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Flavour Constraints on Sequestered Supersymmetry Breaking in String Models

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In a higher world it is otherwise,
but here below to live is to change,
and to be perfect is to have
changed often.

John Henry Newman
Conscience, Consensus and the
Development of Doctrine

Abstract

The Standard Model (SM) is the best theory which describes Nature at the fundamental level. Even if it provides very accurate predictions, the SM features some open problems. A promising extension of the SM is the Minimal Supersymmetric Standard Model (MSSM) provided by supersymmetry. The MSSM can provide solutions for some of the open problems of the SM but it can suffer from the *flavour supersymmetric problem* associated with the presence of large Flavour Changing Neutral Currents (FCNCs). The flavour supersymmetric problem is related to the mechanism which mediates supersymmetry breaking from the hidden to the visible MSSM sector. A popular supersymmetry breaking mediation mechanism is *gravity mediation* but generic gravity mediated models lead to large FCNCs. The conditions to avoid FCNCs in gravity mediated models are called *mirror mediation*. These conditions can naturally be satisfied in 4D string compactifications. 4D string models introduce new scalar particles, the *moduli*, which interact gravitationally with matter and represent natural candidates to mediate supersymmetry breaking to the MSSM. Promising 4D string scenarios are *sequestered models* where the MSSM lives on branes at singularities and the visible sector is sequestered from the sources of supersymmetry breaking in the bulk of the extra dimensions. So one can realise low-energy supersymmetry and all moduli can be heavy enough to avoid any cosmological moduli problem. In this thesis we shall focus on 4D sequestered string scenarios and determine which models can reproduce mirror mediation without the production of large FCNCs. We will find two different classes of sequestered models where only one can be compatible with present flavour constraints on FCNCs. This comparison with observations will provide information on important details of the microscopic theory like the functional dependence on the extra-dimensional volume of the physical Yukawa couplings and the Kähler potential for matter fields.

Contents

Introduction

The Standard Model (SM) of particle physics is the most advanced theoretical description of Nature providing the most sophisticated and precise predictions of all physics. There are however still many questions about fundamental particle physics which are not answered in the SM (gravity, dark matter, the cosmological constant problem, neutrino masses) and there are also many theoretical issues (hierarchy problem, the huge number of free parameters and unification of the fundamental interactions) that provide a good reason to theoretical physicists to try to find a theory beyond the SM. Space-time supersymmetry can provide answers for some of the SM problems (like neutralino dark matter, a stable Higgs mass, gauge coupling unification and radiative electro-weak symmetry breaking) but it could provide also new problems in the phenomenology by the introduction of the new particles. One of these possible problems is the flavour supersymmetric problem which consists in the presence of Flavour Changing Neutral Currents (FCNCs) in the Minimal Supersymmetric Standard Model (MSSM) by the introduction of new particles which carry flavour. Other problems are the presence of colour and charge breaking (CCB) vacua and the unbounded from below (UFB) scalar potential, caused again by the introduction of new scalar particles which can provide new competitive vacuum solutions. The mentioned problems have been heavily studied by supersymmetry theorists encoding the boundaries that a supersymmetric extension of the SM should respect (as studied in Ref. [25,27,29]) to solve these problems. The most important issue of supersymmetry is that it is not realized in Nature, otherwise we would have already seen the superpartners of the SM particles with the same mass. Therefore supersymmetric models have to include a supersymmetry breaking mechanism, and these mechanisms influence the phenomenology of the theory including also the mentioned problems. Even if the supersymmetry has not been already observed (supersymmetric theories could be observed by detecting superpartners at the LHC if their mass is of order the TeV scale) it is still a matter of studies because it can be seen as a consequence of a more fundamental theory: string theory.

String theory is the most ambitious theory for fundamental physics and quantum gravity that describes particles as states of an oscillating string. However, for consistency, the theory should live in 10 dimensions forcing the extra dimensions to be very small (compactified). String theory compactifications lead to a landscape of 4D solutions. Some of these vacua can allow for low energy supersymmetric models that can be very promising for phenomenological applications. In particular, we will focus on a corner of the string landscape where the 4D low-energy effective field theory features an MSSM visible sector with gravity mediated supersymmetry breaking. A potential problem of these string compactifications is the so-called "no scale" structure which leads to the presence of new massless fields which act as a fifth force. The Large Volume Scenario (LVS) is a class of

string compactifications that provide a successful stabilisation of the mentioned flat directions with interesting phenomenological applications to supersymmetry breaking and its mediation to the visible MSSM sector living on D-branes. This is the reason why string LVS models are today object of studies not only for particle physics phenomenology (Ref.[20]). In the LVS framework, sequestered models allow to keep the mass scale of the visible sector superpartners far below the mass of the hidden sector fields, solving automatically the so-called cosmological moduli problem (CMP).
The aim of this thesis is to analyse the amount of flavour changing neutral currents produced in sequestered supersymmetry breaking models within the framework of type IIB LVS string compactifications where supersymmetry is broken by the F-terms of the Kähler moduli whose interaction with MSSM fields mediates this breaking to the visible sector generating non-zero masses for all supersymmetric particles. This thesis is organized into three chapters:

- Chapter 1 is a brief introduction of the basic concepts of supersymmetric field theories and the fundamental principles of superstring theory and type IIB compactifications, paying particular attention to phenomenological models as the Large Volume scenario in order to understand the analysed model.

- Chapter 2 is a brief explanation of supersymmetry breaking and a review of the status of the art of flavour changing neutral currents in the Standard Model and in the Minimal Supersymmetric Standard Model, with the different solutions offered by different mediation mechanisms of supersymmetry breaking. We shall concentrate on the condition for gravity mediated supersymmetry breaking to avoid dangerously large FCNCs. In this chapter we also describe the stricter constraints imposed by the UFB potential and CCB vacua.

- Chapter 3 presents the main results of this thesis, providing a detailed analysis which reproduces the computation for the soft parameters of the two scenarios generated by the sequestered supersymmetry breaking model of [39]. We then analyse if these scenarios respect or not the flavour constraints and the UFB and CCB bounds. We also compare this result with the literature finding a light discrepancy. After determining the acceptability of the model we try to understand its phenomenological applicability.

In the conclusion section, we try to understand if the flavour bounds on the different scenarios provided by the model could give us some clues about the fundamental theory, in particular trying to understand the validity of the hypothesis made in Ref.[16] about the moduli dependence of the physical Yukawa couplings and the Kähler metric for MSSM matter fields.

Chapter 1

Supersymmetry and Superstring Theory

1.1 Supersymmetry

Supersymmetry was historically developed in the string theory framework in order to extend the bosonic string to describe also the fermionic particles producing the so-called Superstring theory. This was made starting from the work of Ramond in 1971 [1]. Only in the second half of the 1970s supersymmetry was applied to field theory resulting in the first supersymmetric field theory proposed by Freedman, Van Nieuwenhuizen, Ferrara in 1976 [2].

1.1.1 One reason for supersymmetry

Supersymmetry can be used to solve some of the open problems of the Standard Model of particle physics like the stabilisation of the Higgs mass [3]. Considering the regularized correction to the mass of the Higgs boson:

$$\Delta m_H^2 = -\frac{|\lambda_f|^2}{8\pi^2}\Lambda_{UV}^2 + ... \tag{1.1}$$

if one looks at the coupling of the Higgs field with a heavier scalar particle the correction is:

$$\Delta m_H^2 = \frac{\lambda_s}{16\pi^2}\left[\Lambda_{UV}^2 - 2m_S \ln(\Lambda_{UV}/m_S) + ...\right]. \tag{1.2}$$

If now one considers a heavier fermionic particle which does not directly couple with the Higgs field the correction due to the gauge interactions is:

$$\Delta m_H^2 = C_H T_F \left(\frac{g^2}{16\pi^2}\right)^2 \left[a\Lambda_{UV}^2 + 24 m_F^2 \ln(\Lambda_{UV}/m_F) + ...\right]. \tag{1.3}$$

So even if one rejects the physical interpretation of the cut off Λ_{UV} the mass of the Higgs is affected by corrections proportional to the masses of possible heavier particles. One possible solution is to tune the counterterm to remove the possible contributions induced by heavier masses, but it may seem unnatural. Supersymmetry provides a natural solution. Let us consider a soft supersymmetry breaking model. In this case, the contribution is:

$$\Delta m_H^2 = m_{soft}^2 \left[\frac{\lambda}{16\pi^2} \ln(\Lambda_{UV}/m_{soft}) + ... \right], \tag{1.4}$$

with m_{soft} the mass of the lightest superparticle. In this way the quadratic dependence on the regulator Λ_{UV} is removed and, to keep the mass of the Higgs at the experimental value, the mass of the superparticle should be around the TeV scale. This provides an experimental constrain to the supersymmetric models and a solution to the instability of the Higgs mass.

1.1.2 Super-Poincaré Group

Supersymmetry can be presented starting from the extension of the concept of Lie algebra. The Lie superalgebra of the operators $\mathcal{O}_a$ is defined as:

$$\mathcal{O}_a\mathcal{O}_b - (-1)^{\eta_a\eta_b}\mathcal{O}_b\mathcal{O}_a = iC_{ab}^l\mathcal{O}_l \tag{1.5}$$

with:

$$\eta_a = \begin{cases} 0 & \text{if } \mathcal{O}_a \text{ is bosonic operator} \\ 1 & \text{if } \mathcal{O}_a \text{ is fermionic operator.} \end{cases} \tag{1.6}$$

Spacetime supersymmetry can be then introduced as an extension of the Poincaré group with the introduction of new fermionic generators.
The generators of the Poincaré group are: P_μ, $L_{\mu\nu}$. The algebra of the Poincaré group is given by:

$$\begin{aligned} &[P_\mu, P_\nu] = 0 \\ &[L_{\mu\nu}, P_\rho] = -i\hbar g_{\mu\rho}P_\nu + i\hbar g_{\nu\rho}P_\mu \\ &[L_{\mu\nu}, L_{\rho\sigma}] = -i\hbar g_{\mu\rho}L_{\nu\sigma} + i\hbar g_{\mu\sigma}L_{\nu\rho} - i\hbar g_{\nu\sigma}L_{\mu\rho} + i\hbar g_{\nu\rho}L_{\mu\sigma}. \end{aligned} \tag{1.7}$$

The new generators Q_α belong to the fundamental representation of the group $SL(2,\mathbb{C})$ (they are left-handed Weyl spinors) and respect the following anti-commutation relation with the elements of the conjugate representation (right-handed Weyl spinors) $\overline{Q}_{\dot\beta}$:

$$\left\{Q_\alpha, \overline{Q}_{\dot\beta}\right\} = 2(\sigma^\mu)_{\alpha\dot\beta}P_\mu. \tag{1.8}$$

The commutation rules with the others generators of the Poincaré group are:

$$\begin{aligned} &[Q_\alpha, L^{\mu\nu}] = (\sigma^{\mu\nu})_\alpha^{\ \beta} Q_\beta \\ &[Q_\alpha, P^\mu] = \left[\overline{Q}^{\dot\alpha}, P^\mu\right] = 0 \\ &\{Q_\alpha, Q_\beta\} = 0. \end{aligned} \tag{1.9}$$

The action of Q_α on a state is to transform it from bosonic to fermionic and vice versa:

$$Q_\alpha \left|B\right\rangle = \left|F\right\rangle. \tag{1.10}$$

The anti-commutation relation (1.8) shows that by double applying the generators to a state it gives back the same initial state but translated and this is the reason why it is a spacetime symmetry.

1.1.3 Supermultiplet

In order to describe the representations of the super-Poincaré algebra with $\mathcal{N} = 1$ supersymmetry one has to define the Casimir operators such that they are invariant under supersymmetry transformation. The usual Casimir operators of the Poincaré group are:

$$\begin{aligned} C_1 &= P^\mu P_\mu \\ C_2 &= W^\mu W_\mu, \qquad W_\mu = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma} P^\nu L^{\rho\sigma}. \end{aligned} \tag{1.11}$$

The first mass Casimir can be still used but the second one based on the Pauli-Lubanski vector W_μ is no more invariant under the super-Poincaré group, as one can see from the equation (1.9). To solve this problem it has been defined a new Casimir operator:

$$\begin{aligned} B_\mu &:= W_\mu - \frac{1}{4}\overline{Q}_{\dot\alpha}(\overline{\sigma}_\mu)^{\dot\alpha\beta} Q_\beta \\ C_{\mu\nu} &:= B_\mu P_\nu - B_\nu P_\mu \\ \tilde{C}_2 &:= C_{\mu\nu} C^{\mu\nu}. \end{aligned} \tag{1.12}$$

We will consider only the massless representation because it is the one which forms the multiplets of the Standard Model. Let us fix the eigenvalue of P_μ $p_\mu = (E, 0, 0, E)$, in this choice of the reference system both C_1 and $\tilde{C}_2$ are equal to zero. Let us consider now the anti-commutation relation (1.8):

$$\left\{Q_\alpha, \overline{Q}_{\dot\beta}\right\} = 2(\sigma^\mu)_{\alpha\dot\beta} P_\mu = 2E(\sigma^0 + \sigma^3)_{\alpha\dot\beta} = 4E \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix}_{\alpha\dot\beta} \tag{1.13}$$

which means that $Q_2 = 0$.
The algebra (1.13) of Q_1 is the one of a fermionic harmonic oscillator, this can be seen by normalizing the generators:

$$a := \frac{Q_1}{2\sqrt{E}}, \qquad a^\dagger := \frac{\overline{Q}_{\dot 1}}{2\sqrt{E}} \tag{1.14}$$

obtaining the usual anti-commutation relation:

$$\left\{a, a^\dagger\right\} = 1. \tag{1.15}$$

The particles states in the supermultiplet are labelled by $|p^\mu, \lambda\rangle$. One can derive the commutation relation between J_3 (the third generator of the rotations) and the annihilation operator a using the spinorial nature of a obtaining $[a, J_3] = \frac{1}{2}a$ [4]. From the relation is possible to see the action of the annihilation operator:

$$J_3(a\,|p^\mu, \lambda\rangle) = \left(\lambda - \frac{1}{2}\right)|p^\mu, \lambda\rangle \tag{1.16}$$

and similarly for $a^\dagger$. So the action of the operator a destroys one half unity of helicity while the action of the operator and $a^\dagger$ creates one half unity of helicity. The supermultiplet can be built starting from the Clifford vacuum $|\Omega\rangle$, defined such that $a\,|\Omega\rangle = 0$

and by the commutation relation $a^\dagger a^\dagger |\Omega\rangle = 0$. The only states in the supermultiplet will be:

$$|\Omega\rangle = |p^\mu, \lambda\rangle \,, \qquad a^\dagger |\Omega\rangle = |p^\mu, \lambda + 1/2\rangle \,. \tag{1.17}$$

Adding also the CPT conjugate the full supermultiplet will be:

$$|p^\mu, \pm\lambda\rangle \,, \qquad |p^\mu, \pm(\lambda + 1/2)\rangle \,. \tag{1.18}$$

The resulting supermultiplet will contain for each half integer helicity particle also one integer helicity particle.

1.1.4 Superfields

The supersymmetric quantum field theories are based on the concept of superfield which is an extension of the usual field definition, with a well defined law of transformation under the action of the super-Poincaré group. The superfields are defined not on the usual spacetime but on the superspace that can be defined starting from a general element of the super-Poincaré group:

$$g = \exp\left[i\left(\omega^{\mu\nu}L_{\mu\nu} + a^\mu P_\mu + \theta^\alpha Q_\alpha + \overline{\theta}_{\dot\alpha}\overline{Q}^{\dot\alpha}\right)\right]. \tag{1.19}$$

Using the parameters of the supersymmetry transformation one can define the superspace as the coset:

$$\mathcal{M}_{superspace} = \left\{\omega^{\mu\nu}, a^\mu, \theta^\alpha, \overline{\theta}_{\dot\alpha}\right\} / \left\{\omega^{\mu\nu}\right\} \tag{1.20}$$

where θ^α and $\overline{\theta}_{\dot\alpha}$ are Grassmann variables. The superfields can then be expressed as Taylor expansion on the Grassmann variables, the most general one is the scalar superfield:

$$\begin{aligned} S(x^\mu, \theta_\alpha, \overline{\theta}_{\dot\alpha}) =& \varphi(x) + \theta\psi(x) + \overline{\theta}\overline{\chi} + \theta\theta M(x) + \overline{\theta\theta}N(x) + (\theta\sigma^\mu\overline{\theta})V_\mu(x) \\ &+ (\theta\theta)\overline{\theta\lambda}(x) + (\overline{\theta\theta})\theta\rho(x) + (\theta\theta)(\overline{\theta\theta})D(x) \end{aligned} \tag{1.21}$$

it is not a irreducible representation of the super-Poincaré group but it is the most general possible superfield, the variations of the components under the supersymmetry transformation are:

$$\begin{aligned} &\delta\varphi(x) = \epsilon\psi(x) + \bar\epsilon\bar\chi(x), && \delta\psi(x) = 2\epsilon M + \sigma^\mu\bar\epsilon\left(i\partial_\mu\varphi - V_\mu\right), \\ &\delta\bar\chi(x) = 2\bar\epsilon N - \epsilon\sigma^\mu\left(i\partial_\mu\varphi - V_\mu\right), && \delta M = \bar\epsilon\bar\lambda - \frac{i}{2}\partial_\mu\psi\sigma^\mu\bar\epsilon, \\ &\delta N = \epsilon\rho + \frac{i}{2}\epsilon\sigma^\mu\partial_\mu\bar\chi, && \delta D = \frac{i}{2}\partial_\mu(\epsilon\sigma^\mu\bar\lambda - \rho\sigma^\mu\bar\epsilon), \\ &\delta V_\mu = \epsilon\sigma_\mu\bar\lambda + \rho\sigma_\mu\bar\epsilon + \frac{i}{2}(\partial^\nu\psi\sigma_\mu\bar\sigma_\nu\epsilon - \bar\epsilon\bar\sigma_\nu\sigma_\mu\partial^\nu\bar\chi), \\ &\delta\bar\lambda = 2\bar\epsilon D + \frac{i}{2}(\bar\sigma^\nu\sigma^\mu\bar\epsilon)\partial_\mu V_\nu + i\bar\sigma^\mu\epsilon\partial_\mu M, \\ &\delta\rho = 2\epsilon D - \frac{i}{2}(\sigma^\nu\bar\sigma^\mu\epsilon)\partial_\mu V_\nu + i\sigma^\mu\bar\epsilon\partial_\mu N. \end{aligned} \tag{1.22}$$

It is important to notice that the D field variation is a total derivative so, D is not a dynamical field but it is an auxiliary field that can be removed from the equation of motion and it forms the famous so-called D-term.

The superfields which are irreducible representations of the super-Poincaré group are:

- Chiral, defined such that $\overline{\mathcal{D}}_{\dot{\alpha}}\Phi = 0$ ($\overline{\mathcal{D}}_{\dot{\alpha}} := -\overline{\partial}_{\dot{\alpha}} - i\theta^{\beta}(\sigma^{\mu})_{\beta\dot{\alpha}}\partial_{\mu}$):

$$\begin{aligned}\Phi(x^{\mu},\theta_{\alpha},\overline{\theta}_{\dot{\alpha}}) =&\varphi(x) + \sqrt{2}\theta\psi(x) + \theta\theta F^{1}(x) + i\theta\sigma^{\mu}\overline{\theta}\partial_{\mu}\varphi(x)\\ &- \frac{i}{\sqrt{2}}(\theta\theta)\partial_{\mu}\psi(x)\sigma^{\mu}\overline{\theta} - \frac{1}{4}(\theta\theta)(\overline{\theta\theta})\partial_{\mu}\partial^{\mu}\varphi(x),\end{aligned} \tag{1.23}$$

it is possible also to define an anti-chiral superfield such that $\mathcal{D}_{\alpha}\bar{\Phi} = 0$, it is important also to notice that any holomorphic function of a chiral superfield is a chiral superfield.

- Vector or real, defined such that $V(x,\theta,\overline{\theta}) = V^{\dagger}(x,\theta,\overline{\theta})$

$$\begin{aligned}V(x^{\mu},\theta_{\alpha},\overline{\theta}_{\dot{\alpha}}) =&C(x) + i\theta\chi(x) - i\overline{\theta}\overline{\chi} = \frac{i}{2}\theta\theta(M(x)+iN(x)) - \frac{i}{2}\overline{\theta\theta}(M(x)-iN(x))\\ &+ \theta\sigma^{\mu}\overline{\theta}V_{\mu}(x) + i\theta\theta\overline{\theta}\left(-i\overline{\lambda}(x) + \frac{i}{2}\overline{\sigma}^{\mu}\partial_{\mu}\chi(x)\right) +\\ &- i\overline{\theta\theta}\theta\left(i\lambda(x) - \frac{i}{2}\sigma^{\mu}\partial_{\mu}\overline{\chi}\right) + \frac{1}{2}(\theta\theta)(\overline{\theta\theta})\left(D(x) - \frac{1}{2}\partial_{\mu}\partial^{\mu}C(x)\right).\end{aligned} \tag{1.24}$$

Notice that if Λ is a chiral superfield it is possible to define the gauge transformation:

$$V' = V - \frac{i}{2}\left(\Lambda - \Lambda^{\dagger}\right) \tag{1.25}$$

with the non abelian generalization:

$$e^{2qV'} = e^{iq\Lambda^{\dagger}}e^{2qV}e^{-iq\Lambda}. \tag{1.26}$$

1.1.5 Action

The dynamics of the superfields is described by, as for the usual fields, an action, whose variation under the supersymmetry transformation is negligible. In details it is based on the D and F terms of respectively a real superfield named Kähler potential (which is a real arbitrary function of Φ and $\Phi^{\dagger}$) and a chiral superfield named superpotential (which is an arbitrary holomorphic function of Φ). The action is defined as:

$$S = \int d^4x \int d^4\theta (K(\Phi,\Phi^{\dagger}) + \xi V) + \int d^4x \int d^2\theta (W(\Phi) + f(\Phi)W^{\alpha}W_{\alpha} + h.c.). \tag{1.27}$$

- $K(\Phi,\Phi^{\dagger})$ is the Kähler potential with the general form $K = \Phi^{\dagger}e^{2qV}\Phi$ where Vs are the gauge fields and Φs are the matter fields.

