely constrained. The other mass differences are constrained more weakly or there is no experimental constraint.

Hence, it is important to study sfermion masses in the particle mixing scenario through the FI breaking, because that happens in generic string models and moreover that has the possibility for deriving realistic quark/lepton mass matrices. One of the important aspects in our flavour scenario is that the FI breaking generates D-term contributions to soft scalar masses [16 – 22]. Furthermore, some physical particles may appear combined with other states, and this introduces another modification to scalar masses [19]. The aim of this paper is to point out that even if a three-generation standard-like model has originally flavour-independent soft scalar masses, the particle mixing contribution generates non-universality among them. Of course, depending on the size of these contributions, the FCNC problem which was apparently absent in the original stringy state basis may reappear after the particle mixing.

Let us consider the simplest example, a model with Yukawa superpotential

$$W = (\lambda_1 C_1 f + \lambda_2 C_2 \xi_2)\,\xi_1\,, \quad (3)$$

where f denotes a (would-be) standard model matter field, $\xi_{1,2}$ denote two extra matter fields (triplets, doublets or singlets), $C_{1,2}$ are the fields developing large VEVs denoted by $\langle C_{1,2}\rangle = c_{1,2}$, and $\lambda_{1,2}$ are the Yukawa couplings. In the following we will use the notation $\tilde{c}_{1,2} \equiv \lambda_{1,2}c_{1,2}$. Clearly the 'old' physical particle f will combine with ξ_2. It is now straightforward to diagonalize the scalar squared mass matrix arising from the supersymmetric mass terms in eq. (3), $|\partial W/\partial\phi_\alpha|^2$, and the soft scalar masses, $m^2_\alpha|\phi_\alpha|^2$,

$$|\tilde{c}_1 f + \tilde{c}_2\xi_2|^2 + m^2_f|f|^2 + m^2_{\xi_2}|\xi_2|^2 + \left(|\tilde{c}_1|^2 + |\tilde{c}_2|^2 + m^2_{\xi_1}\right)|\xi_1|^2\,, \quad (4)$$

to find two very massive (one of them is trivially ξ_1) and one light combination. The latter, f', is obviously due to the mixing between f and ξ_2, and has a mass

$$m^2_{f'} = \frac{1}{2}\left\{m^2_f + m^2_{\xi_2} + \left(m^2_f - m^2_{\xi_2}\right)\frac{|\tilde{c}_2|^2 - |\tilde{c}_1|^2}{|\tilde{c}_2|^2 + |\tilde{c}_1|^2}\right\}\,, \quad (5)$$

where we have neglected in the above formula contributions suppressed by $O(m^2_{f,\xi_2}/|\tilde{c}_{1,2}|^2)$. It is worth noticing here that in the case $m_{f,\xi_2} = 0$, one has $m_{f'} = 0$, i.e. one recovers the limit of unbroken supersymmetry where the combination f' is massless.

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Clearly, the particle mixing generates an additional contribution to soft scalars masses, depending on the soft masses of extra matter ξ's, and also on the mixing angles $\tilde{c}$'s. For the above example (5), only in the particular case when $m_f = m_{\xi_2}$ this contribution is vanishing and one recovers $m_{f'} = m_f$. In most of realistic models the matter fields appear in three copies with the same quantum numbers reproducing the three generations of the standard model[4], and this seems to imply that there would be flavour-independent soft scalar masses, since $m_{f^i} = m_f$, $m_{\xi_2^i} = m_{\xi_2}$, and therefore $m_{f'^i} = m_{f'}$ as apparently deduced from the above example. Actually, the situation is more involved.

Let us consider the following explicit case: a Z_3 orbifold with two Wilson lines [24, 25] where all chiral fields appear automatically in three copies, and therefore three-generation standard-like models have been constructed [8−10]. For more concreteness, one introduces non-vanishing Wilson lines for the first and second T^2's, but not for the third T^2. Hence, the degeneracy of massless spectrum on three fixed points of the third T^2 is not resolved, while degeneracy for the first and second T^2's is resolved by non-vanishing Wilson lines.