- $W(\Phi)$ is the superpotential that can be Taylor expanded as function of Φ

$$W(\Phi) = W_0 + \lambda\Phi + \frac{m}{2}\Phi^2 + \frac{g}{3}\Phi^3. \tag{1.28}$$

[1]The transformation under supersymmetry of F is a total derivative $\delta F = i\sqrt{2}\overline{\epsilon}\overline{\sigma}^{\mu}\partial_{\mu}\psi$, this makes F an auxiliary field that can be removed, we talk about F-term of a chiral superfield.

- W_α is the field strength for the gauge fields and it is defined:

$$W_\alpha := -\frac{1}{4}(\overline{\mathcal{D}\mathcal{D}})\mathcal{D}_\alpha V, \tag{1.29}$$

  with the non abelian generalization:

$$W_\alpha := -\frac{1}{8q}\overline{\mathcal{D}\mathcal{D}}\left(e^{-2qV}\mathcal{D}_\alpha e^{2qV}\right). \tag{1.30}$$

- $f(\Phi)$ is the gauge kinetic function which encodes the gauge coupling.
- ξV is the Fayet-Iliopoulos term only for abelian gauge theories.

The most important term of the action is the scalar potential which has two contributions, one coming from the F-term of the superpotential and the other coming from the D-term of the FI term and the Kähler potential (for gauge theories). The F-term scalar potential is defined as:

$$V_F(\varphi) = -FF^* \tag{1.31}$$

the F-term is defined by its equations of motion:

$$F = -\left.\frac{\partial W^*}{\partial \Phi}\right|_{\Phi=\varphi}. \tag{1.32}$$

The resulting scalar potential is:

$$V_F(\varphi) = \left|\frac{\partial W^*}{\partial \Phi}\right|^2_{\Phi=\varphi} \geq 0 \tag{1.33}$$

with φ the scalar field component of Φ, the scalar potential obtained is positive defined. The D-term potential, which is also positive defined, is defined as:

$$V_D(\varphi) = \frac{1}{8}\left(\xi + 2q|\varphi|^2\right)^2 \tag{1.34}$$

where q is the charge of φ under the gauge interaction.
The D-term is defined as:

$$D = \xi + 2q|\varphi|. \tag{1.35}$$

As for the SM, the scalar potential plays a very important role in the supersymmetric theories because it can describe the spontaneous breaking of the supersymmetry.
The $\mathcal{N} = 1$ supersymmetric theory is equivalent to a usual quantum field theory with the same coupling and mass for the fields which are components of the same superfield. For example the Wess-Zumino model is:

$$\begin{aligned} K &= \Phi\Phi^\dagger \\ W &= \frac{m}{2}\Phi^2 + \frac{g}{3}\Phi^3 \end{aligned} \tag{1.36}$$

computing the Lagrangian one gets:

$$\mathcal{L} = \partial_\mu\varphi\partial^\mu\varphi^* + i\bar{\psi}\bar{\sigma}^\mu\partial_\mu\psi - (m\varphi + g\varphi^2)^2 - \left(\frac{m}{2} + g\varphi\right)\psi\psi - \left(\frac{m}{2} + g\varphi^*\right)\bar{\psi}\bar{\psi}. \tag{1.37}$$

Non-Renormalization Theorem

The most important theorem in supersymmetric field theory is the non-renormalization theorem which starting from the symmetries of the system demonstrates the possible ways of renormalization of the supersymmetric action [5]. The summarized results are:

- the superpotential $W(\Phi)$ is not renormalized perturbatively (but it can be affected by non-perturbative corrections);
- the gauge kinetic function $f(\Phi)$ is renormalized only at one loop;
- the Kähler potential $K(\Phi, \Phi^\dagger)$ is renormalized perturbatively at any order;
- the FI term ξ is not renormalized.

1.1.6 Supergravity

Supergravity can be defined in the same way of gravity, which is automatically obtained making local the Poincaré symmetry trough the metric $g_{\mu\nu}(x)$. In the same way one can obtain supergravity making local the supersymmetry and obtaining the "gravitino" $\psi^\alpha_\mu(x)$, for consistency also the Poincaré symmetry should become local producing automatically gravity. The supergravity $\mathcal{N}=1$ supermultiplet includes: $(e^\mu_a, \psi^\mu_\alpha, M, b_a)$ where e^μ_a are the tetrads (vectors spanning the tangent space $g_{\mu\nu} = g_{ab}e^a_\mu e^b_\nu$), M is a complex scalar auxiliary field and b_a is a real vector auxiliary field. The action describing the pure supergravity theory is:

$$\begin{aligned} S_{SUGRA} = -\frac{1}{2}\int d^4x\sqrt{-g}\Big[R + \frac{1}{2}\varepsilon^{\mu\nu\rho\sigma}(\bar{\psi}_\mu\bar{\sigma}_\nu D_\rho\psi_\sigma - \psi_\mu\sigma_\nu D_\rho\bar{\psi}_\sigma) + \\ + \text{Auxiliary fields terms}\Big] . \end{aligned} \tag{1.38}$$

Once one introduces fields coupled with gravity in the description it has to properly covariatise the action of the theory which couple with gravity. The Kähler potential and the superpotential have to respect the Kähler symmetry:

$$\begin{aligned} K(\Phi,\Phi^\dagger) &\to K'(\Phi,\Phi^\dagger) = K(\Phi,\Phi^\dagger) + h(\Phi) + h^\dagger(\Phi^\dagger) \\ W(\Phi) &\to W'(\Phi) = e^{-h(\Phi)}W(\Phi) \end{aligned} \tag{1.39}$$

with $h(\Phi)$ and holomorphyc function of Φ.
The resulting theory has the following F-term scalar potential:

$$V_F = e^K\left(K^{i\bar{j}}D_iWD_{\bar{j}}\overline{W} - 3|W|^2\right). \tag{1.40}$$

with $D_i = \partial_i W + W\partial_i K$ the Kähler covariant derivative. The F-term will be:

$$F^i = e^{K/2}K^{i\bar{j}}D_{\bar{j}}\overline{W}. \tag{1.41}$$

The important fact is that in supergravity the F-term scalar potential is not positive defined, as in global supersymmetry, but it has the negative term $3|W|^2$ which can lower the cosmological constant.

1.1.7 Minimal Supersymmetric Standard Model

The Minimal Supersymmetric Standard Model is the minimal supersymmetric extension of the Standard Model (with $\mathcal{N} = 1$). In the MSSM for each existing particle it is associated a new particle called the superpartner which has opposite fermion-boson nature. The fields which are involved in the MSSM are exposed in Table 1.1 and 1.2.

Superfields	Spin 0	Spin 1/2	SU(3)$_C$	SU(2)$_L$	U(1)$_Y$
Q_i	$\left(\widetilde{u}_L, \widetilde{d}_L\right)_i$	$(u_L, d_L)_i$	3	2	1/6
u_i^c	$\widetilde{u}_{R,i}^c$	$u_{R,i}^c$	3	1	-2/3
d_i^c	$\widetilde{d}_{R,i}^c$	$d_{R,i}$	3	1	1/3
L_i	$(\widetilde{e}_L, \widetilde{\nu}_L)_i$	$(e_L, \nu_L)_i$	1	2	-1/2
e_i^c	$\widetilde{e}_{R,i}^c$	$e_{R,i}^c$	1	1	1
ν_i^c	$\widetilde{\nu}_{R,i}^c$	$\nu_{R,i}^c$	1	1	0
H_u	(H_u^+, H_u^0)	$\left(\widetilde{H}_u^+, \widetilde{H}_u^0\right)$	1	2	1/2
H_d	$\left(H_d^0, H_d^-\right)$	$\left(\widetilde{H}_d^0, \widetilde{H}_d^-\right)$	1	2	-1/2

Table 1.1: Chiral superfields in the MSSM.

Superfields	Spin 1/2	Spin 1	SU(3)$_C$	SU(2)$_L$	U(1)$_Y$
G	$\widetilde{G}$	G	8	1	0
W	$\widetilde{W}^\pm, \widetilde{W}^0$	$W^\pm, W^0$	1	3	0
B	$\widetilde{B}^0$	B^0	1	1	0

Table 1.2: Vector superfields in the MSSM.

The new Higgs fields are fundamental to keep the trace relation $\mathrm{Tr}U(1)_Y = 0$ that removes the gravitational anomalies.
The action of the MSSM is made by:

- the Kähler potential $K = \Phi^\dagger e^{2qV}\Phi$ where $V \equiv g_3 T^a G_a + g_2 \frac{\sigma_i}{2} W_i + g_1 \frac{Y}{2} B$ is the collection of the gauge fields and Φs are the matter fields;
- the gauge kinetic functions $f_a = \tau_a$ where $\mathrm{Re}(\tau_a) = \frac{4\pi}{g_a^2}$ which determine the gauge couplings of the theory;
- The FI term is set to zero otherwise it would break charge and color $\xi = 0$;
- The superpotential is:

$$W = Y_{\alpha\beta\gamma}\Phi^\alpha\Phi^\beta\Phi^\gamma + \mu H^u H^d + W_{BL} \tag{1.42}$$

  W_{BL} is a Barion and Lepton number cancellation term that would induce a proton decay. In order to preserve the stability of the proton in the theory, it can be removed by introducing a new discrete symmetry: the R-parity. The R-parity is a discrete $\mathbb{Z}_2$ symmetry that has eigenvalue 1 for usual particles and -1 for superpartners. $Y_{\alpha\beta\gamma}$ are the Yukawa couplings and μ is the mass term for the Higgs named μ-term.

Unification

In the MSSM another advantage is the unification of the gauge couplings, whose running is modified by the new supersymmetric particles. The superparticles introduce new loop corrections in the RG flow of the couplings resulting in the intersection of them at $M_{GUT} \sim 10^{16}$GeV as shown in Figure 1.1.

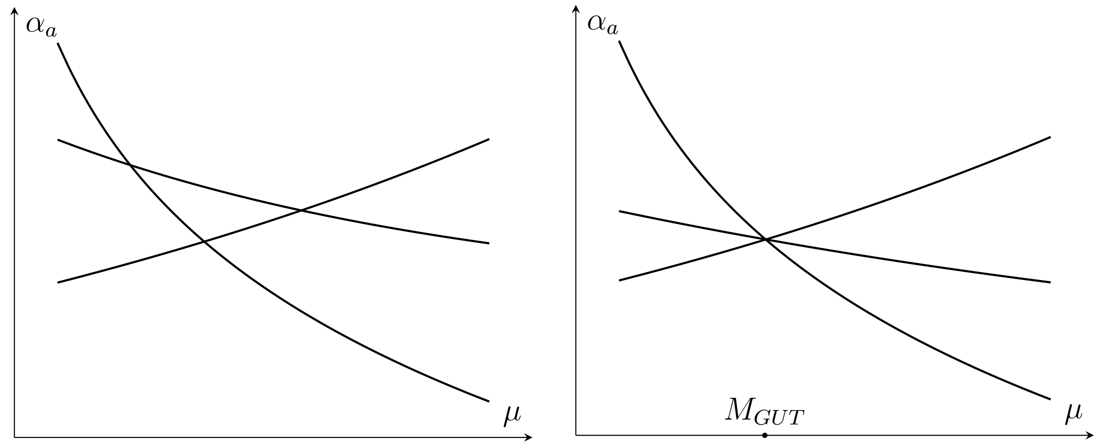

Figure 1.1: (Left) The gauge couplings qualitative behaviour in the SM. (Right) The gauge couplings qualitative behaviour in the MSSM.

The couplings unify at the value $\alpha_{GUT} \simeq 1/25$.
The big advantage of the unification is to provide a natural solution to the behaviour of the couplings in the SM, where they approach each others but do not intersect in a common value.

1.2 Superstring Theory

String theory was historically developed to describe hadrons but was yet abandoned for the Quantum Chromodynamics (QCD) which describes the fundamental interactions between quarks that form the hadrons. In 1974 Joel Scherk and John H. Schwarz decided to propose string theory as a fundamental theory for particle physics considering the particles as strings with length $l_s = l_p$. The first theory developed was able to describe only bosonic states and in fact, it is called bosonic string theory.
The bosonic string is described by the Polyakov action:

$$S = \frac{1}{4\pi\alpha'} \int d^2\sigma \sqrt{-g} g^{\alpha\beta} \partial_\alpha X^\mu \partial_\beta X^\nu \eta_{\mu\nu}, \tag{1.43}$$

with X^μ the worldsheet fields with $\mu = 0, ..., D-1$, σ are the worldsheet coordinates (τ time-like coordinate and σ the space-like coordinate along the string), $g^{\alpha\beta}$ is the worldsheet metric, $\eta_{\mu\nu}$ is the Minkowski background metric and α' is associated to the string tension and in particular it is related to the string length.
More than the absence of fermionic states the bosonic string theory has a very important issue: the ground state excitation of the string has negative mass. The mass squared is [6]:

$$M^2 = -\frac{1}{\alpha'} \frac{D-2}{6} \tag{1.44}$$

where D is the dimension of the spacetime which is not fixed. This result states that the bosonic string theory does not have a stable vacuum. The second issue came out by the first excited state mass [6]:

$$M^2 = \frac{4}{\alpha'} \left(1 - \frac{D-2}{24}\right) \tag{1.45}$$

the first exited state lies in a representation of $SO(D-2)$ and in order to preserve the Lorentz $SO(1, D-1)$ symmetry it has to be a massless representation so $D = 26$.
The introduction of fermions removes the tachyonic problem and change the number of spacetime dimensions.
To describe also fermions in the theory a fermionic field has to be introduced on the worldsheet obtaining the superstring Brink-di Vecchia-How action [7]:

$$\begin{aligned} S =& \frac{1}{4\pi\alpha'} \int d^2\sigma \sqrt{-g} \left[g^{\alpha\beta} \partial_\alpha X^\mu \partial_\beta X^\nu \eta_{\mu\nu} + \frac{i}{2} \bar{\psi}^\mu \partial \psi^\nu \eta_{\mu\nu} + \right. \\ & \left. + \frac{i}{2} (\bar{\chi}_\alpha \gamma^\beta \gamma^\alpha \psi^\mu) \left(\partial_\beta X^\nu - \frac{i}{4} \bar{\chi}_\beta \psi^\nu \right) \eta_{\mu\nu} \right] \end{aligned} \tag{1.46}$$

where ψ^μ is a D-plet of Majorana fermions and χ_A is an auxiliary Majorana field necessary for local supersymmetry. Through the gauge symmetries the auxiliary fields can be fixed out obtaining the Ramond-Neveu-Schwarz action [7]:

$$S = \frac{1}{4\pi\alpha'} \int d^2\sigma \left(\eta^{\alpha\beta} \partial_\alpha X^\mu \partial_\beta X_\mu - i \bar{\psi}^\mu \gamma^\alpha \partial_\alpha \psi_\mu \right). \tag{1.47}$$

The RNS action presents a local $\mathcal{N} = 1$ 2D supersymmetry so it is invariant under the transformation [8]:

$$\begin{aligned} \delta X^\mu &= \bar{\varepsilon} \psi^\mu \\ \delta \psi^\mu &= -i \gamma^\alpha \varepsilon \partial_\alpha X^\mu \end{aligned} \tag{1.48}$$

with ε infinitesimal spinor.

1.2.1 Normal Mode Expansion

The spectrum of the string is obtained performing a canonical quantization starting from the normal mode expansion of the string fields. The normal mode expansion can be done using the lightcone coordinates on the worldsheet:

$$\sigma_\pm = \tau \pm \sigma. \tag{1.49}$$

In this coordinates the components of the Majorana field ψ^μ are:

$$\psi^\mu(\tau,\sigma) = \begin{pmatrix} \psi_-^\mu \\ \psi_+^\mu \end{pmatrix}, \tag{1.50}$$

with a real Majorana 2D representation of the Clifford algebra:

$$\gamma^0 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \gamma^1 = \begin{pmatrix} 0 & i \\ i & 0 \end{pmatrix}. \tag{1.51}$$

In this new coordinates the action will become:

$$S = \frac{1}{4\pi\alpha'} \int d^2\sigma \left(\partial_+ X^\mu \partial_- X_\mu + i\psi_+^{\mu\dagger}\partial_-\psi_{+\mu} - i\psi_-^{\mu\dagger}\partial_+\psi_{-\mu} \right). \tag{1.52}$$

This action describes two independent excitations both for X^μ and ψ^μ fields, in fact:

$$\begin{aligned} X^\mu(\sigma,\tau) &= X_R^\mu(\sigma_-) + X_L^\mu(\sigma_+) \\ \psi^\mu(\sigma,\tau) &= \begin{pmatrix} \psi_-(\sigma_-) \\ \psi_+(\sigma_+) \end{pmatrix}. \end{aligned} \tag{1.53}$$

By the fact that we are interested only in closed strings (open strings are described in Type I string theory) the boundary conditions can be periodic or anti-periodic. For the bosonic X^μ excitations the boundary conditions are periodic:

$$X^\mu|_{\sigma=0} = X^\mu|_{\sigma=\pi}. \tag{1.54}$$

The resulting mode expansion in $\sigma_\pm$ coordinates system is [6]:

$$\begin{aligned} X_L^\mu(\sigma_+) &= \frac{1}{2}x^\mu + \frac{1}{2}\alpha' p^\mu \sigma_+ + i\sqrt{\frac{\alpha'}{2}} \sum_{n\neq 0} \frac{1}{n}\tilde{\alpha}_n^\mu e^{-in\sigma_+}, \\ X_R^\mu(\sigma_-) &= \frac{1}{2}x^\mu + \frac{1}{2}\alpha' p^\mu \sigma_- + i\sqrt{\frac{\alpha'}{2}} \sum_{n\neq 0} \frac{1}{n}\alpha_n^\mu e^{-in\sigma_-}. \end{aligned} \tag{1.55}$$

A few comments on the scalar mode expansion from [6]:

- the two excitations do not satisfy the boundary conditions individually but their sum does,

- x^μ and p^μ are the position and the momentum of the string centre of mass,
- the reality condition of X^μ imposes the constrains on the coefficients of the Fourier modes:
$$\alpha_n^\mu = (\alpha_{-n}^\mu)^*, \qquad \tilde{\alpha}_n^\mu = (\tilde{\alpha}_{-n}^\mu)^*. \tag{1.56}$$

For the spinorial field the boundary conditions can be both periodic or anti-periodic:

$$\psi_+|_{\sigma=0} = \pm\psi_+|_{\sigma=\pi}, \quad \psi_-|_{\sigma=0} = \pm\psi_-|_{\sigma=\pi}. \tag{1.57}$$

The boundary conditions are named: Ramond (R) for + sign and Neveu-Schwarz (NS) for − sign. In this way for closed strings the spectrum will present R-R, NS-NS, NS-R, R-NS sectors of different states.
The mode expansions are [7]:

$$\begin{aligned} \psi_-^\mu &= \sum_{n\in\mathbb{Z}} d_n^\mu e^{-2in\sigma_-} \quad \text{or} \quad \sum_{r\in\mathbb{Z}+1/2} b_r^\mu e^{-2ir\sigma_-} \\ \psi_+^\mu &= \sum_{n\in\mathbb{Z}} \tilde{d}_n^\mu e^{-2in\sigma_+} \quad \text{or} \quad \sum_{r\in\mathbb{Z}+1/2} \tilde{b}_r^\mu e^{-2ir\sigma_+} \end{aligned} \tag{1.58}$$

where the sum over the integers is associated to the R sector and the sum over the integers plus one half is associated to the NS sector.
The choice of a metric locally conformally flat imposes the constraints [6]:

$$T_{\alpha\beta} = 0 \tag{1.59}$$

in the lightcone gauge the constraints are:

$$\begin{aligned} T_{++} &= \partial_+ X^\mu \partial_+ X_\mu + \frac{i}{2}\psi_+\partial_+\psi_+ = 0 \\ T_{--} &= \partial_- X^\mu \partial_- X_\mu + \frac{i}{2}\psi_-\partial_-\psi_- = 0 \end{aligned} \tag{1.60}$$

$T_{+-} = T_{-+} = 0$ because the energy-momentum tensor is traceless.
From the gauge fixing of the auxiliary fields the constraints on the charge conserved under supersymmetry transformation are:

$$\begin{aligned} J_{++} &= \psi_+\partial_+ X = 0 \\ J_{--} &= \psi_-\partial_- X = 0. \end{aligned} \tag{1.61}$$

A fundamental operator in the string quantum theory is the sum of oscillator modes which for open strings are [6,7]:

$$\begin{aligned} L_m^{NS} &= \frac{1}{2}\sum_{n=-\infty}^{\infty} \alpha_{m-n}\alpha_n + \frac{1}{2}\sum_{r=-\infty}^{\infty}\left(r + \frac{1}{2}m\right) b_{-r}b_{m+r} = 0 \\ L_m^{R} &= \frac{1}{2}\sum_{n=-\infty}^{\infty} \alpha_{m-n}\alpha_n + \frac{1}{2}\sum_{n=-\infty}^{\infty}\left(n + \frac{1}{2}m\right) b_{-n}b_{m+n} = 0. \end{aligned} \tag{1.62}$$

These operators are important because in the quantum theory they will form the physical condition for states:

$$L_m |phys\rangle = 0, \quad m > 0. \tag{1.63}$$

For the 0 component there is a order ambiguity that can be solved introducing:

$$(L_0 - a) |phys\rangle = 0 \tag{1.64}$$

with a a quantity that is fixed in the quantum theory.

1.2.2 String Spectrum

The string spectrum is fundamental to understand the possible states described by the Type IIB string theory. In order to quantize the string states is important, first, to move to the lightcone gauge:

$$\begin{aligned} X^{\pm} &= \sqrt{\frac{1}{2}}(X^0 \pm X^{D-1}) \\ \psi^{\pm} &= \sqrt{\frac{1}{2}}(\psi^0 \pm \psi^{D-1}), \end{aligned} \tag{1.65}$$

with the gauge fixing conditions:

$$\begin{aligned} &X^+(\sigma, \tau) = x^+ + p^+ \tau \\ &\psi^+ = 0. \end{aligned} \tag{1.66}$$

In this gauge the physical constraints will be [7,8]:

$$\begin{aligned} \partial_+ X^- &= \frac{1}{p^+}\left(\partial_+ X^i \partial_+ X_i + \frac{i}{2}\psi^i \partial_+ \psi_i\right) \\ \psi^- &= \frac{2}{p^+}\psi^i \partial_+ X_i \end{aligned} \tag{1.67}$$

The satisfaction of the Lorentz algebra in the lightcone gauge is not forgone, in fact, for NS boundary conditions, the satisfaction means that $[J^{i-}, J^{j-}]$ has to vanish. The expression for $[J^{i-}, J^{j-}]$ is [8]:

$$\begin{aligned} \left[J^{i-}, J^{j-}\right] =&(p^+)^{-2} \sum_{m=1}^{\infty} \left[m\left(1 - \frac{D-2}{8}\right) - \frac{1}{m}\left(2a_{NS} - \frac{D-2}{8}\right)\right] \\ &(\alpha^i_{-m}\alpha^j_m - \alpha^j_{-m}\alpha^i_m) = 0. \end{aligned} \tag{1.68}$$

The result is that to preserve the Lorentz symmetry, differently of the bosonic string, $D = 10$ and $a_{NS} = 1/2$. The same computation can be done for the R states giving $a_R = 0$.

After the canonical quantization procedure (the coefficients of the Fourier modes are promoted to operators) in the lightcone gauge some of the resulting states are incompatible with the symmetries of the worldsheet. In order to keep only the states compatible with the symmetries of the worldsheet it has been defined the GSO (Gliozzi-Scherk-Olive) projectors [7]:

$$\begin{aligned} &G_{NS} = (-1)^{\sum_{r=1/2}^{\infty} b^i_{-r} b^i_r}, && \text{for NS states} \\ &G_R = \Gamma_{11}(-1)^{\sum_{n=1}^{\infty} d^i_{-n} d^i_n} && \text{for R states,} \end{aligned} \tag{1.69}$$

with b and d are the NS and R creation operators and Γ_{11} is the 10 dimensional version of the γ_5 matrix. The GSO projectors remove the states with eigenvalue -1 for NS while for the R sector the states to be removed is a matter of convention.
The states of the NS sector are bosonic while the states of the R sector are fermionic [9], so the combinations in the closed string spectrum are fermionic for NS-R and R-NS and bosonic for R-R and NS-NS.
For this thesis we are interested in the Type IIB string theory which describes the right-moving and the left-moving R sector ground states with the same chirality (producing a chiral theory). The resulting spectrum is a $\mathcal{N}=2$ and $D=10$ supergravity multiplet with the following states:

- in the NS-NS sector: graviton, two-form, dilaton;
- in the NS-R sector: gravitino, dilatino;
- in the R-NS sector: gravitino, dilatino;
- in the R-R sector: scalar, two-form, four-form.

The resulting spectrum will enter in the effective field theory.

1.2.3 Effective Field Theory

The effective field theory of Type IIB string theory can be obtaining by considering the description of the spectrum of the string ground state. The same can be obtained by considering the interactions of the string with the background fields of the ground state and requiring the preservation of the conformal symmetry also at the quantum level, imposing the β functions to be equal to zero.
The 10D $\mathcal{N}=2$ supergravity action of Type IIB string theory is [10]:

$$\begin{aligned}
S_{IIB} &= S_{NS} + S_R + S_{CS}, \\
S_{NS} &= \frac{1}{2\kappa_{10}^2}\int d^{10}x\sqrt{-G}e^{-2\Phi}\left(R + 4\partial_\mu\Phi\partial^\mu\Phi - \frac{1}{2}|H_3|^2\right), \\
S_R &= -\frac{1}{4\kappa_{10}^2}\int d^{10}x\sqrt{-G}\left(|F_1|^2 + |\tilde{F}_3|^2 + \frac{1}{2}|\tilde{F}_5|^2\right), \\
S_{CS} &= -\frac{1}{4\kappa_{10}^2}\int C_4\wedge H_3\wedge F_3,
\end{aligned} \tag{1.70}$$

with:

$$\begin{aligned}
\tilde{F}_3 &= F_3 - C_0\wedge H_3 \\
\tilde{F}_5 &= F_5 - \frac{1}{2}C_2\wedge H_3 + \frac{1}{2}B_2\wedge F_3.
\end{aligned} \tag{1.71}$$

$F_{p+1} = dC_p$ is the field strength associated to the p-form for R-R fields and $H_{p+1} = dB_p$ is the field strength associated to the p-form for NS-NS fields. The Chern-Simons term seems to brake the gauge invariance by the direct use of C_4 but thanks to the Bianchi identities of the field strengths it is gauge invariant. It is useful to see only the bosonic part of the action because it is the one which is important in the compactification procedure but there are also the fermionic partners.

1.2.4 D-brane

A p-brane can be seen as generalization of the coupling between a point-like particle and a potential:

$$\int A_\mu dx^\mu \to \int B_{M_1 \dots M_{p+1}} dx^{M_1} \wedge \dots \wedge dx^{M_{p+1}} \tag{1.72}$$

with $B_{M_1 \dots M_{p+1}}$ an antisymmetric tensor of rank $p+1$.
D-branes were initially introduced as boundary conditions for the open string, in fact, Dirichlet boundary conditions for an open bosonic string are:

$$X^i = \text{constant}, \quad i = p+1, \dots, 9. \tag{1.73}$$

As open string's boundary conditions the D-brane is an infinitely extended rigid sheet in p spacial dimensions. Later it was discovered that D-branes are dynamical objects with a central role in string phenomenology.
The spectrum of the Dp-brane can be obtained considering an open string bounded on the Dp-brane with the corresponding states in the spectrum that propagate in $(p+1)$-dimensional volume [11]. The spectrum is a representation of the reduced $SO(p-1)$ Lorentz group consisting in:

- vector gauge boson A_μ,
- $9-p$ real scalar fields ϕ^i,
- spinor λ_α,

resulting in a $U(1)$ supermultiplet with 16 supersymmetries in $p+1$ dimensions.
The effective action of a Dp-brane can be derived by the amplitude of the exchanging of closed strings between two Dp-branes. The resulting action is given by a combination of the DBI (Dirac-Born-Infeld) action for non-linear electrodynamics and a Chern-Simmons term [11,12]:

$$\begin{aligned} S_p &= S_{DBI} + S_{CS} \\ S_{DBI} &= -\frac{(\alpha')^{(p+1)/2}}{(2\pi)^p} \int_{V_{p+1}} d^{p+1}x e^{-\Phi} \sqrt{-\det(P[G+B] - 2\pi\alpha' F)} \\ S_{CS} &= \frac{(\alpha')^{(p+1)/2}}{(2\pi)^p} \int_{V_{p+1}} P\left[\sum_q C_q\right] \wedge e^{2\pi\alpha' F - B} \wedge \hat{A}(R), \end{aligned} \tag{1.74}$$

where $F_{\mu\nu}$ is the field strength of the gauge field, Φ is the dilaton in the NS-NS sector, B_2 is the two-form from the NS-NS sector, $G_{\mu\nu}$ is the background metric, C_qs are the R-R q-forms and $\hat{A}(R)$ is the A-roof polynomial defined as:

$$\hat{A}(R) = 1 - \frac{1}{24(8\pi^2)} \text{tr} R^2 + \dots \tag{1.75}$$

with R the spacetime curvature. $P[E]$ is the "pull-back" of the spacetime tensor E into the brane worldvolume defined as:

$$P[E]_{\mu\nu} = E_{\mu\nu} + E_{\mu i}\partial_\nu \phi^i + \partial_\mu \phi^i E_{i\nu} + \partial_\mu \phi^i \partial_\nu \phi^j E_{ij}, \tag{1.76}$$

with ϕ^i the scalars in the Dp-brane spectrum.
In Type IIB string theory the RR-sector is made by even-degree forms which will couple to odd p Dp-branes allowing only odd p Dp-branes.

1.3 Type IIB Compactification

The explained theory is in 10 dimensions but the physical world is in 4 dimensions. This means that the other six dimensions have to be compactified in a compact manifold so the ten-dimensional spacetime can be decomposed as the product of the usual four-dimensional Minkowski spacetime and the compactified microscopic extra dimensions:

$$\mathcal{M}_{10} = \mathcal{M}_4 \times X_6. \tag{1.77}$$

To keep a four-dimensional $\mathcal{N} = 1$ spacetime supersymmetry X_6 has to be a complex Kähler and Ricci-flat manifold, it has to be a Calabi-Yau manifold.
The 4D theory is obtained by a Kaluza-Klein reduction of the 10D theory by expanding all the fields into modes of the X_6 that respect the Laplacian equation which are in one to one correspondence with the harmonic forms on X_6 [10]. The harmonic forms on X_6 are counted by the dimension of the non trivial cohomologies of the Calabi-Yau [13]. The properties of the Calabi-Yau strongly affect its Hodge decomposition leaving only the cohomology groups expressed in Table 1.3.

Cohomology group	Dimension	Basis
$H^{(1,1)}$	$h^{(1,1)}$	ω_A
$H^{(2,2)}$	$h^{(1,1)}$	$\tilde{\omega}^A$
$H^{(3)}$	$2h^{(2,1)} + 2$	(α_G, β^G)
$H^{(2,1)}$	$h^{(2,1)}$	χ_K
$H^{(3,3)}$	1	vol

Table 1.3: Cohomology groups on X_6 and their basis.

The ansatz for the background metric in the Kaluza-Klein reduction from the 10D theory to the effective 4D theory is:

$$ds^2 = \eta_{\mu\nu} dx^\mu dx^\nu + g_{i\bar{j}}(y) dy^i dy^{\bar{j}}. \tag{1.78}$$

The fluctuations around the background metric describe the four-dimensional graviton $g_{\mu\nu}$ and the deformations of the Calabi-Yau are decomposed in $h^{(1,1)}$ real Kähler structure deformations v^A and $h^{(2,1)}$ complex structure deformations z^A.
The Kähler structure deformations are the deformations of the Kähler form $J = iG_{i\bar{j}} dy^i \wedge d\overline{y}^{\bar{j}}$ which can be expanded as:

$$J = v^A \omega_A, \qquad A = 1, ..., h^{(1,1)}, \tag{1.79}$$

and the complex structure deformations, which are in one to one correspondence with the $(1,2)$-harmonic forms, are defined by:

$$\delta g_{ij} = \frac{i}{||\Omega_3||^2} \overline{z}^K (\bar{\chi}_K)_{i\bar{i}\bar{j}} \Omega_3^{\bar{i}\bar{j}}{}_j, \tag{1.80}$$

with $||\Omega_3||^2 = \frac{1}{3!} \Omega_{3ijk} \hat{\Omega}_3^{ijk}$. The forms in the string spectrum expanded on the basis of the cohomology groups are:

$$\begin{aligned} B_2 &= \hat{B}_2(x) + b^A(x)\omega_A, \quad C_2 = \hat{C}_2(x) + c^A(x)\omega_A, \quad A = 1, ..., h^{(1,1)} \\ C_4 &= D_2^A(x) \wedge \omega_A + V^G(x) \wedge \alpha_G - U_G(x) \wedge \beta^G + \rho_A(x)\tilde{\omega}^A, \quad G = 0, ..., h^{(1,2)}. \end{aligned} \tag{1.81}$$

The resulting spectrum is a collection of $\mathcal{N}=2$ supermultiplets:

- gravity multiplet $(G_{\mu\nu}, V^0)$,
- $h^{(2,1)}$ vector multiplets (V^K, z^K),
- $h^{(1,1)}$ hypermultiplets (v^A, b^A, c^A, ρ^A),
- double-tensor multiplet $(\hat{B}_2, \hat{C}_2, \Phi, C_0)$.

Orientifold Projection

The resulting theory, from manifold compactification, has $\mathcal{N}=2$ supersymmetry and in order to truncate the supermultiplets and obtain a $\mathcal{N}=1$ supersymmetryc theory it is required to perform a orientifold projection. The orientifold projection is performed considering the Type IIB on X_6 and then modding out by the orientifold action $\Omega\mathcal{R}$, where Ω is (practically) the worldsheet parity and $\mathcal{R}$ is a $\mathbb{Z}_2$ discrete symmetry of X_6. The points fixed under $\mathcal{R}$ form the Op-planes which span the Minkowski spacetime and wrap cycles on X_6.
Before seeing the effective action let us see the supermultiplet from the orientifold projection. We will consider the O3/O7 planes landscape which arise from projection of the form $(-1)^{F_L}\Omega\mathcal{R}$ [10]. The behaviour of the bosonic states under the action of Ω and $(-1)^{F_L}$ are expressed in Table 1.4.

Field	Ω	$(-1)^{F_L}$
Φ	even	even
$G_{\mu\nu}$	even	even
B_2	odd	even
C_0	odd	odd
C_2	even	odd
C_4	odd	odd

Table 1.4: Type IIB spectrum under the two operations.

The action of $\mathcal{R}$ in the O3/O7 planes landscape is:

$$\begin{aligned} \mathcal{R}^*J &= J \\ \mathcal{R}^*\Omega_3 &= -\Omega_3 \end{aligned} \tag{1.82}$$

where J is the X_6 Kähler form, Ω_3 is the holomorphic 3-form and $\mathcal{R}^*$ is the pull-back of the map $\mathcal{R}$. To preserve the invariance of the states they have to transform as:

$$\begin{aligned} &\mathcal{R}^*\Phi = \Phi, \qquad &&\mathcal{R}^*G_{\mu\nu} = G_{\mu\nu}, \qquad &&\mathcal{R}^*B_2 = -B_2, \\ &\mathcal{R}^*C_0 = C_0, \qquad &&\mathcal{R}^*C_2 = -C_2, \qquad &&\mathcal{R}^*C_4 = C_4. \end{aligned} \tag{1.83}$$

Since $\mathcal{R}$ is an holomorphic involution the cohomology group splits into two eigenspace:

$$H^{(p,g)} = H^{(p,q)}_+ \oplus H^{(p,q)}_- . \tag{1.84}$$

Cohomology group		Dimension		Basis	
$H_+^{(1,1)}$	$H_-^{(1,1)}$	$h_+^{(1,1)}$	$h_-^{(1,1)}$	ω_α	ω_a
$H_+^{(2,2)}$	$H_-^{(2,2)}$	$h_+^{(1,1)}$	$h_-^{(1,1)}$	$\tilde{\omega}^\alpha$	$\tilde{\omega}^a$
$H_+^{(3)}$	$H_-^{(3)}$	$2h_+^{(2,1)}+2$	$2h_-^{(2,1)}+2$	$(\alpha_\gamma, \beta^\lambda)$	(α_g, β^l)
$H_+^{(2,1)}$	$H_-^{(2,1)}$	$h_+^{(2,1)}$	$h_-^{(2,1)}$	χ_κ	χ_k
$H_+^{(3,3)}$	$H_-^{(3,3)}$	1	0	vol	

Table 1.5: Cohomology groups and their basis after the orientifold projection.

By the fact that J is invariant under the action of $\mathcal{R}^*$ the only remaining Kähler deformations are the ones from $H_+^{(1,1)}$ so:

$$J = v^\alpha \omega_\alpha, \qquad \alpha = 1, ..., h_+^{(1,1)}. \tag{1.85}$$

By the opposite argument the only remaining complex structure deformations are the ones from $H_-^{(2,1)}$:

$$\delta g_{ij} = \frac{i}{||\Omega_3||^2} \bar{z}^k (\bar{\chi}_k)_{i\bar{i}\bar{j}} \Omega_3^{\bar{i}\bar{j}}{}_j, \qquad k = 1, ..., h_-^{(1,2)}. \tag{1.86}$$

The new decompositions of the forms of the string spectrum will be:

$$\begin{aligned} &B_2 = b^a(x)\omega_a, \quad C_2 = c^a(x)\omega_a, \quad a = 1, ..., h_-^{(1,1)} \\ &C_4 = D_2^\alpha(x) \wedge \omega_\alpha + V^\gamma(x) \wedge \alpha_\gamma - U_\gamma(x) \wedge \beta^\gamma + \rho_\alpha(x)\tilde{\omega}^\alpha, \gamma = 0, ..., h_+^{(1,2)}, \alpha = 1, .., h^{(1,1)}. \end{aligned} \tag{1.87}$$

The new fields are arranged in the $\mathcal{N} = 1$ 4D super-multiplets:

- gravity multiplet $g_{\mu\nu}$,
- $h_+^{(2,1)}$ vector multiplets (V^λ),
- $h_-^{(2,1)}$ chiral multiplet z^k,
- $h_-^{(1,1)}$ chiral multiplet (b^a, c^a),
- chiral multiplet (Φ, l),
- $h_+^{(1,1)}$chiral/linear multiplet (v^α, ρ_α).

The chiral supermultiplets can be arranged to form the supergravity moduli chiral superfields:

$$\begin{aligned} &T_\alpha = -i \int \mathcal{J}_C \wedge \omega_\alpha, \quad \mathcal{J}_C = C_4 + \frac{i}{2} e^{-\Phi} J \wedge J + (C_2 - iSB_2) \wedge B_2, \\ &U_\gamma = i \int \Omega_3 \wedge \alpha_\gamma \\ &S = e^{-\Phi} + iC_0 \\ &G^a = c^a - iSb^a. \end{aligned} \tag{1.88}$$

Kähler Potential

The resulting 4D Kähler potential for the moduli of the $\mathcal{N} = 1$ supergravity theory is [11]:

$$\begin{aligned}\kappa_4 K_{IIB} &= -\log\left(-i\int \Omega_3 \wedge \Omega_3\right) - \log\left(S+\overline{S}\right) - 2\log\left(e^{-\frac{3}{2}\Phi}\int J\wedge J\wedge J\right) = \\ &= -\log\left(-i\int \Omega_3 \wedge \Omega_3\right) - \log\left(S+\overline{S}\right) - 2\log\left(\mathcal{V}\right)\end{aligned} \tag{1.89}$$

where $\mathcal{V}$ is the volume of X_6 in the Einstein frame.
The resulting Kähler potential is the leading order in α', in fact the 10D effective theory can be affected by higher perturbation terms in α' which lead to the R^4 term [14]:

$$S_{R^4} = (\alpha')^4 \int d^{10}x t^8 t^8 R^4. \tag{1.90}$$

After the Kaluza-Klein reduction the α' modify the Kähler potential for the Kähler moduli introducing the new term [11]:

$$K_{\alpha'} = -2\log\left[e^{-\frac{3}{2}\Phi}\left(\int J\wedge J\wedge J + \frac{\xi}{2}\right)\right] = -2\ln\left(\mathcal{V} + \frac{\xi \mathrm{Re}(S)^{3/2}}{2}\right) \tag{1.91}$$

with $\xi = -\frac{\zeta(3)}{16\pi^3}\chi(X_6)$ where $\chi(X_6)$ is the Euler number of X_6. This correction is important in particular models to stabilize the Kähler moduli.

Gauge Kinetic Function

The gauge kinetic function depends on the Dp-brane configuration and on the volume of the cycle Π_{p-3} of the Calabi-Yau wrapped by the Dp-brane. The flux of $F_{\mu\nu}$ associated to the Dp-brane respect a Dirac quantization condition [11]:

$$m\int_{\Pi_{p-3}} F = 2\pi n, \qquad m, n \in \mathbb{Z} \tag{1.92}$$

m^i is the number of wrappings of the Dp-brane along the Π_{p-3} cycle and n^j is the quantized the magnetic flux.
In fact if one expand the DBI action respect to the gauge field strength tensor $F_{\mu\nu}$ one gets:

$$\frac{1}{g^2_{\mathrm{D}p}} = e^{-\Phi}\frac{(\alpha')^{(3-p)/2}}{(2\pi)^{p-2}}\mathrm{Vol}(\Pi_{p-3}). \tag{1.93}$$

The result from the Kaluza-Klein dimensional reduction for magnetized D7-brane is:

$$\begin{aligned}f_a^{D7} =& \frac{(\alpha')^{-2}}{(2\pi)^5}\left[e^{-\Phi}\int_{\Pi_a}\mathrm{Re}\left(e^{-i\varphi_a}e^{J+i2\pi\alpha'(B+2\pi\alpha' F)}\right) + \right.\\ &\left. + i\int_{\Pi_a}\sum_k C_{2k}e^{2\pi\alpha'(B+2\pi\alpha' F)}\right]\end{aligned} \tag{1.94}$$

with Π_a the holomorphic 4-cycle and ϕ_a the so-called BPS phase. Considering the wrapping number m^i and the magnetic flux n^j of the D7-brane the result of the gauge kinetic function is:

$$2\pi f_i^{D7} = n_i^j n_i^k S - m_i^j m_i^k T_i, \quad i \neq j \neq k \neq i. \tag{1.95}$$

For the D3-brane the computation is much easier because, as one can see from (1.92), the volume dependence vanishes and it remains only the dilaton dependence:

$$2\pi f^{D3} = S. \tag{1.96}$$

The above description has to be extended if one wants to consider models with D-branes at the singularities but this kind of new contributions can be neglected for the purpose of this thesis.

Superpotential from Fluxes

In the previous section the compactification has been considered without allowing non-trivial geometrical fluxes for the p-forms in the string spectrum. If one allows the presence of non trivial fluxes they are quantized following:

$$\frac{1}{(2\pi)^2\alpha'}\int_\Sigma F_{(3)} = m_\Sigma, \qquad \frac{1}{(2\pi)^2\alpha'}\int_\Sigma H_{(3)} = n_\Sigma, \quad n_\Sigma, m_\Sigma \in \mathbb{Z}, \tag{1.97}$$

with Σ are 3-cycles on the Calabi-Yau X_6.
The resulting superpotential from the compactification in O3/O7 configuration is the GVW superpotential:

$$W(U,S) = \int \Omega_3 \wedge G_{(3)} \tag{1.98}$$

with the complex three form.

$$G_{(3)} = F_{(3)} - SH_{(3)}. \tag{1.99}$$

The superpotential generated depends only on the dilaton S and the complex structure moduli U_γ and is independent on the Kähler moduli.