Let us recall first that the FI breaking induces additional terms to soft scalar masses[5] due to F-terms, namely, the so-called D-term contribution [16−22]. In particular, the presence of an anomalous U(1) after compactification generates the dilaton-dependent FI term, that is, the D-term of the anomalous U(1) is given by

$$D^A = \frac{\delta^A_{GS}}{S+S^*} + \sum_\alpha (T+T^*)^{n_\alpha} q^A_\alpha |\phi_\alpha|^2 \ , \tag{6}$$

where the first term corresponds to the dilaton-dependent Fayet-Iliopoulos term with the GS coefficient δ^A_{GS} proportional to the value of the anomaly, and the second one is the usual D-term with the U(1) charges q^A_α of the fields ϕ_α. Then, some of these fields (with vanishing hypercharges), let us call them C_β, develop large VEVs along the D-flat direction in order to cancel the FI term, inducing the D-term contribution to the soft scalar masses of the observable fields. Totally, the soft scalar mass squared is given by [18]

$$m^2_\alpha = m^2_{3/2}\left\{1 + n_\alpha \cos^2\theta + q^A_\alpha \frac{\sum_\beta (T+T^*)^{n_\beta} q^A_\beta |C_\beta|^2 \left[(6-n_\beta)\cos^2\theta - 5\right]}{\sum_\beta (T+T^*)^{n_\beta} (q^A_\beta)^2 |C_\beta|^2}\right\}. \tag{7}$$

Here, for simplicity we have assumed that the fields C_β with VEVs have no other U(1) charges. The first two terms are the usual contributions, with the angle θ parameterizing the direction of the goldstino in the dilaton S/overall modulus T field space [26], and the modular weights with typical values $n_\alpha = -1(-2)$ for untwisted(twisted) matter fields. The third term is the D-term contribution, which is proportional to U(1) charge q^A_α. Obviously, if the observable fields have vanishing U(1) charges, $q^A_\alpha = 0$, this contribution is also vanishing. However, the observable fields have usually non-vanishing charges in explicit

[4]Recently, in [23] a new type of models has been constructed, where the three generations consist of singlets and doublets under the D_4 discrete flavour symmetry. The D_4 singlets correspond to the third generation, while the D_4 doublets correspond to the first and second generations.

[5]Let us remark that there is no additional contributions to gaugino masses and A-terms when Higgs fields relevant to such symmetry breaking have less F-term than those of dilaton and moduli fields.

models [8 − 10], and the effect of this contribution must be taken into account in the analysis. The natural size of D-term contributions is of $O(m_{3/2}^2)$, while in certain models [22] it may be enhanced.

As we can see in the above formula, the D-term contribution generates an additional non-universality among soft scalar masses, depending on q_α^A. For an illustrating example, in the simplest case that only a single field C develop a VEV in order to cancel the FI term, the above result reduces to the following form:

$$m_\alpha^2 = m_{3/2}^2 \left\{ 1 + n_\alpha \cos^2\theta + \frac{q_\alpha^A}{q_C^A} \left[(6 - n_C)\cos^2\theta - 5 \right] \right\} \,, \tag{8}$$

where q_C^A and n_C are the U(1) charge and modular weight of the field C, respectively. Notice that even in the dilaton-dominated case ($\sin\theta = 1$) the soft scalar masses are non-universal

$$m_\alpha^2 = m_{3/2}^2 \left(1 - 5\frac{q_\alpha^A}{q_C^A} \right) \,. \tag{9}$$

Nevertheless, for our realistic flavour scenario, where all matter fields appear automatically in three copies (generations), the value of q_α^A is the same for all of them. Also, since the three generations appear in the same twisted (untwisted) sector, they have the same modular weights n_α. Thus we have flavour-independent soft scalar masses. Of course, we still can have non-universal soft masses within the same generation. This is obviously harmless from the FCNC viewpoint. Let us then come back now to the particle mixing issue, and study the modifications to this result.

Since all chiral fields appear automatically in three copies in this realistic flavour scenario, the Yukawa superpotential for the example in (3) must be modified as [15]

$$\begin{aligned} W = {} & \varepsilon' g N \, (\xi_1^1 \; \xi_1^2 \; \xi_1^3) \begin{pmatrix} C_1^1 & C_1^3\varepsilon_3 & C_1^2\varepsilon_3 \\ C_1^3\varepsilon_3 & C_1^2 & C_1^1\varepsilon_3 \\ C_1^2\varepsilon_3 & C_1^1\varepsilon_3 & C_1^3 \end{pmatrix} \begin{pmatrix} f^1 \\ f^2 \\ f^3 \end{pmatrix} \\ & + \varepsilon'' g N \, (\xi_1^1 \; \xi_1^2 \; \xi_1^3) \begin{pmatrix} C_2^1 & C_2^3\varepsilon_3 & C_2^2\varepsilon_3 \\ C_2^3\varepsilon_3 & C_2^2 & C_2^1\varepsilon_3 \\ C_2^2\varepsilon_3 & C_2^1\varepsilon_3 & C_2^3 \end{pmatrix} \begin{pmatrix} \xi_2^1 \\ \xi_2^2 \\ \xi_2^3 \end{pmatrix} \,, \end{aligned} \tag{10}$$