Matter Metric

The matter metric is defined as the metric of the matter fields space, the Kähler potential generated for the matter fields is:

$$K_{matter} = \tilde{K}_{\alpha\bar{\beta}} C^\alpha \overline{C}^{\bar{\beta}} \tag{1.100}$$

where C^αs are the chiral matter fields. The value of $\tilde{K}_{\alpha\bar{\beta}}$ depends on the configuration used in the model to generate chiral fields from open strings. In the O3/O7 orientifold models the possible configurations are: stack D7-branes with transversal fluctuations, stack of D7-branes with parallel fluctuations, intersection of D7-branes, stack of D3-branes and D3-brane coinciding with D7-brane.
The exact value of $\tilde{K}_{\alpha\bar{\beta}}$ is difficult to compute, in fact, it has been calculated only for toroidal orientifold compactification [11] but for a model of branes at resolved singularities or for large volume models the form could be deduced by the geometrical properties

of the Calabi-Yau and from the properties of the superpotential [15].
Starting by the superpotential of the chiral superfields C^α which is independent on the possible configurations:

$$W_{matter} = Y_{\alpha\beta\gamma}(U,S)C^\alpha C^\beta C^\gamma \tag{1.101}$$

also, a mass term for the Higgs field is allowed from the symmetries. $Y(U,S)_{\alpha\beta\gamma}$ are the unnormalized Yukawa couplings. Due to the holomorphy of the superpotential the Yukawa couplings can depend only on the complete moduli but the imaginary part of the Kähler moduli enjoys a shift symmetry:

$$\mathrm{Im}(T) \to \mathrm{Im}(T) + a \tag{1.102}$$

to preserve this symmetry the Yukawa couplings can not depend on the Kähler moduli but in this way, they will not depend on the volume of the Calabi-Yau. The non-renormalization theorem prevents also the possibility of a dependence induced by perturbative corrections which are not allowed to the superpotential.
The physical normalized Yukawa couplings will be:

$$\hat{Y}_{\alpha\beta\gamma} = e^{K/2}\frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha\tilde{K}_\beta\tilde{K}_\gamma)^{\frac{1}{2}}}. \tag{1.103}$$

By the fact that the matter fields are localized on one of the smaller cycles, the locality implies that the physical Yukawa couplings are independent on the whole bulk geometry and so they are independent on the whole volume $\mathcal{V}$. This means that the matter metric can be obtained from the normalization condition imposing that $e^{K/2}/(\tilde{K}_\alpha\tilde{K}_\beta\tilde{K}_\gamma)^{\frac{1}{2}}$ is independent on the volume. The result is:

$$K_{\alpha\tilde{\beta}} = \frac{h_{\alpha\tilde{\beta}}(U,S)}{\mathcal{V}^{2/3}} \tag{1.104}$$

where $h_{\alpha\tilde{\beta}}(U,S)$ encodes the dependence on the dilaton and the complex structure moduli.

1.3.1 Moduli Stabilisation

The theory described in (1.73) as the no-scale structure in fact the related scalar potential does not depend on the Kähler moduli, they are flat directions of the potential leaving them massless. A well-defined model needs a mechanism to stabilize the Kähler moduli otherwise, without a mass, they will act as a new gravitational force with an unlimited range of action that would spoil the prediction for gravitational interactions.
This can be achieved by using perturbative corrections to the Kähler potential as seen in (1.75) that would brake the no-scale structure:

$$K^{T_i\overline{T}_j}K_{T_i}K_{\overline{T}_j} = 3 \tag{1.105}$$

with the correction

$$K^{T_i\overline{T}_j}K_{T_i}K_{\overline{T}_j} = 3 + \frac{3}{4}\frac{\hat{\xi}}{\mathcal{V}}, \tag{1.106}$$

or by non-perturbative corrections to the superpotential as the one arising from ED3-brane instantons [17] which take the form:

$$W^i_{n.p.} = A_i(U,S)e^{-a_i T_i}. \tag{1.107}$$

The dilaton and the complex structure moduli are stabilized by the flux induced superpotential that will provide the supersymmetric stabilisation conditions [18]:

$$D_S W = 0, \qquad D_{U^\gamma} W = 0. \tag{1.108}$$

KKLT Model

A first example of moduli stabilization trough non-perturbative effects is the KKLT model [19] whose Kähler moduli sector is described by:

$$\begin{aligned} K &= -3\ln\left[-i(T-\overline{T})\right] \\ W &= W_0^{flux} + Ae^{-aT} \end{aligned} \tag{1.109}$$

with the definition $\tau = iT$ the F-term scalar potential is defined as:

$$V_F = \frac{aAe^{-a\tau}}{2\tau^2}\left(\frac{1}{3}\tau a A e^{-a\tau} + W_0 + Ae^{-a\tau}\right). \tag{1.110}$$

As can be seen in the Figure 1.2 the F-term scalar potential provides an AdS minimum which stabilizes the Kähler moduli T.

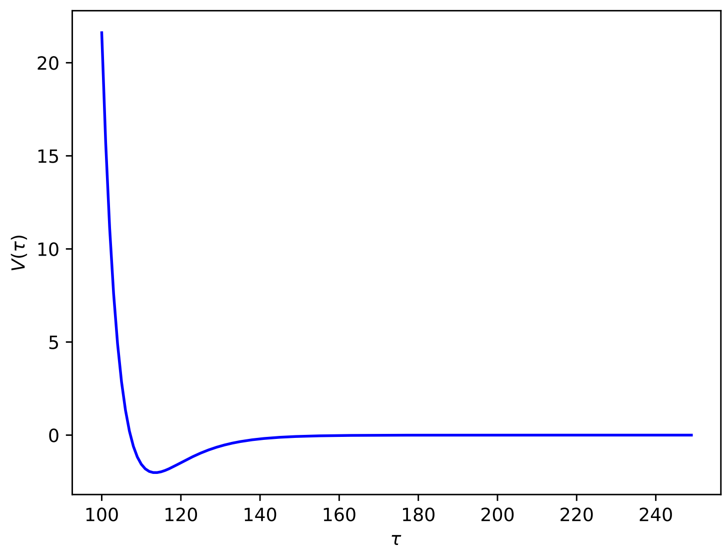

Figure 1.2: The KKLT potential (multiplied for 10^{15}) with the AdS minimum ($W_0 = -10^{-4}$, $A = 1$, $a = 0.1$).

The minimum of the model as can be seen by Figure 1.2 is a AdS one since $\langle V \rangle < 0$ but it can be uplifted through the different procedures in order to obtain a more cosmologically accurate model.

1.3.2 Large Volume Scenario

The large volume scenario is a class of string compactification that extends the KKLT model introducing a new minimum with exponentially large volume and provides the stabilisation for the Kähler moduli combining perturbative corrections to the Kähler potential with non-perturbative corrections to the superpotential.
One of the advantages of the large volume scenario can be seen by the relation $M_P^2 \sim \frac{M_S^2 \mathcal{V}}{g_s}$. In order to keep the string mass M_S hierarchically small, the large volume is fundamental. The Kähler potential takes the following form:

$$K = -2\ln\left(\mathcal{V} + \frac{\xi \mathrm{Re}(S)^{3/2}}{2}\right) - \log(S+\overline{S}) - \log\left(-i\int \Omega\wedge\overline{\Omega}\right) \tag{1.111}$$

with the volume:

$$\mathcal{V} = \frac{1}{6}\int J\wedge J\wedge J = \frac{1}{6}k_{ijk}t^i t^j t^k. \tag{1.112}$$

Where k_{ijk} are the triple intersection number of X_6 and t^i are 2-cycles volumes. The Kähler moduli can be expanded as: $T_i = \tau_i + ib_i$, the imaginary part is an axion but the real part is the 4-cycle volume which is related to the 2-cycle volume by the relation:

$$\tau_i = \frac{\partial \mathcal{V}}{\partial t^i} = \frac{1}{2}k_{ijk}t^j t^k. \tag{1.113}$$

The superpotential of the model is made by the usual flux term plus non-perturbative corrections induced by ED3 instantons for each Kähler moduli:

$$W = W_0 + \sum_i A_i e^{-a_i T_i} \tag{1.114}$$

the sum is over the cycles generating non-perturbative corrections with $a_i = 2\pi$ for brane instantons (this could be also generated by gaugino condensation producing $a_i = 2\pi/N$), A_i depends on the complex structure moduli and the position of the D3-brane.
The perturbative correction, as already said, brakes the no-scale structure and combined with the non-perturbative terms in the superpotential will provide a F-term scalar potential that stabilise the Kähler moduli [20]:

$$\begin{aligned} V =& e^K \Big[K_0^{jk}\left(a_j A_j a_k \overline{A}_k e^{-(a_j T_j + a_k T_k)} - \left(a_j A_j e^{-a_j T_j}\overline{W}\partial_{\overline{T}_k}K_0 + a_k\overline{A}_k e^{-a_k \overline{T}_k} W\partial_{T_j}K_0\right)\right) + \\ &+ 3\hat{\xi}|W|^2 \frac{\left(\hat{\xi}^2 + 7\hat{\xi}\mathcal{V} + \mathcal{V}^2\right)}{\left(\mathcal{V}-\hat{\xi}\right)\left(2\mathcal{V}+\hat{\xi}\right)^2}\Big] \end{aligned} \tag{1.115}$$

with $\hat{\xi} = \xi \mathrm{Re}(S)^{3/2}$. The large volume limit $\mathcal{V}\to\infty$ is taken respecting the condition that N_{small} Kähler moduli remain small producing the following Kähler potential and superpotential:

$$\begin{aligned} K &= K_{cs} - \ln(S+\overline{S}) - 2\ln\left(\mathcal{V} + \frac{\hat{\xi}}{2}\right), \\ W &= W_0 + \sum_{j=1}^{N_{small}} A_j e^{-a_j T_j}. \end{aligned} \tag{1.116}$$

This condition provides an AdS non-supersymmetric minimum.
In order to respect the large volume scenario the Calabi-Yau has to satisfy the conditions exposed by [20]:

1 The Euler number of the Calabi-Yau has to be negative by the fact that $\hat{\xi}$ must be positive to make the potential goes to zero from the below.

2 The Calabi-Yau must have at least one blow-up mode associated to a 4-cycle modulus that resolves a point-like singularity in order to have the non-perturbative contribution to the superpotential.

3 The 4-cycles are fixed as small by the combination of non-perturbative corrections and the α' correction.

4 All the other 4-cycles can not be small, possible non-perturbative corrections are sent to zero making the moduli big.

5 Non blow-up Kähler moduli (except the overall volume mode) remain flat directions.

6 To stabilise these moduli are crucial the string loop corrections giving g_s terms which are dominant over non-perturbative corrections.

1.3.3 Single-Hole Swiss Cheese

An example of large volume scenario is the Single-Hole Swiss Cheese $\mathbb{CP}^4_{[1,1,1,6,9]}$ Calabi-Yau which has two 4-cycles τ_s and τ_b providing the volume:

$$\mathcal{V} = \tau_b^{3/2} - \tau_s^{3/2}. \tag{1.117}$$

with τ_b associated to a big cycle and τ_s associated to a small cycle, in fact in the large volume limit $\tau_b \to \infty$ and τ_s remains small.
Neglecting the dilaton and the complex structure moduli (which can be stabilised by the fluxes) the theory is described by the following Kähler potential and superpotential:

$$\begin{aligned} K &= -2\ln\left(\mathcal{V} + \frac{\hat{\xi}}{2}\right) \\ W &= W_0 + A_s e^{-a_s T_s} \end{aligned} \tag{1.118}$$

with $\hat{\xi} = 2\xi/g_s^{3/2}$. The Kähler potential contains also the α' corrections and the superpotential is made by the flux superpotential and by the ED3-brane instanton which generate the non-perturbative correction which depends only the small cycle moduli.
For simplicity it is useful to use as variables $\mathcal{V}$ and τ_s. With these variables the F-term scalar potential is:

$$V_F = \frac{3}{4}\frac{\hat{\xi}W_0^2}{\mathcal{V}^3} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^2} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}} \tag{1.119}$$

which provides a minimum for $\mathcal{V} \sim W_0 e^{a_s \tau_s}$ and $\tau_s \sim \hat{\xi}^{2/3}$, the minimum coincides with exponentially large volume in a consistent way with the large volume limit.
The resulting minimum provides an AdS vacuum energy:

$$\langle V_F \rangle = -\frac{9\hat{\xi} W_0^2}{16\mathcal{V}^3 a_s \tau_s}. \tag{1.120}$$

The two Kähler moduli acquire mass leaving no flat directions in the potential. The two masses can be related to the gravitino mass $m_{3/2}$ (it will be explained in Chapter 2) producing the spectrum made by:

$$\begin{aligned} m_{\tau_s} &\sim m_{3/2} \\ m_{\tau_b} &\sim \sqrt{m_{3/2}^3/M_P}. \end{aligned} \tag{1.121}$$

The associated F-terms are:

$$F^{T_b} = -\frac{2\tau_b W_0}{\mathcal{V}} \sim m_{3/2}^{1/3} M_P^{5/3}, \qquad F^{T_s} = -\frac{3W_0}{2a_s \mathcal{V}} \sim m_{3/2} M_P. \tag{1.122}$$

The explicit computations are exposed in Appendix A.

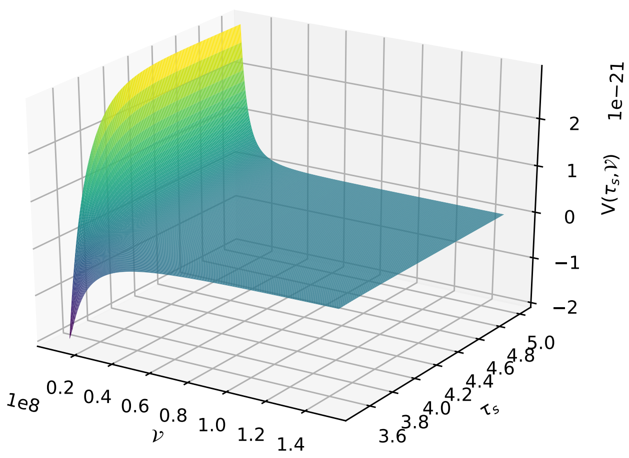

Figure 1.3: LVS potential in the single hole Swiss cheese model (with $W_0 = 20$, $A_s = 1$, $s = 2\pi$, $\hat{\xi} = 0.01$).

As one can see from the Figure 1.3 the potential exhibits an AdS for exponentially large volume.

Chapter 2

Supersymmetry Breaking and FCNC

2.1 Supersymmetry Breaking

Supersymmetry is not realized in nature so a very important feature is the breaking of supersymmetry. As for other symmetries, supersymmetry is broken if the vacuum is not invariant under supersymmetry transformation:

$$Q_\alpha \left|vac\right\rangle \neq 0. \tag{2.1}$$

If now one considers the anti-commutation relation of the supersymmetry generators and one contracts it with $\overline{\sigma}$ the result is:

$$(\overline{\sigma}^\nu)^{\dot{\beta}\alpha} \left\{Q_\alpha, \overline{Q}_{\dot{\beta}}\right\} = 4P^\nu \tag{2.2}$$

considering now the $\nu = 0$ component the final result is:

$$(\overline{\sigma}^0)^{\dot{\beta}\alpha} \left\{Q_\alpha, \overline{Q}_{\dot{\beta}}\right\} = 4P^0 = 4E. \tag{2.3}$$

Since $Q_\alpha Q_\alpha^\dagger + Q_\alpha^\dagger Q_\alpha$ is positive than $\left\langle vac\right| Q_\alpha Q_\alpha^\dagger + Q_\alpha^\dagger Q_\alpha \left|vac\right\rangle$ and $\left\langle vac\right| E \left|vac\right\rangle$ are positive, in this way one can see that in broken supersymmetry the energy is positive.
As the usual symmetry breaking provides a massless scalar field called the Goldstone boson, the supersymmetry breaking provides a massless field which is spinorial and it is called goldstino.
The effect of the supersymmetry breaking is to split the mass spectrum of bosons and fermions in the superfields. From a phenomenological point of view there are two possible ways of brake supersymmetry:

- soft-supersymmetry breaking: the masses of particles and super-partners are splitted but the couplings are kept equal, preserving part of the miraculous cancellation;
- hard-supersymmetry breaking: both masses and couplings are splitted removing the miraculous cancellation.

2.1.1 Global Supersymmetry Breaking

Supersymmetry is broken by he presence of a positive vacuum energy and as seen in Chapter 1 there are two possible scalar potentials in the supersymmetric Lagrangian, the D-term potential and the F-term potential. In this way there are two different mechanisms to brake global supersymmetry which can also be mixed in the same model.

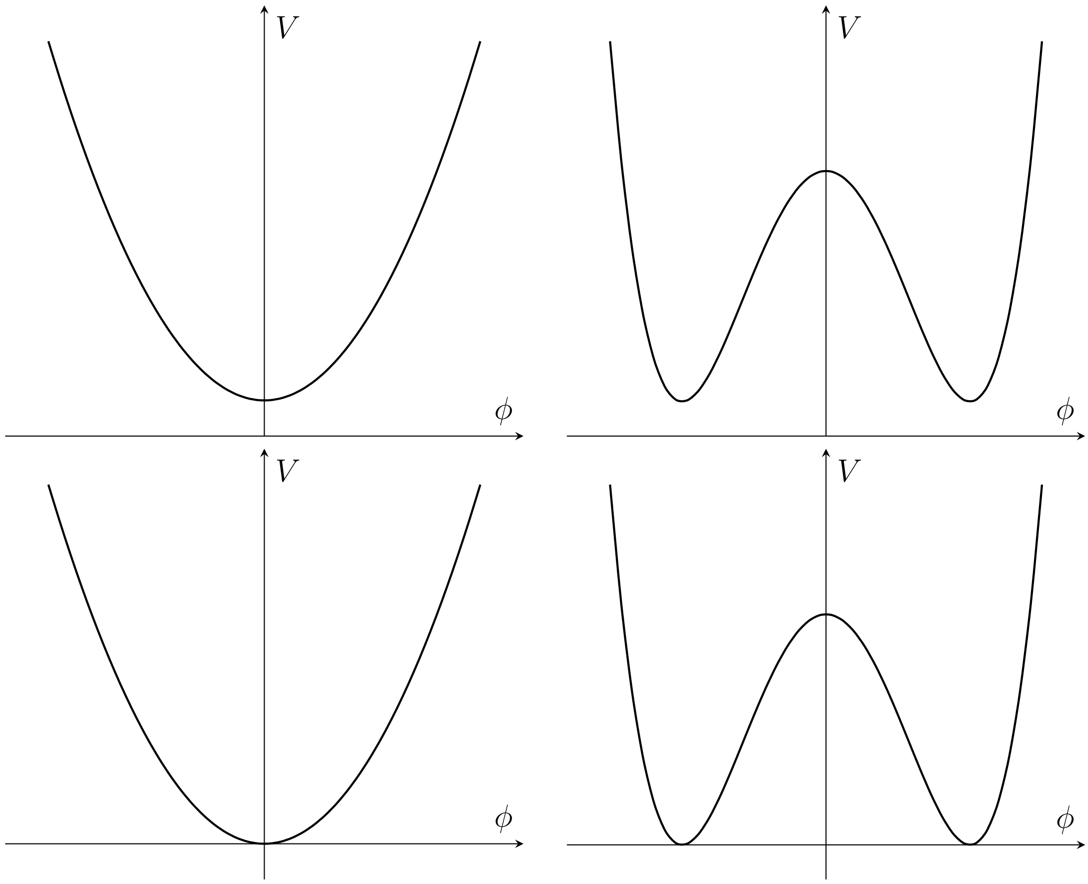

Figure 2.1: The possibles scalar potential for SUSY and GS breaking. (Upper left) SUSY breaking, GS preserving vacuum; (Upper right) SUSY breaking, GS breaking vacuum; (Down left) SUSY preserving, GS preserving vacuum; (Down right) SUSY preserving, GS breaking vacuum.

F-term Breaking

The F-term breaking can be studied starting from the supersymmetry transformation of the components of a chiral superfields Φ:

$$\begin{aligned} \delta\varphi &= \sqrt{2}\epsilon\psi \\ \delta\psi &= \sqrt{2}\epsilon F + i\sqrt{2}\sigma^\mu\bar{\epsilon}\partial_\mu\varphi \\ \delta F &= i\sqrt{2}\bar{\epsilon}\bar{\sigma}^\mu\partial_\mu\psi. \end{aligned} \tag{2.4}$$

Due to the Lorentz invariance the only field that can acquire a vacuum expectation value is the scalar auxiliary field F which makes ψ the Goldstone fermion. The resulting masses acquired by the particles provide a mass gap of the form:

$$\Delta m^2 \propto \langle F \rangle. \tag{2.5}$$

D-term Breaking

The D-term breaking can be studied starting from the supersymmetry transformation of the components of a vector superfields V:

$$\begin{aligned}\delta\lambda &= 2\bar{\epsilon}D + \frac{1}{2}\bar{\sigma}^{\nu}\sigma^{\mu}\bar{\epsilon}\partial_{\mu}A_{\nu} \\ \delta A_{\mu} &= \epsilon\sigma_{\mu}\bar{\lambda} + \lambda\sigma_{\mu}\bar{\epsilon} \\ \delta D &= \frac{i}{2}\left(\epsilon\sigma^{\mu}\partial_{\mu}\bar{\lambda} - \partial_{\mu}\lambda\sigma^{\mu}\bar{\epsilon}\right).\end{aligned} \tag{2.6}$$

Also in this case due to the Lorentz invariance the only field that can acquire a vacuum expectation value is the scalar auxiliary field D which makes λ the Goldstone fermion. The mass gap relation in this kind of models is:

$$\Delta m^2 \propto \langle D \rangle \tag{2.7}$$

Despite there are different ways to brake global supersymmetry at tree level there is a universal relation for the mass spectrum which is:

$$\mathrm{STr}(M^2) = \sum_j (-1)^j(2j+1)m_j^2 = 0. \tag{2.8}$$

The supertrace mass relation (2.8) shows the issue of the tree-level supersymmetry breaking: one partner is heavier and the other is lighter. This problem is phenomenological unacceptable because there should exist superparticles with lower mass. By the definition of the F-term $\overline{F}_i = -\partial_i W$ and by the non-renormalization theorem one gets that in order to avoid the supertrace mass relation (2.8) supersymmetry has to be broken by non-perturbative effects.

2.1.2 Supersymmetry Breaking in Supergravity

Supersymmetry can also be broken in local supersymmetry (supergravity). In this case, the Goldstone fermion is "eaten up" by the gravitino which develops a non-zero mass given by:

$$m_{3/2}^2 = e^K |W|^2. \tag{2.9}$$

In this case the supertrace relation is modified giving [21]:

$$\mathrm{STr}(M^2) = \sum_j (-1)^j(2j+1)m_j^2 = 2(N-1)m_{3/2}^2 \tag{2.10}$$

where N is the number of chiral superfields.
Due to the definition of the F-term in supergravity, there are two possible ways to provide a non-vanishing vacuum expectation value out of the tree level, using perturbative and non-perturbative corrections:

$$F^i = e^{K/2}K^{i\bar{j}}D_{\bar{j}}\overline{W} = e^{K/2}K^{i\bar{j}}\Bigg(\underbrace{\partial_{\bar{j}}\overline{W}}_{\text{Non-Perturbative}} + \underbrace{\overline{W}K_{\bar{j}}}_{\text{Perturbative}}\Bigg). \tag{2.11}$$

A particular feature of supergravity is the non-positive definite scalar potential which allows negative or zero cosmological constant, providing a way to obtain a very low cosmological constant to match cosmological observations.

2.1.3 Supersymmetry Breaking in the MSSM

Spontaneous supersymmetry breaking is not phenomenological acceptable because as shown by the equation (2.8) this would imply lighter partners which have not been observed, the way to solve this problem is to break supersymmetry in a hidden sector which is a singlet under the SM gauge group [1]. Supersymmetry breaking is than mediated to the visible sector by a different mechanism. The different kinds of mediation are:

- Gravity mediated: supersymmetry breaking is mediated by gravitational interactions, so the scale of supersymmetry breaking can be roughly estimated by the mass gap:

$$\Delta m = \frac{M_{SB}^2}{M_P}. \tag{2.12}$$

  Using the value of the Planck mass ($M_P \approx 10^{18}$GeV) and imposing the mass gap at the TeV scale $\Delta m \approx 1$TeV, the scale of the supersymmetry breaking is $M_{SB} \approx 10^{11}$GeV while the gravitino mass scale is $m_{3/2} \approx 1$TeV.

- Gauge mediated: supersymmetry breaking is mediated by a new gauge interaction under which matter fields are charged, in this case, the mass gap is proportional to the mass scale of supersymmetry breaking $\Delta m \approx M_{SB} \approx 1$TeV the gravitino mass is very low $m_{3/2} \approx 1$meV.

- Anomaly mediated: supersymmetry breaking is mediated by loops corrections. This kind of contribution is always present but generally it is negligible with respect to the others.

The three mediations are not necessarily unique but can be mixed giving different contributions.
Despite the presence of different mechanisms to mediate supersymmetry breaking, to perform a soft-breaking, the supersymmetry breaking part of the Lagrangian is model independent. The soft-breaking Lagrangian is defined:

$$-\mathcal{L}_{soft} = \frac{1}{2}(M_{1/2})_a \lambda^a \lambda^a + h.c. + (m_0)^2_{ij} \varphi_i^* \varphi_j + A_{ijk} \varphi_i \varphi_j \varphi_k + h.c. + B\mu H_u H_d + h.c. \tag{2.13}$$

The parameters are:

- $(M_{1/2})_a$ are the gaugino masses;

- $(m_0)_{ij}$ are the scalar particle masses;

- A_{ijk} are the trilinear terms;

- $B\mu$ is the mass term for the Higgs fields.

2.2 Flavour Changing Neutral Currents

Flavour is the way in which the three different families of quarks and leptons are labelled in the Standard Model:

$$\begin{aligned}
&3 \text{ up-type quarks } u, c, t;\\
&3 \text{ down-type quarks } d, s, b;\\
&3 \text{ type of charged leptons } e, \mu, \tau;\\
&3 \text{ type of neutrinos } \nu_e, \nu_\mu, \nu_\tau.
\end{aligned} \tag{2.14}$$

Flavour is an ambiguous quantity because the Standard Model without the Yukawa couplings would manifest a $U(3)$ flavour symmetry [22] and also in the complete Standard Model it is conserved for neutral process but is changed in charged boson exchange phenomenon. In the Standard Model this evidence arises from the breaking of the $SU(2)_L \times U(1)_Y$ symmetry, moving from the gauge eigenstates to the mass eigenstates the interacting Lagrangian for the Electroweak sector of quarks is [23]:

$$\begin{aligned}
J_Y^\mu =& \frac{1}{3}\overline{u}_L^i\gamma^\mu u_L^i + \frac{1}{3}\overline{d}_L^i\gamma^\mu d_L^i + \frac{4}{3}\overline{u}_R^i\gamma^\mu u_R^i - \frac{2}{3}\overline{d}_R^i\gamma^\mu d_R^i\\
J_3^\mu =& \frac{1}{2}\overline{u}_L^i\gamma^\mu u_L^i - \frac{1}{2}\overline{d}_L^i\gamma^\mu d_L^i\\
J_C^\mu =& 2\overline{u}_L^i\gamma^\mu V_{ij} d_L^j\\
\mathcal{L}_{Int} =& \left[\frac{1}{2}g_1 cos\theta_W J_Y^\mu + g_2 sin\theta_W J_3^\mu\right] A_\mu + \left[-\frac{1}{2}g_1 sin\theta_W J_Y^\mu + g_2 cos\theta_W J_3^\mu\right] Z_\mu +\\
&+ gW_\mu J_C^\mu + h.c..
\end{aligned} \tag{2.15}$$

The charged current J_C^μ in equation (2.15) respect to the currents J_Y^μ and J_3^μ, which couple to neutral vector bosons, presents a flavour changing behaviour parametrized by the matrix V which is the Cabibbo-Kobayashi-Maskawa matrix. The CKM matrix is in general not diagonal and produce the diagrams in Figure 2.2.

$$V = \begin{pmatrix} V_{ud} & V_{us} & V_{ub}\\ V_{cd} & V_{cs} & V_{cb}\\ V_{td} & V_{ts} & V_{tb} \end{pmatrix}. \tag{2.16}$$

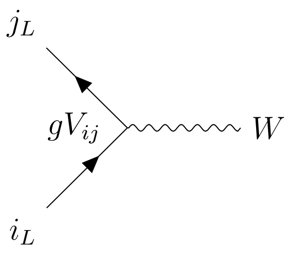

Figure 2.2: Flavour changing vertex with $i = u,\ c,\ t$ and $j = d,\ s,\ b$.

The FCNCs are so forbidden at the tree level by the definition of the interacting Lagrangian (2.15) but they could appear at the loop level. The CKM matrix is unitary and it is important in order to avoid the FCNCs at the loop level. This is explained by the so-called GIM mechanism.
For the leptons the FCNCs are completely forbidden by the absence of a flavour changing matrix by the fact that the neutrinos remain massless.

2.2.1 GIM-Mechanism

The GIM mechanism was introduced by Glashow, Iliopoulos and Maiani in the 1970 [24] to explain the experimental suppression of the decay of the K^0, at that days the discovered quarks were only 3: u, d, s. The dominant channel of the decay is given by the diagram in Figure 2.3.

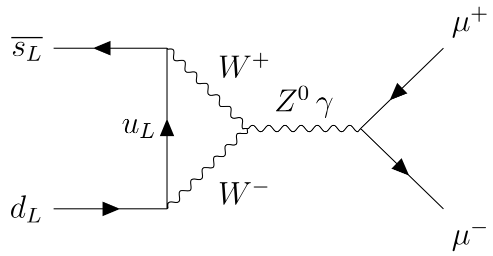

Figure 2.3: Dominant loop contribution to the decay of K^0 into $\mu^+\mu^-$.

The amplitude of this diagram is:

$$\Gamma^{diag}_{K^0\to\mu^+\mu^-} \propto g^2 V^*_{us} V_{ud} \tag{2.17}$$

The GIM mechanism explanation of the suppression is obtained introducing a new quark c analogous of the u but heavier, with CKM matrix unitary. By the introduction of the new charm quark c with unitary CKM matrix the new amplitude is:

$$\Gamma^{diag}_{K^0\to\mu^+\mu^-} \propto g^2 \left(V^*_{us} V_{ud} + V^*_{cs} V_{cd}\right) = 0. \tag{2.18}$$

Thanks to the introduction of the new quark the flavour changing neutral currents are suppressed also at loop level and the experimental measurements for K^0 decay are verified. This mechanism can be extended for 3 families of quarks as in the Standard Model.

2.2.2 FCNCs in the MSSM

In the MSSM the introduction of new particles which carry flavour has to be treated carefully because can introduce new possible FCNCs contributions, in this landscape is important to consider the possible deviation of the FCNCs results from the Standard Model ones.
The first possible source of flavour changing process is the, already mentioned in Chapter 1, barion and lepton number brake superpotential W_{BL} which allows FCNCs at the tree level:

$$W_{BL} = \lambda_1 LLe^c + \lambda_2 LQd^c + \lambda_2 LQd^c + \lambda_3 u^c d^c d^c + \mu' LH_u. \tag{2.19}$$

As already mentioned in Chapter 1 this term in the MSSM superpotential can be removed introducing a new discrete symmetry the R-parity which acts on a state as:

$$|\alpha\rangle \to |\alpha'\rangle = R\,|\alpha\rangle \qquad \text{with} \qquad R = (-1)^{3(B-L)+2S} \tag{2.20}$$

S is the spin, B is the baryon number and L is the lepton number. This symmetry implies that for each vertex of a diagram involving supersymmetric particles the number of superparticles must be even.
The R-parity forbids FCNCs at the tree level but they are not forbidden at the loop level. A complete analysis of the FCNCs in MSSM has been exposed in [25]. For the purpose of the thesis we will consider only the decay of the $K^0-\overline{K}^0$ for the quarks sector whose dominant term, in the Standard Model, is given by the diagram in Figure 2.4.

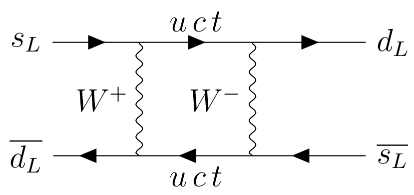

Figure 2.4: SM leading contribution to $K^0 - \overline{K}^0$ decay.

The amplitude dominant term is $(V_{ts}V_{td})^2\frac{m_t^2}{M_W^4}$ [25] which depends on the mass of the top because it is the heavier quark.
The squarks in the MSSM will produce a new contribution for the $K^0 - \overline{K}^0$ whose dominant diagram is traced in Figure 2.5.

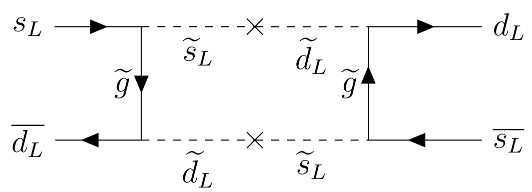

Figure 2.5: Supersymmetric contribution to $K^0 - \overline{K}^0$ decay.

The amplitude depends on $\Delta m^2_{\tilde{s}\tilde{d}}/m^2_{\tilde{q}}$ [26], where $m_{\tilde{q}}$ is the flavour universal mass term of the squarks and $\Delta m_{\tilde{s}\tilde{d}}$ is the non universal mass term between strange and down squarks. The SM model contribution well explains the experimental results so in order to preserve the SM description the flavour non universality of the masses should respect [26]:

$$\frac{\Delta m^2_{\tilde{s}\tilde{d}}}{m_0^2} \lesssim m_0 V_{ts} V_{td} \frac{1}{M_W} \sim 10^{-3}\left(\frac{m_0}{500GeV}\right). \tag{2.21}$$

So to avoid the FCNCs the scalar particles masses have to be flavour universal:

$$(m_0)_{ij} = m_0 \delta_{ij}. \tag{2.22}$$

The new supersymmetric particles allow the onset of FCNCs also in the leptonic sector while they are forbidden in the standard model. Let us consider as explicative example the decay of a μ^+ into a e^+ and a γ, this is given by the SM diagram in Figure 2.6 which is forbidden by the suppression of FCNCs in the leptonic sector. In the MSSM there is a new contribution given by the loop diagram (Figure 2.7) which involves the sleptons, also in this case the amplitudes is proportional to [27]: $\Delta m^2_{\tilde{\mu}\tilde{e}}/m^6_{\tilde{l}}$. So both quark and lepton sectors are affected by FCNCs in the MSSM and they need the flavour universality to respect the experimental results of the SM.

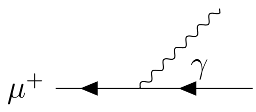

Figure 2.6: Standard Model contribution to MEG.

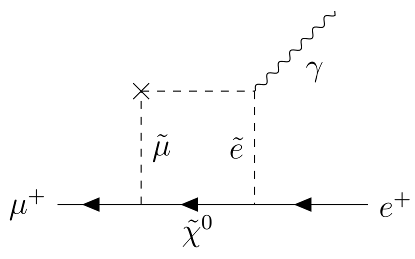

Figure 2.7: Supersymmetric contribution to MEG [28].

The non-flavour universal mass terms are not only generated by the supersymmetry breaking mechanism but can be generated by the Higgs mechanism trough the trilinear A_{ijk} term as can be seen in the soft-breaking Lagrangian [29]:

$$\begin{aligned} -\mathcal{L}_{soft} =& (m^2_Q)_{ij}\overline{Q}_{L_i}Q_{L_j} + (m^2_{u_R})_{ij}\overline{u}_{R_i}u_{R_j} + (m_{d_R})_{ij}\overline{d}_{R_i}d_{R_j}+ \\ &+ \left[A^u_{ij}\overline{Q}_{L_i}H^u u_{R_j} + A^d_{ij}Q_{L_i}H^d d_{R_j} + h.c.\right] + ... \end{aligned} \quad (2.23)$$

The acquiring of vacuum expectation value to the Higgs fields will give a non universal mass contribution to the squarks:

$$(\Delta m^2_{LR}{}^{(u)})_{ij} = A^d_{ij}\langle H^u_0\rangle, \qquad (\Delta m^2_{LR}{}^{(d)})_{ij} = A^d_{ij}\langle H^d_0\rangle, \quad (2.24)$$

the new off-diagonal mass terms will contribute to the FCNCs with the same bound of the one generated by the supersymmetry breaking. The same problem arises considering the coupling of sleptons with the Higgs field which has the same form.
In order to avoid also the contributions from the Higgs mechanism is necessary to impose that the trilinear terms are proportional to the Yukawa couplings:

$$A_{ijk} = A_0 Y_{ijk}. \quad (2.25)$$

The universality problem is strongly related to the mediation mechanism of the supersymmetry breaking in fact, it is differently solved by the previously described mediations.

2.2.3 Colour Charge Breaking vacua

The trilinear terms have an additional constrain from colour and charge breaking vacua. The introduction of scalar fields, which carries colour and charge, introduces the possibility of new vacuum that brake the colour and charge symmetry which can be competitive respect to the real colour and charge preserving vacuum.
Let us consider the off-diagonal trilinear scalar coupling defining $|e_L| = |\mu_R| = |H_0^d| \equiv a$ [29]:

$$\begin{aligned} V =& m_{e_L}^2 \left|e_L\right|^2 + m_{\mu_R}^2 \left|\mu_R\right|^2 + m_{H^d}^2 \left|H^d\right|^2 + \left|\lambda_e \bar{e}_L H_0^d\right|^2 + \left|\lambda_\mu H_0^d \mu_R\right|^2 - 2\left|A_{12}^{(l)}\right| \bar{e}_L H_0^d \mu_R = \\ =& \left(m_{e_L}^2 + m_{\mu_R}^2 + m_{H^d}^2\right) a^2 + \left(\lambda_e^2 + \lambda_\mu^2\right) a^4 - 2\left|A_{12}^{(l)}\right| a^3. \end{aligned} \tag{2.26}$$

One can easily see that the potential has a deep CCB minimum at $a \sim \left|A_{12}^{(l)}\right| / \lambda_\mu^2$. The minimum disappears for:

$$\left|A_{12}^{(l)}\right|^2 \leq \frac{8}{9}\left(\lambda_\mu^2 + \lambda_e^2\right)\left(m_{e_L}^2 + m_{\mu_R}^2 + m_{H^d}^2\right). \tag{2.27}$$

For diagonal trilinears the condition is similar $\left|A_{11}^{(l)}\right|^2 \leq 3\lambda_e^2\left(m_{e_L}^2 + m_{e_R}^2 + m_{H^d}^2\right)$, $\left|A_{22}^{(l)}\right|^2 \leq 3\lambda_\mu^2\left(m_{\mu_L}^2 + m_{\mu_R}^2 + m_{H^d}^2\right)$. Generalizing for the others particles:

$$\begin{aligned} \left|A_{ij}^{(u)}\right|^2 &\leq \lambda_{u_k}^2\left(m_{u_{Li}}^2 + m_{u_{Rj}}^2 + m_{H^u}^2\right), \quad k = \mathrm{Max}(i,j) \\ \left|A_{ij}^{(d)}\right|^2 &\leq \lambda_{d_k}^2\left(m_{d_{Li}}^2 + m_{d_{Rj}}^2 + m_{H^d}^2\right), \quad k = \mathrm{Max}(i,j) \\ \left|A_{ij}^{(l)}\right|^2 &\leq \lambda_{e_k}^2\left(m_{e_{Li}}^2 + m_{e_{Rj}}^2 + m_{H^d}^2\right), \quad k = \mathrm{Max}(i,j). \end{aligned} \tag{2.28}$$

These conditions avoid the formation of a deep CCB minimum which would make the vacuum unstable causing a decay in the deepest vacuum braking the colour and charge symmetry. The CCB conditions are more restrictive respect to the FCNCs ones as derived in [24].

2.2.4 Unbounded From Below potential

Another constrain on the trilinear terms can arise from the possibility of unbounded from below directions in the scalar potential, this can be seen considering the scalar fields $|e_L|^2 = |\mu_R|^2 = |H_0^d|^2 + |\nu_\tau|^2 \equiv a^2$ in this case the scalar potential is only [29]:

$$\begin{aligned} V =& a^2\left[m_{e_L}^2 + m_{\mu_R}^2 + m_{\nu_\tau}^2 - 2\left|A_{12}^{(l)}\right| H_0^d + \left|\lambda_\mu H_0^d\right|^2\right] + \\ &+ \left(m_{H^d}^2 - m_{\nu_\tau}^2\right)\left|H_0^d\right|^2. \end{aligned} \tag{2.29}$$

If one minimizes the potential with respect to the Higgs field gets:

$$H_0^d = \frac{\left|A_{12}^{(l)}\right| a^2}{\lambda_\mu^2 a^2 + \left(m_1^2 - m_{\nu_\tau}^2\right)}, \tag{2.30}$$

leading to the potential:

$$V = a^2 \left[m^2_{e_L} + m^2_{\mu_R} + m_{\nu_\tau} - \left| A^{(l)}_{12} \right|^2 \frac{a^2}{\lambda^2_\mu a^2 + \left(m^2_{H^d} - m^2_{\nu_\tau} \right)} \right]. \tag{2.31}$$

In order to get a positive potential one has to impose the condition:

$$\left| A^{(l)}_{12} \right|^2 \frac{a^2}{\lambda^2_\mu a^2 + \left(m^2_{H^d} - m^2_{\nu_\tau} \right)} \leq \left(m^2_{e_L} + m_{\mu_R} + m_{\nu_\tau} \right) \tag{2.32}$$

In the limit $a \gg \frac{(m^2_{H^d} - m^2_{\nu_\tau})}{\lambda^2_\mu}$ the bound becomes:

$$\left| A^{(l)}_{12} \right|^2 \leq \lambda^2_\mu \left(m^2_{e_L} + m^2_{\mu_R} + m^2_{\nu_\tau} \right). \tag{2.33}$$

This condition can be extended to the others trilinear terms leading to:

$$\begin{aligned} \left| A^{(u)}_{ij} \right|^2 &\leq \lambda^2_{u_k} \left(m^2_{u_{Li}} + m^2_{u_{Rj}} + m^2_{e_{Lp}} + m^2_{e_{Rq}} \right), & k =& \text{Max}(i,j),\ p \neq q \\ \left| A^{(d)}_{ij} \right|^2 &\leq \lambda^2_{d_k} \left(m^2_{d_{Li}} + m^2_{d_{Rj}} + m^2_{\nu_m} \right), & k =& \text{Max}(i,j) \\ \left| A^{(l)}_{ij} \right|^2 &\leq \lambda^2_{e_k} \left(m^2_{e_{Li}} + m^2_{e_{Rj}} + m_{\nu_m} \right), & k =& \text{Max}(i,j),\ m \neq i,j. \end{aligned} \tag{2.34}$$

The conditions are important to get a positive potential in the minimum. These conditions are at the leading order comparable to the ones for the CCB vacua.

2.2.5 Naturalness

Before seeing the ways to obtain the flavour universality in the different mediations of the supersymmetry breaking, it is important to analyse another way to suppress the FCNCs: to set the mass of the scalar particles at very high energy scale (respect to the TeV scale). By the fact that the FCNCs amplitudes depend on $1/m_0^2$, the high value of the mass would make negligible the supersymmetric contributions.
This would give problems with naturalness. Naturalness is a criterion on the parameters of a theory that has been well expressed by Susskind as a criterion "which requires the observable properties of a theory to be stable against minute variations of the fundamental parameters" [30]. In fact one can consider the dependence of M_z respect to the parameters at the unification scale which is [27]:

$$\begin{aligned} M_z^2 =& c_\mu \mu^2 + c_{H_d} m^2_{H_d0} + c_{H_u} m^2_{H_u0} + c_t m^2_{t0} + c_{\tilde{t}} m^2_{\tilde{t}0} + \\ &+ C_{M_{1/2}} M^2_{1/20} + c_{AM} A_{t0} M_{1/20} + c_A A^2_{t0}, \end{aligned} \tag{2.35}$$

where the masses are the masses of the particles of the MSSM at the gran unification scale (GUT) and the coefficients c are constants, of the order of the unity, results of the RG equations. One defines the fine tuning amount f_i as:

$$f_i = \frac{M_z^2}{c_i i^2} \tag{2.36}$$

with i the possible parameters in the equation (2.35). So if one sets the mass of the stop $m_{\tilde{t}0}$ to a very high value respect to the TeV scale, for example considering the hill case of 100% of non-universality, one gets that the mass of the squarks should be at the order $m_0 \gtrsim 5\ 10^5$ GeV. One can compute the amount of fine-tuning needed to get the value of $M_z \approx 91$ GeV getting $f_{m_0} \lesssim 3\ 10^{-8}$ with a very huge amount of fine-tuning out of the 10% naturalness criterion [31].
This kind of solution would spoil naturalness and even if naturalness is a discussed criterion from an epistemological point of view [32]. It is also one of the reasons to introduce supersymmetry at low energy to stabilise the Higgs mass as seen in Chapter 1.