where the flavour index i of f^i and $\xi_{1,2}^i$ correspond to three fixed points on the third T^2. The magnitude of Yukawa couplings have been calculated explicitly in heterotic orbifold models [27]. Suppressed Yukawa couplings are obtained depending on distances among fixed points[6]. Furthermore, the coupling selection rule allows only two types of combinations of fixed points on each T^2; 1) all of three correspond to the same fixed point on T^2, and 2) all of three correspond to three different fixed point each other on T^2. In the latter case, the Yukawa coupling includes the suppression factor ε_i depending on the volume of the i-th T^2 as approximately $\epsilon_i \approx 3\, e^{-\frac{2\pi}{3}T_i}$, where T_i is the moduli parameter corresponding to the volume of the i-th T^2, while the former case does not lead to such suppression

[6] More exactly, the values of the Yukawa couplings are given by a Jacobi theta function of the moduli fields.

factor. For a typical value $T_i = O(1)$ one obtains $\epsilon_i \sim 0.1$. In addition, g in (10) is the unification coupling constant and N is a quantity proportional to the root square of the volume of the unit cell for the Z_3 lattice. Also we have $gN \sim 1$. The factors $\varepsilon', \varepsilon''$ can take different values

$$\varepsilon', \varepsilon'' = 1 \ , \ \varepsilon_1 \ , \ \varepsilon_2 \ , \ \varepsilon_1\varepsilon_2 \ \ , \tag{11}$$

depending on the particular case.

Now, in order to simplify the analysis following [15], let us consider the following VEVs for the $C^i_{1,2}$ fields[7];

$$\begin{aligned} c_1^1 \equiv c_1 \ , \qquad & c_1^2 = c_1^3 = 0 \ , \\ c_2^1 = c_2^3 = 0 \ , \qquad & c_2^2 \equiv c_2 \ . \end{aligned} \tag{12}$$

Here, we expect $c_1 \sim c_2$ naturally. Then (10) gives rise to the following superpotential:

$$W = gN \left\{ \left(\varepsilon' C_1 f^1 + \varepsilon'' C_2 \varepsilon_3 \xi_2^3\right) \xi_1^1 + \left(\varepsilon' C_1 \varepsilon_3 f^3 + \varepsilon'' C_2 \xi_2^2\right) \xi_1^2 + \left(\varepsilon' C_1 \varepsilon_3 f^2 + \varepsilon'' C_2 \varepsilon_3 \xi_2^1\right) \xi_1^3 \right\} \ . \tag{13}$$

Following the discussion for (3) we can deduce straightforwardly that the masses for the three generations of physical particles f' are

$$\begin{aligned} m^2_{f'^1} &= \frac{1}{2}\left\{ m_f^2 + m_{\xi_2}^2 + \left(m_f^2 - m_{\xi_2}^2\right) \frac{|\hat{c}_2\varepsilon_3|^2 - |\hat{c}_1|^2}{|\hat{c}_2\varepsilon_3|^2 + |\hat{c}_1|^2} \right\} \ , \\ m^2_{f'^2} &= \frac{1}{2}\left\{ m_f^2 + m_{\xi_2}^2 + \left(m_f^2 - m_{\xi_2}^2\right) \frac{|\hat{c}_2|^2 - |\hat{c}_1|^2}{|\hat{c}_2|^2 + |\hat{c}_1|^2} \right\} \ , \\ m^2_{f'^3} &= \frac{1}{2}\left\{ m_f^2 + m_{\xi_2}^2 + \left(m_f^2 - m_{\xi_2}^2\right) \frac{|\hat{c}_2|^2 - |\hat{c}_1\varepsilon_3|^2}{|\hat{c}_2|^2 + |\hat{c}_1\varepsilon_3|^2} \right\} \ , \end{aligned} \tag{14}$$

where

$$\hat{c}_1 \equiv \varepsilon' c_1 \quad , \qquad \hat{c}_2 \equiv \varepsilon'' c_2 \ , \tag{15}$$

and the soft masses $m_f = m_{f^i}$, $m_{\xi_2} = m_{\xi_2^i}$, with $i = 1, 2, 3$, are given by (7).

Now the particle mixing contribution generates an additional non-universality among soft scalar masses of different generations. Thus there may appear dangerous FCNC effects from this lack of universality. The question now is whether or not these contributions can be avoided in order not to have FCNC problems. Notice that if we manage to have $m_f = m_{\xi_2}$, then universality $m_{f'^i} = m_f$ with $i = 1, 2, 3$ is recovered. In the present case of the Z_3 orbifold with soft masses given by (7) this means that the modular weights n and anomalous U(1) charges q^A of the fields f and ξ_2 must be the same. For particles in twisted sectors the modular weights are always the same, however they can have different values of q^A. Thus the lack of universality is the most general situation mainly because of the D-term contributions. Nevertheless, at least an explicit example of a standard-like model can be found [8−10] where the relevant couplings producing the mixings have the fields f and ξ_2 with the same q^A, and therefore flavour-independent scalar masses are obtained.

[7] In principle we are allowed to do this since the cancellation of the FI D-term only imposes $\sum_\alpha (T + T^*)^{n_\alpha} q^A_\alpha (|c^1_\alpha|^2 + |c^2_\alpha|^2 + |c^3_\alpha|^2) = const$, and therefore flat directions arise.

The reason for this result is that in this model there are only two possible values of q^A. Thus from (3) we deduce that $q_f^A = q_{\xi_2}^A = q_{C_1}^A = q_{C_2}^A = -q_{\xi_1}^A/2$. Therefore, this discussion implies that a way to avoid the FCNC constraints is to construct string models such that the matter fields f and ξ_2 have the same U(1) charges like in [8−10].

Clearly, this situation does not hold in generic string model. For example, we can find in appendix A of [7] two other models where more possibilities for the values of q^A for the different fields are present[8]. Hence, it is important to examine whether there is another way out to avoid the dangerous FCNC problem in a generic situation, that is, the matter fields f and ξ_2 have different U(1) charges.

Let us then discuss whether it is possible to suppress non-degeneracy of sfermion masses due to the particle mixing in these orbifold models, in the most general situation with $m_f \neq m_{\xi_2}$. For that let us consider three patterns for the values of $\hat{c}_1$, $\hat{c}_2$; 1) $\hat{c}_1 \sim \hat{c}_2$, 2) $\hat{c}_1 \ll \hat{c}_2$ and 3) $\hat{c}_1 \gg \hat{c}_2$, following [15]. When $\hat{c}_1 \sim \hat{c}_2$ one obtains from (14) the following non-universality: $m^2_{f'^1} \sim m^2_{\xi_2}$, $m^2_{f'^3} \sim m^2_f$, $m^2_{f'^2} \sim \frac{1}{2}(m_f^2 + m^2_{\xi_2})$. Clearly, depending on the U(1) charges of the fields f and ξ_2 the degree of non-universality can be important, as discussed above. This case may face the dangerous FCNC problem.

The model with $\hat{c}_1 \ll \hat{c}_2$ may also be realized. For example, this is the case when $\varepsilon'' = 1$, $\varepsilon' = \varepsilon_1\varepsilon_2$, and therefore using (15) one obtains $\hat{c}_2 = c_2$ and $\hat{c}_1 = c_1\varepsilon_1\varepsilon_2$. Since one expects $c_1 \sim c_2$, as obtained in explicit models [8−10], $\hat{c}_1$ is much smaller than $\hat{c}_2$. As a consequence, the three generations have $m^2_{f'^i} \sim m^2_f$. This result is simply understood from the fact that in this case the mixing between matter fields f^i and ξ_2^i almost vanishes for all of flavour indices i ($i = 1, 2, 3$), and that all of the three light generations approximately correspond to f^i. That is not interesting for the purpose to derive realistic fermion mass matrices, because in this case there is no particle mixing effect on fermion mass matrices of three light generations, and these matrices even after the FI breaking are almost the same as those before the FI breaking. Nevertheless, there is a subtlety in some cases, as for example when $\varepsilon'' = 1$, $\varepsilon' = \varepsilon_{1,2}$ or $\varepsilon'' = \varepsilon_1$, $\varepsilon' = \varepsilon_1\varepsilon_2$, since now $\hat{c}_1 \sim \hat{c}_2\varepsilon_3$. Then, still we have two generations, $i = 2, 3$ in (14), with $m^2_{f'^i} \sim m^2_f$, but the other has $m^2_{f'^1} \sim \frac{1}{2}(m_f^2 + m^2_{\xi_2})$. This may not give rise to a FCNC problem if we assign the first two generations of the standard model to $f'^{2,3}$, and the third one to f'^1. In this case, two of three light modes $f'^{2,3}$ approximately correspond to $f^{2,3}$, and the other light mode f'^1 to a mixture between f^1 and ξ_2^3. That would lead to non-trivial fermion mass matrices because of the particle mixing.