2.2.6 Gauge-mediated solution

The gauge-mediated supersymmetry is based on the breaking of supersymmetry in a hidden sector and then it is mediated to the visible sector thanks to new superfields which are charged under the MSSM gauge group.
There are many ways to brake the supersymmetry in the hidden sector but the easiest possible gauge-mediated model [33] is made by introducing a new goldstino chiral superfield X which acquires a vacuum expectation value in the scalar M and F auxiliary components:

$$\langle X \rangle = \langle M \rangle + \theta\theta \langle F \rangle. \tag{2.37}$$

The messenger sector is made by N_f flavours of chiral superfields Φ_i and $\overline{\Phi}_i$ which are charged under the gauge group of the MSSM. The superpotential which couples X and Φ is:

$$W = \lambda_{ij} \overline{\Phi}_i X \Phi_j. \tag{2.38}$$

By the acquisition of vacuum expectation value of X, thanks to the direct coupling, Φ and $\overline{\Phi}$ get masses. The spinorial components form a Dirac field with mass λM while the scalar components get mass matrix:

$$m_{\Phi\overline{\Phi}} = \begin{pmatrix} (\lambda M)^\dagger (\lambda M) & (\lambda F)^\dagger \\ (\lambda F) & (\lambda M)(\lambda M)^\dagger \end{pmatrix}. \tag{2.39}$$

Going in the mass eigenstates $(\Phi + \overline{\Phi}^\dagger)/\sqrt{2}$ and $(\overline{\Phi} - \Phi^\dagger)/\sqrt{2}$ the mass eigenvalues are $(\lambda M)^2 \pm (\lambda F)$.
The ordinary particles do not directly acquire masses, by the fact that they do not couple with X, but, by the gauge interactions with Φs, they acquire masses by loop corrections. Thanks the non renormalization of the superpotential all the relevant contributions to the soft-terms are in the gauge and matter wave-function normalizations S and Z_i. The soft Lagrangian will be:

$$-\mathcal{L}_{soft} = \frac{1}{2} M_{1/2} \lambda_a \lambda_a + h.c. + m_i^2 \varphi_i^\dagger \varphi_i + A_i \varphi_i \partial_{\varphi_i} W(\varphi) + h.c. \tag{2.40}$$

with the parameters [33]:

$$M_{1/2} = -\frac{1}{2} \frac{\partial \ln S(X,t)}{\partial \ln X}\bigg|_{X=M} \frac{F}{M}, \tag{2.41}$$

$$m_i^2 = -\left.\frac{\partial^2 \ln Z_i(X, X^\dagger, t)}{\partial \ln X \partial \ln X^\dagger}\right|_{X=M} \frac{FF^\dagger}{MM^\dagger}, \tag{2.42}$$

$$A_i(t) = \left.\frac{\partial \ln Z_i(X, X^\dagger, t)}{\partial \ln X}\right|_{X=M} \frac{F}{M}, \tag{2.43}$$

with $t = \ln M^2/Q^2$ and Q is the definition scale of the soft-terms. The gaugino masses are generated by the loop in Figure 2.8 while the scalar masses are generated by the loop in Figure 2.9.

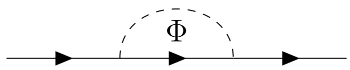

Figure 2.8: Generation of gaugino masses at loop level.

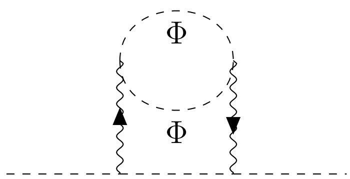

Figure 2.9: Generation of scalar masses at 2-loop level.

As shown by the definitions the scalar masses are flavour universal thanks to the symmetries of gauge interaction. Also the trilinear terms respect the flavour universality because, from the Lagrangian (2.40), they depend on the Yukawa couplings (by the derivative of the superpotential). In this mediation the FCNCs are naturally avoided.
A possible source of non-universality is given by Planckian effects. If supersymmetry is broken by chiral field ($X = \theta^2 F_0$) one can not avoid in the Lagrangian the operator [26]:

$$\Delta\mathcal{L} \sim \int d^4\theta \frac{1}{M_P^2} X^\dagger X \Psi_i^\dagger \Psi_i \tag{2.44}$$

which gives a flavour non-universal mass term:

$$\Delta m_i^2 \sim \frac{F_0^2}{M_P^2} \tag{2.45}$$

the experimental bound for F_0 to avoid FCNCs is:

$$\sqrt{F_0} \lesssim 10^{10} \left(\frac{m_i^2}{500\text{GeV}}\right)^3 \text{GeV}. \tag{2.46}$$

The FCNCs are strongly suppressed in the gauge mediated models by the fact that the pure gauge contribution generates flavour universal terms and the spoil term arise from gravitational contribution (which is suppressed by the Planck mass).

2.2.7 Anomaly-mediated solution

The anomaly-mediated supersymmetry breaking is based on the breaking of the super-Weyl invariance in the supergravity by the loop corrections of a super Yang-Mills theory. Let us consider a super Yang-Mills theory coupled to supergravity. Its Lagrangian is:

$$\begin{aligned}\mathcal{L} =&\sqrt{-g}\{\int d^4\theta h(Q^\dagger, e^{-V}Q)\Phi^\dagger\Phi + \int d^2\theta(\Phi^3 W(Q) + \tau(Q)W_\alpha W^\alpha) + h.c.+ \\ &- \frac{1}{6}h(\tilde{q}^\dagger, \tilde{q})(R + ...)\}\end{aligned} \tag{2.47}$$

with Q the chiral superfield in the visible sector, $\tilde{q}$ the scalar component, $h = -3e^{-K/3}$ and Φ is the auxiliary scalar field $\Phi := 1 + \theta^2 F_\Phi$ from the off-shell supergravity multiplet. The Lagrangian (2.47) is equivalent to the usual supergravity Lagrangian mentioned in Chapter 1, one can see this by performing the Weyl transformation which removes Φ:

$$g_{\mu\nu} \to e^{K/3} g_{\mu\nu}. \tag{2.48}$$

The Lagrangian is given by:

$$\begin{aligned}\mathcal{L} =&\sqrt{-g}\{\frac{1}{2}R + K_{ij}(\tilde{q}^\dagger\tilde{q})D_\mu\tilde{q}^{i\dagger}D^\mu\tilde{q}^j - V(\tilde{q}^\dagger\tilde{q})+ \\ &- \tau(\tilde{q})(F_{\mu\nu}F^{\mu\nu} + iF_{\mu\nu}\tilde{F}^{\mu\nu}) + h.c + \text{fermionic terms}\}\end{aligned} \tag{2.49}$$

with $K_{ij} = \partial_{\tilde{q}^i}\partial_{\tilde{q}^j}K$ and $V(\tilde{q}^\dagger\tilde{q})$ the supergravity scalar potential.
Φ can acquire supersymmetry breaking expectation value but Lorentz preserving. The fact that the matter field of the visible sector does not couple with Φ correspond to a super-Weyl invariance. If the super-Weyl symmetry was exact than the acquiring of VEV of Φ would not affect the visible sector because it would have been removed by a Weyl transformation. In the presence of a super Yang-Mills sector the super-Weyl transformation is anomalous producing the shift of the τ gauge kinetic function [35]:

$$\tau \to \tau - 2b_0 \ln(\Phi) \tag{2.50}$$

with b_0 the one loop coefficient. Solving the integral in the Lagrangian one gets a gaugino mass which is:

$$M_{1/2} = -b_0 g^2 F_\Phi. \tag{2.51}$$

The scalar mass can be understood considering the Lagrangian and integrating out the hidden sector:

$$\mathcal{L}_{eff} = \int d^4\theta Q^\dagger e^{-V} Q\Phi^\dagger\Phi + \int d^2\theta \left[\Phi^3\left(m_0 Q^2 + Y_0 Q^3\right) + \frac{1}{g_0^2}W_\alpha W^\alpha\right] + h.c. \tag{2.52}$$

The Φ field can be removed rescaling the Q field:

$$Q\Phi \to Q. \tag{2.53}$$

Q is a chiral superfield and Φ is a background chiral superfield so the rescaling is allowed by the supersymmetry. The resulting Lagrangian is:

$$\mathcal{L}_{eff} = \int d^4\theta Q^\dagger e^{-V} Q + \int d^2\theta \left[\left(m_0^2 + Y_0 Q^3\right) + \frac{1}{g_0^2}W_\alpha W^\alpha\right] + h.c. \tag{2.54}$$

In the classical field theory Φ can be removed by the rescaling (it appears only in the quadratic term but also the μ-term is allowed by supersymmetry) removing also the supersymmetry breaking term but the quantum functional integral measure is not invariant under the rescaling providing a mass term from the loop corrections. The Φ dependence provide a scalar mass term at the second loop correction of the form:

$$m_0^2(\mu) = -\frac{1}{2}|F_\Phi|^2 \left(\frac{\partial\gamma}{\partial g}\beta_g + \frac{\partial\gamma}{\partial y}\beta_y \right), \tag{2.55}$$

and it is flavour universal.
Due to the same properties of the renormalization and to the holomorphy of the superpotential, the trilinear terms are proportional to the Yukawa couplings:

$$A_{ikj} = \frac{1}{2}(\gamma_i + \gamma_j + \gamma_k) Y_{ijk} F_\Phi \tag{2.56}$$

β and γ are the beta and gamma functions of the renormalization group.
Also in this case the flavour universality problem is naturally solved by the mediation mechanism thanks to the universal nature of the loops corrections. The problem of the anomaly mediated supersymmetry breaking is the mass spectrum at the TeV scale [35]:

$$\begin{aligned} m^2_{sleptons} &= -1.3 \times 10^{-5} |F_\phi|^2 \\ m^2_{squarks} &= 5.5 \times 10^{-4} |F_\phi|^2. \end{aligned} \tag{2.57}$$

In this mediation scenario, the flavour supersymmetric problem is automatically solved but the resulting spectrum has negative mass for sleptons so one need other mechanisms contributions to avoid tachyonic solutions.

2.2.8 Gravity-mediated solution

In gravity mediated supersymmetry breaking the hidden sector is made by some chiral superfields Φ_i, Φ_is couple with the visible sector by the matter metric. The Φ_i fields acquire a non-vanishing F-term VEVs that break supersymmetry. Generally it happens thanks to non-perturbative effects in the superpotential. A general model is described by:

$$\begin{aligned} K &= \hat{K}(\Phi_i, \overline{\Phi}_i) + \tilde{K}_{\alpha\overline{\beta}}(\Phi_i, \overline{\Phi}_i) C^\alpha \overline{C}^\beta + ..., \\ W &= \hat{W}(\Phi_i) + \mu_{\alpha\overline{\beta}}(\Phi_i) C^\alpha \overline{C}^\beta + Y_{\alpha\beta\gamma} C^\alpha C^\beta C^\gamma + ..., \\ f &= f(\Phi_i). \end{aligned} \tag{2.58}$$

An example for the hidden sector is the no-scale supergravity model where the hidden sector is made by two chiral fields T and S and the dynamic is described by:

$$\begin{aligned} \hat{K}(S,T) &= -\ln\left(S + \bar{S}\right) - 3\ln(T + \bar{T}) \\ \hat{W}(S) &= W_0 + Ae^{-aS}. \end{aligned} \tag{2.59}$$

And the matter metric is: $\bar{K}_{\alpha\bar{\beta}} = \frac{3}{T+\bar{T}}$. S is stabilised supersymmetrically such that $F^S = 0$ but T no, giving $\bar{F}^{\bar{T}} = -\frac{W}{\sqrt{(S+\hat{S})(T+\hat{T})}}$.

After the supersymmetry breaking expanding the F-term scalar potential for the C^α fields one gets the soft breaking terms [36]:

$$\begin{aligned}\tilde{m}_{\alpha\overline{\beta}} &= (m_{3/2}^2 + V_0)\tilde{K}_{\alpha\overline{\beta}} - \overline{F}^{\overline{m}} F^n \left(\partial_{\overline{m}}\partial_n \tilde{K}_{\alpha\overline{\beta}} - (\partial_{\overline{m}}\tilde{K}_{\alpha\overline{\gamma}})\tilde{K}^{\overline{\gamma}\delta}(\partial_n \tilde{K}_{\delta\overline{\beta}})\right) \\ \tilde{A}_{\alpha\beta\gamma} &= e^{\hat{K}/2} F^m \left[\hat{K}_m Y_{\alpha\beta\gamma} - \left((\partial_m \tilde{K}_{\alpha\overline{\rho}})\tilde{K}^{\overline{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma)\right)\right].\end{aligned} \tag{2.60}$$

There could also be D-term contributions if the Φ_is are charged under some gauge field charge, expanding the D-term potential the resulting contributions are [37]:

$$\begin{aligned}m_{\alpha\overline{\beta}}^2 &= \sum_i \left[-\frac{1}{2} g_i^2 D_i^2 K_{\alpha\overline{\beta}} + g_i^2 D_i \partial_\alpha \partial_{\overline{\beta}} D_i\right] \\ A_{\alpha\beta\gamma} &= \sum_i \left[-\frac{1}{2} g_i^2 D_i^2 D_\alpha D_\beta \left(K_\gamma + \frac{W_\gamma}{|W|}\right) + g_i^2 D_i D_\alpha D_\beta D_\gamma D_i\right],\end{aligned} \tag{2.61}$$

with D_i the D-term associated to the Φ_i field.
From equations (2.59) and (2.60) one can see that the flavour universality is not automatically respected, this problem has lead to the necessity of a flavour analysis in order to build phenomenological relevant models. To get the flavour universality the supersymmetry breaking and the flavour physics should be decoupled.

2.3 Mirror Mediation

In the study [38] Joseph Conlon had derived the conditions that a supersymmetry breaking gravity mediated model has to respect in order to respect the flavour universality.

1 The hidden sector fields have to be factorized into two classes of fields Ψ_i and χ_j.

2 The Kähler potential should be a direct sum of terms depending only on the real part of Ψ_i and χ_j

$$K(\Psi + \overline{\Psi}, \chi, \overline{\chi}) = K_1(\Psi + \overline{\Psi}) + K_2(\chi, \overline{\chi}). \tag{2.62}$$

The reality condition is necessary to avoid relative phases between different A-terms.

3 The Yukawa couplings should depend only on the χ fields and the gauge kinetic functions should depend linearly only on the Ψ fields.

$$\begin{aligned}Y_{\alpha\beta\gamma}(\Psi, \chi) &= Y_{\alpha\beta\gamma}(\chi) \\ f_a(\Psi, \chi) &= \sum_i \lambda_i \Psi_i\end{aligned} \tag{2.63}$$

The linearity allows to have universal gaugino mass phases aligned with those of the A-terms.

4 The matter metric should factorise in terms of the real part of Ψ_i and χ_j

$$K_{\alpha\overline{\beta}}(\Psi, \overline{\Psi}, \chi, \overline{\chi}) = h(\Psi + \overline{\Psi}) k_{\alpha\overline{\beta}}(\chi, \overline{\chi}). \tag{2.64}$$

5 The supersymmetry breaking should be induced by the Ψ fields while the χ fields are stabilised supersymmetrically:

$$D_\Psi W \neq 0, D_\chi W = 0. \tag{2.65}$$

These conditions are important because in this way there are generated two decoupled sectors: Φ which is responsible for the supersymmetry breaking and χ which is related to the flavour physics. One can sees that these conditions generate flavour universal soft-terms looking at the unnormalized scalar masses and the unnormalized trilinears:

$$\begin{aligned}\tilde{m}_{\alpha\overline{\beta}} =&(m_{3/2}^2 + V_0)K_{\alpha\overline{\beta}} - \overline{F}^{\overline{\Psi}_j}F^{\Psi_i}\left(\partial_{\overline{\Psi}_j}\partial_{\Psi_i}K_{\alpha\overline{\beta}} - (\partial_{\overline{\Psi}_j}K_{\alpha\overline{\gamma}})K^{\overline{\gamma}\delta}(\partial_{\Psi_i}K_{\delta\overline{\beta}})\right) = \\ =&\left(((m_{3/2}^2 + V_0)h - \overline{F}^{\overline{\Psi}_j}F^{\Psi_i}\left(\partial_{\overline{\Psi}_j}\partial_{\Psi_i}h - \frac{\partial_{\overline{\Psi}_j}h\partial_{\Psi_i}h}{h}\right)\right)k_{\alpha\overline{\beta}}(\chi,\overline{\chi})\end{aligned} \tag{2.66}$$

$$A_{\alpha\beta\gamma} = e^{K/2}Y_{\alpha\beta\gamma}(\chi)\left(F^\Psi\partial_\Psi K(\Psi,\overline{\Psi}) - 3\frac{F^\Psi\partial_\Psi h(\Psi,\overline{\Psi})}{h(\Psi,\overline{\Psi})}\right). \tag{2.67}$$

These conditions are ad hoc conditions in order to obtain flavour universal masses. There is not a deeper meaning and this is not theoretical satisfying but there is a more fundamental theory: String Theory that naturally respects the Mirror Mediation conditions.

2.3.1 Mirror Mediation in String Theory

If now one extends the analysis of Mirror Mediation to the Type IIB string compactification, as made in [38], one can see that the conditions are respected naturally solving the flavour supersymmetric problem in gravity mediated supersymmetry breaking from String Theory.

Moduli

Let us start from the moduli sector. In Type IIB compactification there are 3 types of moduli: S the dilaton, U the complex structure moduli (that can be many) and T the Kähler moduli (that can be many) as seen in Chapter 1. The moduli are independent in the Kähler potential for O3/O7 set-up in the large volume regime:

$$K = -\log\left(i\int\Omega_3\wedge\overline{\Omega}_3\right) - \log\left(S+\overline{S}\right) - 2\log\left(\mathcal{V}\right). \tag{2.68}$$

The moduli form two different and independent classes as can be seen by the Kähler potential (2.68):

$$\begin{aligned}\Psi_i &\longleftrightarrow T_i \\ \chi_i &\longleftrightarrow U_i, S.\end{aligned}$$

One possible mixing term is the α' correction from the 10D action which will modify the Kähler moduli potential mixing the dilaton and the Kähler moduli dependence:

$$K_{KM} = -2\log\left[\mathcal{V} + \xi\mathrm{Re}(S)^{3/2}\right]. \tag{2.69}$$

This kind of mixing term induced by the perturbative correction spoil the Mirror Mediation condition but it could be reconstructed, at least, at the leading order.
The mixed Kähler metric will be:

$$
\begin{aligned}
K_{i\bar{j}} &= \begin{pmatrix} \frac{3}{8}\frac{3\tau_b-\tau_b^{-1/2}\mathcal{Y}}{\mathcal{Y}^2} & \frac{9}{8}\frac{(\tau_b\tau_s)^{1/2}}{\mathcal{Y}^2} & -\frac{9}{4}\frac{\xi(2s)^{1/2}\tau_b^{1/2}}{\mathcal{Y}^2} \\ \frac{9}{8}\frac{(\tau_b\tau_s)^{1/2}}{\mathcal{Y}^2} & \frac{3}{8}\frac{3\tau_s+\tau_s^{-1/2}\mathcal{Y}}{\mathcal{Y}^2} & \frac{9}{4}\frac{\xi(2s)^{1/2}\tau_s^{1/2}}{\mathcal{Y}^2} \\ -\frac{9}{4}\frac{\xi(2s)^{1/2}\tau_b^{1/2}}{\mathcal{Y}^2} & \frac{9}{4}\frac{\xi(2s)^{1/2}\tau_s^{1/2}}{\mathcal{Y}^2} & \frac{1}{(2s)^2}\left[1+\frac{3}{2}\xi(2s)^{3/2}\frac{2\xi(2s)^{3/2}-\mathcal{V}}{\mathcal{Y}^2}\right] \end{pmatrix} \sim \\
&\sim \begin{pmatrix} \mathcal{O}\left(\frac{1}{\mathcal{V}^{4/3}}\right) & \mathcal{O}\left(\frac{1}{\mathcal{V}^{5/3}}\right) & \mathcal{O}\left(\frac{1}{\mathcal{V}^{5/3}}\right) \\ \mathcal{O}\left(\frac{1}{\mathcal{V}^{5/3}}\right) & \mathcal{O}\left(\frac{1}{\mathcal{V}}\right) & \mathcal{O}\left(\frac{1}{\mathcal{V}^{2}}\right) \\ \mathcal{O}\left(\frac{1}{\mathcal{V}^{5/3}}\right) & \mathcal{O}\left(\frac{1}{\mathcal{V}^{2}}\right) & \mathcal{O}\left(1\right) \end{pmatrix}.
\end{aligned}
\tag{2.70}
$$

In the Kähler metric it is possible to see that the mixing terms of order $\mathcal{O}\left(\frac{1}{\mathcal{V}^{5/3}}\right)$ and $\mathcal{O}\left(\frac{1}{\mathcal{V}^{2}}\right)$ can be neglected respect to the diagonal terms in the large volume limit. So it is possible to see that LVS models, which stabilise the Kähler moduli through α' corrections, are compatible with the Mirror Mediation conditions.

Yukawa Couplings

In type IIB compactification the imaginary part of the Kähler moduli enjoys a shift symmetry:

$$\mathrm{Im}(T) \to \mathrm{Im}(T) + a \tag{2.71}$$

so only the real part of the Kähler moduli can enter into the action so due to the holomorphy of the superpotential, the superpotential could not depend on the Kähler moduli. Then the Yukawa couplings should depends only on the dilaton and the complex structure moduli as already demonstrated in the Chapter 1:

$$Y_{\alpha\beta\gamma} = Y_{\alpha\beta\gamma}(S, U). \tag{2.72}$$

The Mirror Mediation condition is then respected by the Yukawa couplings.

Gauge Kinetic Function

As we have seen in Chapter 1 the gauge kinetic function comes from the dimensional reduction of the DBI-CS action of the Dp-brane. It depends on what kind of brane we are looking at, in the analysis of this thesis is useful to see what happens for D3-brane (the one entering in the model we will look at in the Chapter 3) and D7-brane (the one mainly used for phenomenological models). For D7-brane the result is:

$$2\pi f_i^{D7} = n_i^j n_i^k S - m_i^j m_i^k T_i, \tag{2.73}$$

as we can see if the S dependence is negligible the only dependence of the gauge kinetic function is the Kähler moduli respecting the Mirror Mediation condition. Now we look at the D3-brane whose gauge kinetic function is:

$$2\pi f_{D3} = S, \tag{2.74}$$

as we can see this result brakes the Mirror Mediation condition and in fact, it will be a problem in the D3-brane model that we have analysed.

Factorisation of the Matter Metric

As already mentioned in Chapter 1 the matter metric is very hard to compute for general models but it has been estimated for Dp-branes at singularities in the large volume regime. At the leading order, it can be factorized as:

$$\tilde{K}_\alpha = \frac{h_\alpha(U,S)}{\mathcal{V}^{2/3}}, \tag{2.75}$$

where $h_\alpha(U,S)$ encodes the dependence on the other moduli and it is flavour dependent. In this way we can see that the Mirror Mediation condition on the matter metric is respected for general models, it can be studied also for a more particular models where $h_\alpha(U,S)$ is computable for example in toroidal models.

Supersymmetry Breaking

The last condition of the Mirror Mediation is focussed on supersymmetry breaking. Supersymmetry breaking in Type IIB string flux compactification arises from the GKP model:

$$\begin{aligned} K(T,S,U) &= -3\log(T+\bar{T}) - \log(S+\bar{S}) - \log\left(\int_{\mathcal{M}} \Omega_3 \wedge \bar{\Omega}_3\right) \\ W &= \int_{\mathcal{M}} \Omega_3 \wedge G_{(3)}. \end{aligned} \tag{2.76}$$

The superpotential is induced by the fluxes and the superpotential does not depends on the Kähler moduli at the tree level. From the flux constraints in the 10D theory the moduli are fixed as:

$$\begin{aligned} D_{U_\alpha} W &= \int_{\mathcal{M}} G_{(3)} \wedge \chi_\alpha = 0, \\ D_S W &= \frac{1}{S+\bar{S}} \int_{\mathcal{M}} \overline{G}_{(3)} \wedge \Omega_3 = 0, \\ D_T W &= -\frac{3W}{T+\bar{T}} \neq 0. \end{aligned} \tag{2.77}$$

As one can see, the model respects the Mirror Mediation condition obtaining that supersymmetry is broken by the Kähler moduli and the flavour-related moduli are stabilised supersymmetrically. The presence of non-perturbative corrections in the superpotential to stabilise the Kähler moduli will not break the result but α' correction could.
Following Conlon's work it is possible to see that Type IIB String Theory compactification will automatically respect the Mirror Mediation conditions producing flavour universal soft-terms.

Chapter 3

Sequestered String Models and FCNC Constraints

3.1 Sequestered Supersymmetry Breaking

As mentioned in the Section 1.3.3 the masses of the moduli are generally proportional to the mass of the gravitino and the lower bound for the mass of the moduli from the cosmological moduli problem is of order 50 TeV. With these premises, one would like to have the soft masses below the gravitino and moduli masses, the models which provide this different hierarchy are called sequestered models. In the LVS the sequestering happens when the SM fields are localized in the extra dimensions, like in models where the matter fields are obtained with D3-branes at a singularity [39]. The sequestered hierarchy is then produced by the weak coupling of the dominant F-terms and the matter fields because of their bulk separation. This is the case of the model [39] analysed in this thesis where in addition to the sequestering also the uplifting to the dS vacuum is achieved.

3.2 The Model

The model analysed is the one proposed in [39] which is the result of a Type IIB compactification with O3/O7-planes, the moduli are stabilized following the large volume scenario producing a visible sector sequestered from supersymmetry breaking, the visible sector is realised with proper D-brane configurations on blow-up moduli. The uplifting to a dS vacuum is realised through E(-1) instantons at a singularity whose blow-up mode develops non-vanishing F-term thanks to the dilaton-dependent non-perturbative effect. The Kähler moduli are:

- a big four-cycle T_b which controls the dimension of the Calabi-Yau,

- a small blow-up mode T_s supporting non-perturbative effects,

- the visible sector cycle T_{SM},

- the orientifold projection of T_{SM} G,
- the blow-up mode T_{dS} which supports E(-1) instantons non-perturbative effects.

The Kähler moduli are then decomposed into the real and the imaginary parts:

$$T_b = \tau_b + i\psi_b,\ T_s = \tau_s + i\psi_s,\ T_{SM} = \tau_{SM} + i\psi_{SM},\ G = b + ic,\ T_{dS} = \tau_{dS} + i\psi_{dS},\ \ (3.1)$$

where τ_b, τ_s, τ_{SM} and τ_{dS} are divisor volumes, the ψ's are axions obtained by the reduction of C_4 on the four-cycles, b and c are reduction of respectively B_2 and C_2 on the two-cycles dual to the shrinking ones.
The $\mathcal{N} = 1$ supergravity theory is described by the superpotential:

$$W = W_{flux}(U, S) + A_s(U, S)e^{-a_sT_s} + A_{dS}(U, S)e^{-a_{dS}(S+\kappa_{dS}T_{dS})} + W_{matter}. \qquad (3.2)$$

The first part of the superpotential W_{flux} is the standard superpotential induced by the fluxes that we have mentioned in the Chapter 1. The second term is the one induced by ED3-instantons on the small blow-up cycle (or by gaugino condensation) $A_s(U, S)$ depends on the complex structure moduli U and the dilaton S whose real part set the string coupling $\langle s \rangle = g_s^{-1}$, a_s depends on the D-brane configuration. The third term is the one associated to the uplifting moduli and it is generated by non-perturbative effects. The last term is the visible sector matter superpotential:

$$W_{matter} = \mu(\Phi)H_uH_d + \frac{1}{6}Y_{\alpha\beta\gamma}(\Phi)C^\alpha C^\beta C^\gamma + ... \qquad (3.3)$$

where Φs are the moduli and C^αs are the MSSM superfields. The dots refer to higher dimensional operators which are neglected.
The Kähler potential is:

$$K = -2\ln\left(\mathcal{V} + \frac{\hat{\xi}}{2}\right) - \ln(2s) + \lambda_{SM}\frac{\tau_{SM}^2}{\mathcal{V}} + \lambda_b\frac{b^2}{\mathcal{V}} + K_{cs}(U) + K_{matter} + \lambda_{dS}\frac{\tau_{dS}^2}{\mathcal{V}}. \quad (3.4)$$

The first term is the usual large volume scenario Kähler potential for the Kähler moduli with the α' correction. The second term is the Kähler potential for the dilaton. The λs are $\mathcal{O}(1)$ coefficients. $K_{cs}(U, S)$ is the tree level complex structure Kähler potential. The last term is the Kähler potential for the uplifting Kähler moduli. The matter Kähler potential is:

$$K_{matter} = \tilde{K}_\alpha(\Phi, \overline{\Phi})\overline{C}^{\overline{\alpha}}C^\alpha + \left[Z(\Phi, \overline{\Phi})H_uH_d + h.c.\right] \qquad (3.5)$$

The matter metric is assumed to be flavour diagonal with the only exception of the Higgs bilinear term which is parametrized by the function Z.

$$\tilde{K}_\alpha = \frac{f_\alpha(U, S)}{\mathcal{V}^{2/3}}\left(1 - c_s\frac{\hat{\xi}}{\mathcal{V}} + c_{dS}\frac{\tau_{dS}^2}{\mathcal{V}} + c_{SM}\tau_{SM}^p + c_b b^p\right). \qquad (3.6)$$

p is positive in order to have the matter metric well defined in the singular limit $b, \tau_{SM} \to 0$. The $f_\alpha(U, S)$ is the factor which depends on the dilaton and the complex structure moduli and will be described later in the work.

The gauge kinetic function is obtained by D3-branes at singularities, so it is defined to be:

$$f_a = \delta_a S + \kappa_a T_{SM} \tag{3.7}$$

where δ_a are universal constants for $\mathbb{Z}_n$ singularities but can be non universal for more general types of singularities, we can already see that the dilaton dependence already break the Mirror Mediation condition.

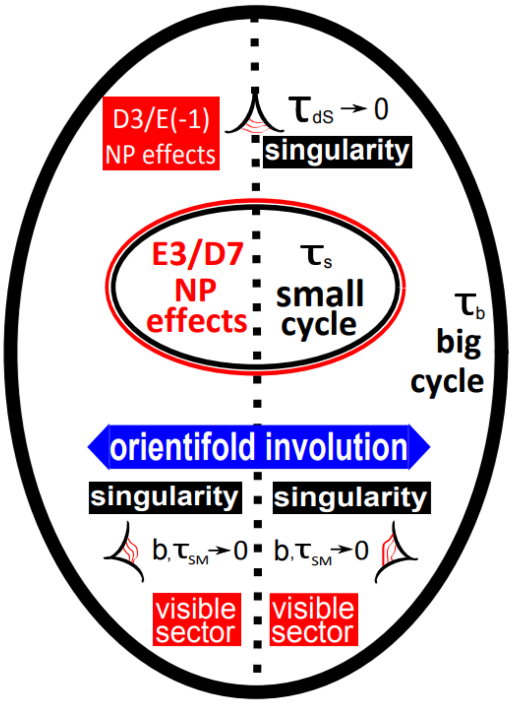

Figure 3.1: Pictorial representation of the Calabi Yau with the blow-up modes (from [39] page 13).

3.2.1 D-term Stabilization

The Kähler moduli that parametrise the location of the visible sector D-brane are stabilised using D-terms. The two moduli T_{SM} and G are charged under two anomalous $U(1)$ symmetries with charges q_1 and q_2, the resulting D-term potential is:

$$V_D = \frac{1}{2Re(f_1)} \left(\sum_\alpha q_{1\alpha} \frac{\partial K}{\partial C^\alpha} C^\alpha - \xi_1 \right)^2 + \frac{1}{2Re(f_2)} \left(\sum_\alpha q_{2\alpha} \frac{\partial K}{\partial C^\alpha} C^\alpha - \xi_2 \right)^2 \tag{3.8}$$

where f_1 and f_2 are the gauge kinetic functions of the two $U(1)$s. The Fayet-Iliopoulos terms are given by:

$$\begin{aligned} \xi_1 &= -\frac{q_1}{4\pi} \frac{\partial K}{\partial T_{SM}} = -\frac{q_1 \lambda_{SM}}{4\pi} \frac{\tau_{SM}}{\mathcal{V}} \\ \xi_1 &= -\frac{q_1}{4\pi} \frac{\partial K}{\partial G} = -\frac{q_1 \lambda_b}{4\pi} \frac{b}{\mathcal{V}}. \end{aligned} \tag{3.9}$$

The D-term potential vanishes in the vacuum in order to have a supersymmetric minimum at $\xi_1 = \xi_2 = 0$, this implies that $\langle \tau_{SM} \rangle = \langle b \rangle = 0$.
The T_{dS} mode can be fixed in the singular regime by minimising the hidden sector D-term potential (focusing for simplicity on canonically normalised hidden fields $\phi_{h,i}$ with charges $q_{h,i}$ under an anomalous $U(1)$).

$$V_D^{dS} = \frac{1}{2Re(f_h)} \left(\sum_i q_{h,i} |\phi_{h,i}|^2 - \xi_h \right)^2 \tag{3.10}$$

with FI-term:

$$\xi_h = -\frac{q_{dS}}{4\pi} \frac{\partial K}{\partial T_{dS}} = -\frac{q_{dS}}{4\pi} \frac{\lambda_{dS} \tau_{dS}}{\mathcal{V}}. \tag{3.11}$$

The hidden matter fields $\phi_{h,i}$ are assumed to be fixed from the related F-term in order to have $\langle \sum_i q_{hid,i} |\phi_{h,i}|^2 \rangle = 0$. The term will be added to the resulting F-term potential to get the total scalar potential.

3.2.2 F-term Stabilization

The dilaton and the complex structure moduli are stabilised supersymmetrically at the flux level providing the conditions:

$$D_S W_{flux}|_{\xi=0} = 0, \quad D_U W_{flux}|_{\xi=0} = 0, \quad \langle W_{flux} \rangle \equiv W_0. \tag{3.12}$$

The minimum provided by this stabilization for the dilaton and the complex structure moduli is shifted by the α' effect and by the T_{dS} dependent term in the superpotential, the shift of the dilaton is parametrised by the function $\omega_S(U,S)$ and for the complex structure by $\omega_{U_i}(U,S)$ providing the covariant derivatives:

$$D_S W \simeq D_S W_{flux}|_{\xi=0} - \frac{3\omega_S(U,S)\hat{\xi} W_0}{4s\mathcal{V}} = -\frac{3\omega_S(U,S)\hat{\xi} W_0}{4s\mathcal{V}}, \tag{3.13}$$

for the complex structure:

$$D_{U_i} W \simeq D_{U_i} W_{flux}|_{\xi=0} - \frac{3\omega_{U_i}(U,S)\hat{\xi} W_0}{4s\mathcal{V}} = -\frac{3\omega_{U_i}(U,S)\hat{\xi} W_0}{4s\mathcal{V}}. \tag{3.14}$$

The explicit forms of the functions ω are visible in the Appendix B.
Considering now the stabilization of the Kähler moduli, as already mentioned in Chapter 1, the α' correction would spoil the no-scale structure of the model providing the following scalar potential (with the stabilization on the vacuum of the axion imaginary part of T_s such that $e^{-a_s Im(T_s)} = -1$):

$$V_F = \frac{1}{2s} \left[\frac{3\hat{\xi} W_0^2}{4\mathcal{V}^3} - 4 \frac{W_0 A_s a_s e^{-a_s \tau_s} \tau_s}{\mathcal{V}^2} + \frac{8}{3} \frac{A_s^2 a_s^2 e^{-2a_s \tau_s} \tau_s^{1/2}}{\mathcal{V}} \right] \tag{3.15}$$

The potential is the usual LVS scalar potential with a exponentially large volume at the minimum but the minimum is modified by the presence of the F-term potential of the T_{dS} moduli which gives a F-term potential (fixing $e^{-i2a_{dS}\varphi_{dS}} = -1$):

$$V_F^{dS} = \frac{(\kappa_{dS} a_{dS} A_{dS})^2}{s\lambda_{dS}} \frac{e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})}}{\mathcal{V}} \tag{3.16}$$

The total potential will be:

$$\begin{aligned} V^{Tot} =& \frac{1}{2s}\left[\frac{3\hat{\xi}W_0^2}{4\mathcal{V}^3} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^2} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}}\right] + \\ &+ \frac{(\kappa_{dS}a_{dS}A_{dS})^2}{s\lambda_{dS}}\frac{e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})}}{\mathcal{V}} + \frac{1}{2s}\left(\frac{q_{dS}}{4\pi}\frac{\lambda_{dS}\tau_{dS}}{\mathcal{V}}\right)^2. \end{aligned} \tag{3.17}$$

The minimum of the potential will be provided in the limit $\epsilon_s = 1/(4a_s\tau_s) \ll 1$ resulting in:

$$\begin{aligned} \langle\tau_s\rangle^{3/2} &= \frac{3}{2}\hat{\xi}\frac{(1-\epsilon_s)^2}{(1-4\epsilon_s)[(4-4\epsilon_s)(1+4\alpha\epsilon_s)-(1-4\epsilon_s)]} = \\ &= \frac{\hat{\xi}}{2}\left[1+\left(\frac{3}{4}-\frac{4}{3}\alpha\right)4\epsilon_s + \left(\frac{3}{4}-\frac{2}{3}\alpha+\frac{16}{3}\alpha^2\right)16\epsilon_s^2 + o(\epsilon_s^3)\right] \\ \langle\mathcal{V}\rangle &= \frac{3W_0\langle\tau_s\rangle^{1/2}}{4a_sA_s}e^{a_s\langle\tau_s\rangle}\frac{1-4\epsilon_s}{1-\epsilon_s} = \frac{3W_0\langle\tau_s\rangle^{1/2}}{4a_sA_s}e^{a_s\langle\tau_s\rangle} + o(\epsilon_s) \\ \langle s\rangle &= \frac{1}{g_s}. \end{aligned} \tag{3.18}$$

With $\alpha = q_{dS}^2\lambda_{dS}\tau_{dS}(1+\lambda_{dS}\tau_{dS})/((4\pi)^2 4W_0A_se^{-a_s\tau_s})$.
The value of τ_{dS} is obtained tuning the scalar potential in order to have a small dS cosmological constant (in the limit $\alpha \ll 1$):

$$\left(\frac{\kappa_{dS}a_{dS}A_{dS}}{W_0}\right)^2 e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})} = \frac{9}{8}\frac{\hat{\xi}\epsilon_s\lambda_{dS}}{\mathcal{V}^2} \tag{3.19}$$

In this way we will fix:

$$\langle\tau_{dS}\rangle = \frac{9\pi a_{dS}\kappa_{dS}W_0^2\hat{\xi}\epsilon_s}{2q_{dS}^2\lambda_{dS}\langle\mathcal{V}\rangle} \tag{3.20}$$

The resulting de Sitter vacuum has the cosmological constant:

$$\langle V^{Tot}\rangle = \frac{1}{2s}\left(\frac{9a_{dS}\kappa_{dS}W_0^2\hat{\xi}\epsilon_s}{8q_{dS}\langle\mathcal{V}\rangle^2}\right)^2 \sim \mathcal{O}(1/\mathcal{V}^4) \tag{3.21}$$

The cosmological constant from the model can also be tuned to be smaller but it has to be compared with the cosmological constant observed value $\Lambda \sim 10^{-122}$ in Planck units.

3.2.3 F-terms

The mass of the gravitino generated by the supersymmetry breaking minimum is:

$$m_{3/2} = e^{K/2}|W| = \frac{W_0}{\sqrt{2s}\mathcal{V}}\left[1+\frac{1}{\mathcal{V}}\left(\lambda_{dS}\tau_{dS} - \frac{\hat{\xi}}{2}\right) + o\left(\frac{1}{\mathcal{V}^2}\right)\right]. \tag{3.22}$$

The mass is poorly affected by the τ_{dS} and $\hat{\xi}$ contributions producing a leading order dependence from the volume of $\mathcal{O}(1/\mathcal{V})$, this is important because all the other quantities

will be traced back to the gravitino mass which will set the scale of the model.
The F-terms related to the LVS Kähler moduli are:

$$\begin{aligned} F^{T_b} &= -\frac{2\tau_b W_0}{(2s)^{1/2}\mathcal{V}} + o\left(\frac{1}{\mathcal{V}^2}\right) \sim \mathcal{O}\left(m_{3/2}^{1/3} M_p^{5/3}\right) \\ F^{T_s} &= -\frac{3W_0}{2\mathcal{V} a_s (2s)^{1/2}} + o\left(\frac{1}{\mathcal{V}^2}\right) + o\left(\frac{1}{a_s^2 \tau_s}\right) \sim \mathcal{O}\left(m_{3/2} M_p\right) \end{aligned} \tag{3.23}$$

The F-terms of the two moduli are at the leading order the same of a usual LVS model with F^{T_b} the dominant term in the supersymmetry breaking.
The F-terms associated to the dilaton S and the complex structure moduli U_i should be zero at the leading order by the supersymmetric flux stabilization but they are affected by the shift of the minimum induced by the α' correction.

$$\begin{aligned} F^S &= e^{K/2} W_0 \frac{3}{2}\frac{\hat{\xi} s}{\mathcal{V}}\left[3 - 2\omega_S(U,S)\right] + o\left(\frac{1}{\mathcal{V}^3}\right) \sim \mathcal{O}\left(m_{3/2}^2\right) \\ F^{U_i} &= e^{K/2} K^{U_i \overline{U}_j} D_{\overline{U}_j} \overline{W} \simeq K^{U_i \overline{U}_j} \frac{\omega_{\overline{U}_j}(U,S)}{2s^2 \omega'_S(U,S)} F^S = \beta^{U_i}(U,S) F^S \sim \mathcal{O}(m_{3/2}^2) \end{aligned} \tag{3.24}$$

with $\omega'_S(U,S) = 3 - 2\omega_S(U,S)$. The β^{U_i} terms are $\mathcal{O}(1)$ in the volume expansion. The two F-terms are hierarchically smaller respect to the ones of the Kähler moduli.
The F-term associated to the new Kähler moduli T_{dS} is:

$$F^{T_{dS}} = \frac{3W_0 \hat{\xi}^{1/2} \epsilon_s^{1/2}}{\sqrt{2}\lambda_{dS}^{1/2} (2s)^{1/2} \mathcal{V}} + o\left(\frac{1}{\mathcal{V}^2}\right) \sim \mathcal{O}\left(\frac{m_{3/2}}{\epsilon_s^{1/4}}\right). \tag{3.25}$$

The last F-terms associated to the moduli are the ones of G and T_{SM} which are fixed to 0:

$$F^G = 0, \qquad F^{T_{SM}} = 0 \tag{3.26}$$

which is very important to provide the sequestering because the dominant F-terms are the ones associated to the moduli weakly coupled with the visible sector.

3.3 Soft-Terms

The soft-supersymmetry breaking Lagrangian of the model is:

$$-\mathcal{L}_{soft} = \frac{1}{2}\left(M_a \lambda^a \lambda^a + h.c.\right) + m_\alpha^2 C^\alpha \bar{C}^{\bar{\alpha}} + \left(\frac{1}{6} A_{\alpha\beta\gamma} \hat{Y}_{\alpha\beta\gamma} C^\alpha C^\beta C^\gamma + B\mu H_u H_d + h.c.\right) \tag{3.27}$$

where all the fields and the parameters are the normalized ones.
Before seeing the form of the Soft-Terms is important to look at the factorization of the matter metric which distinguishes two possible phenomenological scenarios.