Finally, the third pattern arises when $\hat{c}_1 \gg \hat{c}_2$, i.e. $\varepsilon' \gg \varepsilon''$. For example with $\varepsilon'' = \varepsilon_1\varepsilon_2$, $\varepsilon' = 1$, one gets $m^2_{f'^i} \sim m^2_{\xi_2}$. In this case, all of the three light generations approximately correspond to ξ_2^i, and there is no particle mixing effect in fermion mass matrices as the previous case. Thus, this case is not interesting. As in the previous pattern there is the subtlety that for $\varepsilon'' = \varepsilon_{1,2}$, $\varepsilon' = 1$ or $\varepsilon'' = \varepsilon_1\varepsilon_2$, $\varepsilon' = \varepsilon_1$,, one still has $m^2_{f'^i} \sim m^2_{\xi_2}$ for

[8]However, these models are not realistic since fields with vanishing hypercharges cannot be found, and therefore the extra U(1) symmetries cannot be broken. Actually those two models have a very involved U(1) combination for the hypercharge (and for the anomalous U(1)), and we might speculate that realistic models should have simple U(1) combinations giving rise to the hypercharge and the anomalous U(1), and therefore a very limited possibilities for the values of q^A for the different fields.

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$i = 1, 2$, but $m^2_{f'^3} \sim \frac{1}{2}(m^2_f + m^2_{\xi_2})$. In this case, two of three light modes $f'^{1,2}$ approximately correspond to $\xi_2^{3,1}$, respectively, and the other light mode f'^3 is a mixture between f^3 and ξ_2^2. If the first two generations are assigned to $f'^{1,2}$, there may not be the FCNC problem.

Of course, we are interested in cases leading to realistic fermion mass matrices. In [15, 28] quark and lepton masses and mixing angles have been studied in the context of the Z_3 orbifold with mixing between fields due to the FI breaking. In particular, interesting results have been obtained for the structure of quark mass matrices when $\hat{c}_1 \gg \hat{c}_2$, the first two generations correspond to $f'^{1,2} \sim \xi_2^{3,1}$ and the third generation corresponds to the mixture between f^3 and ξ_2^2. Such a case may be free from the FCNC problem as discussed above.

To summarize, we have studied sfermion masses of the flavour structure that all of matter fields appear originally as three copies and the particle mixing through the FI breaking leads to non-trivial quark/lepton mass matrices. These studies are important because such particle mixing usually happens in generic string model, and that is one of the scenarios to derive realistic quark/lepton mass matrices. Our result shows that although sfermion masses are flavour-independent in the original basis, light modes after particle mixing, in general, have flavour-dependent sfermion masses mainly due to the D-term contributions. Therefore, this type of the flavour scenario may face the FCNC problem. One way to avoid it is to construct string models such that mixed states corresponding to three light generations have the same U(1) charges. Another way out is to consider the situation that particle mixing effects are negligible in the first two generations and these fields approximately correspond to the original fields, while in the third generation particle mixing effects are large enough to lead to non-trivial fermion mass matrices. In this case, the sfermion masses between the first and second generations are degenerate. Indeed, it has been shown in [15] that this case leads to interesting quark mass matrices.

Here we have assumed that the scalar fields C's have vanishing U(1) charges except the anomalous U(1), and break only the anomalous U(1) symmetry. However, in generic string models such scalar fields have non-vanishing charges other than only one U(1) symmetry, and anomaly-free and anomalous U(1) symmetries as well as non-abelian symmetries are broken at the same time. Then, complicated D-term contributions are induced as shown in [18]. However, each D-term contribution is proportional to U(1) charges of broken symmetries. For such a complicated case we can repeat our analysis, and the situation is almost the same as in the simple case we have studied here. That is, sizable non-universal sfermion masses are, in general, induced through the particle mixing mainly due to the D-term contributions, even when all of matter fields appear originally as three copies and sfermion masses are flavour-independent before the particle mixing. One way to avoid the FCNC problem is to construct string models that mixed states have the same U(1) charges for symmetries leading to large D-term contributions. Another way is to consider the situation that the particle mixing effects are negligible for the first two generations, while the particle mixing in the third generation would lead to non-trivial quark/lepton mass matrices.

Let us finally remark that, even in the cases where flavour non-universality is present at the string scale, this not need to be necessarily a problem [26, 29] if the effect is substantially

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diluted when taking into account the flavour-independent gluino mass contribution to the renormalization group running. For example, the squark masses squared receive radiative corrections due to the gluino mass $M_{1/2}$ as $\Delta m^2 \sim 6.7 \times M_{1/2}^2$. Thus, if the gluino mass is large sufficiently compared with the D-term contributions, non-universality of sfermion masses through the particle mixing mainly due to the D-term contributions would not be severely dangerous.

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