Factorisation of the matter metric

The general factorization of the matter metric follows what said in the Chapter 1 considering the new terms that arise from the α' correction and the T_{dS} contribution:

$$\tilde{K}_\alpha = \frac{h_\alpha(U,S) e^{K_{cs}/3}}{(2s)^{1/3} \mathcal{V}^{2/3}} \left(1 - c_s \frac{\hat{\xi}}{\mathcal{V}} + c_{dS} K_{dS}\right). \tag{3.28}$$

There can be defined two scenarios considering if the equation $\tilde{K} = e^{K/3}$ holds only at the leading order or also for higher order corrections:

- Local: if the matter metric factorization holds only at leading order in $\mathcal{V}^{-1}$;
- Ultra-local: if the matter metric is factorized in the mentioned way also at higher orders.

In the Ultra-local scenario:

$$\tilde{K}_\alpha = h_\alpha(U,S)e^{K/3} + o\left(\frac{1}{\mathcal{V}^{8/3}}\right) \tag{3.29}$$

and $c_s = c_{dS} = 1/3$. In fact:

$$e^{K/3} = \frac{e^{K_{cs}/3}}{(2s)^{1/3}\mathcal{V}^{2/3}}\left(1 - \frac{\hat{\xi}}{3\mathcal{V}} + \frac{\lambda_{dS}\tau_{dS}^2}{3\mathcal{V}} + \frac{5}{36}\frac{\hat{\xi}^2}{\mathcal{V}^2} + \frac{1}{18}\frac{\lambda_{dS}^2\tau_{dS}^4}{\mathcal{V}^2} - \frac{1}{18}\frac{\hat{\xi}\lambda_{dS}\tau_{dS}}{\mathcal{V}^2} + o\left(\frac{1}{\mathcal{V}^3}\right)\right). \tag{3.30}$$

The scenarios are very important because they distinguish two phenomenological different spectra for the MSSM which can lead to the presence of FCNCs and this will be object of study of the Section 3.4.

3.3.1 Gaugino Masses

The gaugino masses in gravity mediated supersymmetry breaking are:

$$M_a = \frac{1}{2}\frac{F^m\partial_m f_a}{\mathrm{Re}(f_a)}. \tag{3.31}$$

By the fact that on the vacuum $f_a = \delta_a S$ and the δ_a are universal terms (but could be also non universal for different models) the gaugino masses of the model are universal:

$$M_{1/2} = \frac{1}{2}\frac{F^S\partial_S f}{\mathrm{Re}(f)} = \frac{3}{4}\frac{\hat{\xi}W_0}{(2s)^{1/2}\mathcal{V}^2}\omega_S'(U,S) \sim \mathcal{O}\left(\frac{m_{3/2}^2}{M_p}\right). \tag{3.32}$$

The resulting masses are hierarchically smaller than the gravitino mass $m_{3/2}$.

3.3.2 Scalar Masses

The scalar masses are generated by both F-term and D-term contributions by the D-term stabilization of the T_{dS} moduli.

F-term Contribution

The F-term contribution to the scalar masses for diagonal matter metric is:

$$m_\alpha^2|_F = m_{3/2}^2 + \langle V\rangle - \overline{F}^{\overline{m}}F^n\partial_{\overline{m}}\partial_n \log(\tilde{K}_\alpha) \tag{3.33}$$

Local Scenario

The scalar masses in the local scenario are:

$$
\begin{aligned}
m_\alpha^2 &= m_{3/2}^2 + \langle V \rangle - \overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \log(\tilde{K}_\alpha) = \\
&= m_{3/2}^2 - \overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \left[\frac{1}{3} K + \log \left(\frac{1 - c_s \frac{\hat{\xi}}{\mathcal{V}} + c_{dS} K_{dS}}{1 - \frac{\hat{\xi}}{3\mathcal{V}} + \frac{1}{3} K_{dS}} \right) + \log(h_\alpha) \right] = \\
&= -\overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \left[\log \left(\frac{1 - c_s \frac{\hat{\xi}}{\mathcal{V}} + c_{dS} K_{dS}}{1 - \frac{\hat{\xi}}{3\mathcal{V}} + \frac{1}{3} K_{dS}} \right) + \log(h_\alpha) \right] = \\
&= \left(c_s - \frac{1}{3} \right) \left(F^{T_b} \right)^2 \partial_{T_b}^2 \frac{\hat{\xi}}{\mathcal{V}} + o \left(\frac{1}{\mathcal{V}^4} \right) = \\
&= \left(c_s - \frac{1}{3} \right) \frac{5}{\omega_S'(U,S)} m_{3/2} M_{1/2} + o \left(\frac{1}{\mathcal{V}^4} \right) \sim \mathcal{O}(M_{1/2} m_{3/2}).
\end{aligned} \tag{3.34}
$$

In the local scenario the masses of the scalar particles are mainly generated by the F^{T_b} contribution. The masses obtained are flavour universal and are hierarchically between the gravitino mass $m_{3/2}$ and the universal gaugino mass $M_{1/2}$. For $c_s < 1/3$ the resulting masses are negative, the tachyonic solution would provide a non stable vacuum for the MSSM fields.

Ultra-local Scenario

The ultra-local scenario provides a different result:

$$
\begin{aligned}
m_\alpha^2 &= m_{3/2}^2 + \langle V \rangle - \overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \log(\tilde{K}_\alpha) = \\
&= m_{3/2}^2 - \overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \left[\frac{1}{3} K + \log(h_\alpha) \right] = \\
&= -\overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \log(h_\alpha) = \\
&= -(F^S)^2 \left(\partial_S^2 + \beta^U \partial_S \partial_U + \beta^U \beta^{\overline{U}} \partial_U^2 \right) \log(h_\alpha) = \\
&= -s^2 M_{1/2}^2 \left(\partial_s^2 + \beta^{U_i} \partial_s \partial_{u_i} + \beta^{U_i} \beta^{\overline{U}_j} \partial_{u_i} \partial_{u_j} \right) \log(h_\alpha) \sim \mathcal{O}(M_{1/2}^2)
\end{aligned} \tag{3.35}
$$

The masses in the ultra-local scenario are generated by the F^S contribution. The mass spectrum in this scenario is different from the one in the local one and has the masses of the scalar particles at the same order of the mass of the gauginos. The $\log(h_\alpha)$ factor can provide non flavour universal masses, inducing FCNCs.

D-term Contribution

The D-term contribution to the scalar masses is given by:

$$
m_\alpha^2|_D = K_\alpha^{-1} \sum_i g_i^2 D_i \partial_{\alpha\overline{\alpha}}^2 D_i - V_{D,0}. \tag{3.36}
$$

The only term is the one associated to the D-term of T_{dS} associated to the anomalous $U(1)$ symmetry which is:

$$
D_{T_{dS}} = \frac{q_{dS}}{4\pi} \partial_{T_{dS}} K. \tag{3.37}
$$

And $g_i^2 = 1/\mathrm{Re}(f_a) =$. The contribution will be:

$$\begin{aligned} m_\alpha^2|_D &= \frac{q_{dS}}{4\pi} K_\alpha^{-1} g_{T_{dS}}^2 D_{T_{dS}} \partial_{T_{dS}} \partial_C \partial_{\overline{C}} K - V_{D,0} = \\ &= \frac{q_{dS}}{4\pi s} K_\alpha^{-1} D_{T_{dS}} \partial_{T_{dS}} K_\alpha - V_{D,0} = \frac{c_{dS}}{4\pi s} D_{dS} q_{dS} \frac{\tau_{dS}}{\mathcal{V}} - V_{D,0} = (2c_{dS} - 1) V_{D,0}. \end{aligned} \tag{3.38}$$

Adding also $V_{F,0}/3 = -V_{D,0}/3$ the final result is

$$m_\alpha = 2\left(c_{dS} - \frac{1}{3}\right) V_{D,0} \sim \mathcal{O}(\mathcal{V}^{-4}). \tag{3.39}$$

In the local scenario the contribution is $\mathcal{O}(\mathcal{V}^{-4})$ so it is negligible respect to the contribution from the F-term. In the ultra-local scenario $c_s = 1/3$ so the $\mathcal{O}(\mathcal{V}^{-4})$ contribution is removed giving that the dominant contribution is the F-term one.

3.3.3 Trilinear terms

The trilinear A-terms are defined for a diagonal matter metric as (in the leading order contribution $F^m \partial_m \log Y_{\alpha\beta\gamma} = 0$):

$$A_{\alpha\beta\gamma} = F^m \left[K_m \delta_{\alpha\beta\gamma} - \partial_m \log\left(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma\right)\right]. \tag{3.40}$$

Local scenario

In the local scenario the result for the trilinear A-terms is:

$$\begin{aligned} A_{\alpha\beta\gamma} &= F^m \left[K_m \delta_{\alpha\beta\gamma} - \partial_m \log\left(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma\right)\right] = \\ &= F^m \partial_m \left[K - \log\left(e^K h_\alpha h_\beta h_\gamma\right) - 3\log\left(\frac{1 - c_s \frac{\hat{\xi}}{\mathcal{V}} + c_{dS} K_{dS}}{1 - \frac{\hat{\xi}}{3\mathcal{V}} + \frac{1}{3} K_{dS}}\right)\right] = \\ &= F^m \partial_m \left[-\log\left(h_\alpha h_\beta h_\gamma\right) - 3\left(\frac{1}{3} - c_s\right)\frac{\hat{\xi}}{\mathcal{V}} - 3\left(c_{dS} - \frac{1}{3}\right) K_{dS} + o\left(\frac{1}{\mathcal{V}^2}\right)\right] = \\ &= \left[-s\partial_{s,u} \log\left(h_\alpha h_\beta h_\gamma\right) - \frac{6}{\omega_S'(S,U)}\left(\frac{1}{3} - c_s\right)\right] M_{1/2} \sim \mathcal{O}(M_{1/2}) \end{aligned} \tag{3.41}$$

with $\partial_{s,u} = \partial_s + \beta^U \partial_u$. The trilinear terms are mainly generated by the F^S term contribution, they are universal if the term $s\partial_{s,u} \log\left(h_\alpha h_\beta h_\gamma\right)$ is negligible respect to the other terms producing trilinear terms of the order $\mathcal{O}(M_{1/2})$ which are proportional to the Yukawa couplings.

Ultra local Scenario

In the ultra-local scenario the trilinear terms are:

$$\begin{aligned} A_{\alpha\beta\gamma} &= F^m \left[K_m \delta_{\alpha\beta\gamma} - \partial_m \log\left(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma\right)\right] = \\ &= F^m \partial_m \left[K - \log\left(e^K h_\alpha h_\beta h_\gamma\right)\right] = \\ &= F^m \partial_m \left[-\log\left(h_\alpha h_\beta h_\gamma\right)\right] + o\left(\frac{1}{\mathcal{V}^3}\right) = \\ &= -s\partial_{s,u} \log\left(h_\alpha h_\beta h_\gamma\right) M_{1/2} \sim \mathcal{O}(M_{1/2}). \end{aligned} \tag{3.42}$$

The trilinear terms are, also in this scenario, generated by the F^S term but in this case they are not universal because of the cancellation which leaves only the $s\partial_{s,u}\log(h_\alpha h_\beta h_\gamma)$ term but this does not change the hierarchy leaving them at the $\mathcal{O}(M_{1/2})$ order.

3.3.4 Bμ-term

F-term contribution

The last soft-term is the $B\mu$-term. Parametrizing $Z = \gamma\tilde{K}_M$ with $\tilde{K}_M$ the flavour universal part of the matter metric and assuming $\mu = 0$ the term is:

$$\begin{aligned}
B\mu|_F = (\tilde{K}_{H_u}\tilde{K}_{H_d})^{-1/2}\{&(2m_{3/2}^2 + \langle V\rangle)Z - m_{3/2}\overline{F}^{\overline{m}}\partial_{\overline{m}}Z + \\
&+ m_{3/2}F^m\left[\partial_m Z - Z\partial_m log(\tilde{K}_{H_u}\tilde{K}_{H_d})\right] + \\
&-\overline{F}^{\overline{m}}F^n\left[\partial_{\overline{m}}\partial_n Z - (\partial_{\overline{m}}Z)\partial_n log(\tilde{K}_{H_u}\tilde{K}_{H_d})\right]\Big\} = \\
= \frac{1}{Z}\{&(2m_{3/2}^2 + \langle V\rangle)Z - 2Zm_{3/2}F^m\partial_m log(Z) + \\
&-\overline{F}^{\overline{m}}F^n\left[\partial_{\overline{m}}\partial_n Z - 2(\partial_{\overline{m}}Z)\partial_n log(Z)\right]\Big\} = \\
= \frac{\gamma}{\sqrt{f_{H_u}f_{H_d}}}\{&2m_{3/2}^2 - 2m_{3/2}F^i\partial_i log(\tilde{K}_H) + \\
&-F^i\overline{F}^{\overline{j}}\left[\partial_i\partial_{\overline{j}}log(\tilde{K}_M) - \partial_i log(\tilde{K}_M)\partial_{\overline{j}}log(\tilde{K}_H)\right]\Big\} = \\
= \frac{\gamma}{\sqrt{f_{H_1}f_{H_2}}}&\left[m_{3/2}^2 - F^i\overline{F}^{\overline{j}}\partial_i\partial_{\overline{j}}log(\tilde{K}_M) + o\left(\frac{1}{\mathcal{V}^4}\right)\right] = \\
= \frac{\gamma}{\sqrt{f_{H_u}f_{H_d}}}&\left[m^2|_F + o\left(\frac{1}{\mathcal{V}^4}\right)\right]
\end{aligned} \tag{3.43}$$

In the Local and Ultra-local scenario the behaviour of the $B\mu$ term is proportional to the flavour universal term of the scalar masses, reproducing the same hierarchy.

D-term contribution

The D-term contribution is given by:

$$\begin{aligned}
B\mu|_D &= (K_{H_u}K_{H_d})^{-1/2}\left(\sum_i g_i^2 D_{iH_u}\partial_{H_d}D_i - V_{D,0}Z\right) = \\
&= \frac{\gamma}{\sqrt{f_{H_u}f_{H_d}}}K_M^{-1}\left(\frac{q_{dS}}{4\pi}g_{dS}^2 D_{T_{dS}}\partial_{T_{dS}}K_M - V_{D,0}K_M\right) = \\
&= \frac{\gamma}{\sqrt{f_{H_u}f_{H_d}}}m_\alpha^2|_D.
\end{aligned} \tag{3.44}$$

Also for the D-term contribution the behaviour of $B\mu$ is the same as the D-term contribution of the scalar masses, this means that again is negligible respect to the F-term contribution.

3.3.5 Mass Spectra

The two mass spectra for the different scenarios are summarised in Figure 3.2.

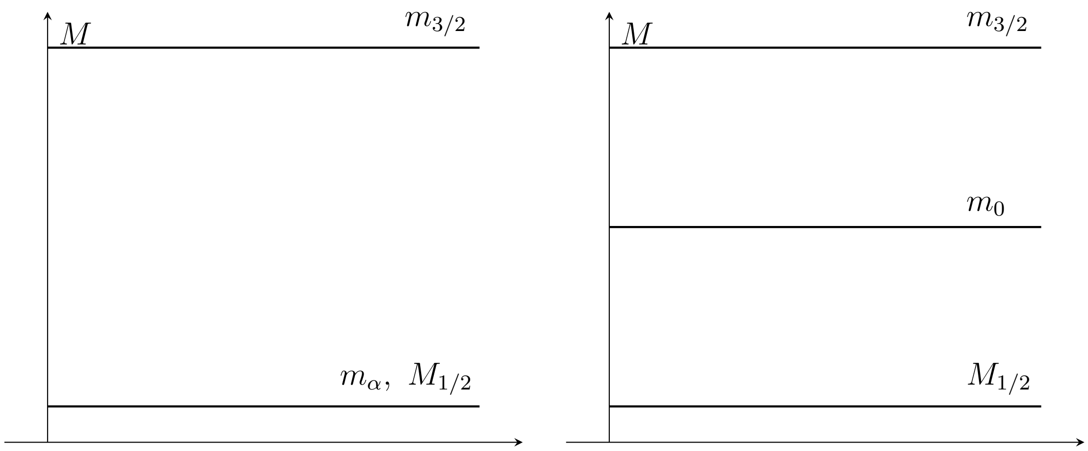

Figure 3.2: The two possible mass spectra of the model. (Left) Ultra-local scenario. (Right) Local scenario.

The Ultra-local scenario spectrum is similar to the one obtained from D7-branes while the Local scenario gives a split supersymmetry breaking spectrum. Considering the compactification volume expressed in [39] $\mathcal{V} \sim 10^7$ the spectrum will be $m_{3/2} \sim 10^{10}$ TeV, $M_{1/2} \sim 1$ TeV and in the local scenario $m_0 \sim 10^4$ TeV.
The split of the supersymmetry in the local scenario influences the renormalization of the mass of the Higgs in the RG group. In fact we have looked at the study [40] where it was computed the RG flow of the Higgs mass in the MSSM with different supersymmetry breaking scenarios. In the split supersymmetry breaking scenario, fixing $M_{1/2} \sim 1$ TeV, the mass of the Higgs is reproduced for the chosen flavour universal mass of the scalar particles for $\tan\beta = 1$. The β angle is defined as:

$$\begin{pmatrix} H \\ A \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} = \begin{pmatrix} -\epsilon H_d^* \\ H_u \end{pmatrix} \tag{3.45}$$

where ϵ is the antisymmetric tensor with $\epsilon_{12} = 1$. H is the SM-like Higgs doublet and A is an heavier Higgs doublet. $\tan\beta$ is then defined by the tuning condition:

$$\tan^2\beta = \left.\frac{m_{H_d}^2 + \mu^2}{m_{H_u}^2 + \mu^2}\right|_{m_0} . \tag{3.46}$$

By this result, we can see that the hierarchy of the model provides a strict bound on the MSSM RG parameters.
The lightest supersymmetric particle in this split MSSM and the best candidate to be experimentally observed is the gluino which has been widely studied from a phenomenological point of view as, for example, done in [41] where the gluino signature are studied in term of its lifetime in fact from [40] the gluino lifetime is:

$$c\tau_{\tilde{g}} = \left(\frac{2\ \text{TeV}}{M_{1/2}}\right)^5 \left(\frac{m_0}{10^7\ \text{GeV}}\right)^4 0.4\ \text{m}. \tag{3.47}$$

By our prediction of the spectrum in the local scenario the gluino lifetime result is $c\tau_{\tilde{g}} \simeq 13$ m, in this Long-lived gluino scenario the best possibility to observe it is when the gluino hadronizes and turn into an R-hadron [41]. The R-hadrons are massive charged particles colour singlet hadronic states.

3.4 Sequestering and Mirror Mediation

Before looking at the FCNCs analysis of the mass spectra and the trilinear terms is possible to see if the five Mirror Mediation conditions are respected by the model.

1 The Hidden sector is made by the modulus T_b, T_s, T_{dS}, S, U, T_{SM}, G providing the distinction of two classes:

- $\Phi_i = T_b, T_s, T_{dS}, T_{SM}, G$;
- $\chi_i = S, U$.

So the model matches the first condition.

2 The Kähler potential is a sum at the zero order in $1/\mathcal{V}$ but the first order corrections spoil the sum form introducing mixing terms between Φ_i and χ_i:

$$K = \underbrace{-2\ln(\mathcal{V}) + \lambda_{dS}\frac{\tau_{dS}^2}{\mathcal{V}} + \lambda_{SM}\frac{\tau_{SM}^2}{\mathcal{V}} + \lambda_b\frac{b^2}{\mathcal{V}}}_{\Phi_i \text{ sector}} - \underbrace{\ln(2s) + K_{cs}(U)}_{\chi_i \text{ sector}} - \underbrace{\frac{\hat{\xi}}{\mathcal{V}} + o\left(\frac{1}{\mathcal{V}^2}\right)}_{\text{mixing terms}}. \tag{3.48}$$

The condition is respected only at the leading order but as exposed in the Chapter 2 the mixing terms in the Kähler metric are suppressed respect to the diagonal ones in the LVS.

3 The Yukawa couplings depends only on the χ_i because of the shift symmetry of the Khäler moduli and the holomorphy of the superpotential. The gauge kinetic function depends both on Φ_i and χ_i, in fact in the model:

$$f_a = \delta_a S + \kappa_a T_{SM}. \tag{3.49}$$

The dominant term in the gauge kinetic function is the dilaton dependence which is from the χ_i class but it should come from the Φ_i class, so the Mirror Mediation condition on the gauge kinetic function is not respected by the model.

4 The factorization of the matter metric is broken by the term of order $1/\mathcal{V}^{5/3}$:

$$K_\alpha = \underbrace{f_\alpha(U,S)}_{\chi \text{ factor}} \underbrace{\frac{1}{\mathcal{V}^{2/3}}\left(1 - c_{SM}\tau_{SM}^p + c_b b^p + c_{dS}\frac{\tau_{dS}^2}{\mathcal{V}}\right)}_{\Phi_i \text{ factor}} - \underbrace{\frac{f_\alpha(U,S)c_s\hat{\xi}}{2\mathcal{V}^{5/3}}}_{\text{mixing term}}. \tag{3.50}$$

The mixing term is suppressed respect to the factorized one so at the leading order the Mirror Mediation condition is respected but at the higher order is broken.

5 The non-supersymmetric stabilization of Φ_i is respected while the supersymmetric one of χ_i is not respected, in fact:

$$\left.\begin{array}{l} F^{T_b} = -\frac{2\tau_b W_0}{(2s)^{1/2}\mathcal{V}} + o\left(\frac{1}{\mathcal{V}^2}\right) \\ F^{T_s} = -\frac{3W_0}{2\mathcal{V}a_s(2s)^{1/2}} + o\left(\frac{1}{\mathcal{V}^2}\right) \\ F^{T_{dS}} = \frac{3W_0\hat{\xi}^{1/2}\epsilon_s^{1/2}}{\sqrt{2}\lambda_{dS}^{1/2}(2s)^{1/2}\mathcal{V}} + o\left(\frac{1}{\mathcal{V}^2}\right) \\ F^{T_{SM}} = 0 \\ F^G = 0 \end{array}\right\} F^{\Phi_i} \neq 0$$
$$\left.\begin{array}{l} F^S = e^{K/2}W_0\frac{3}{2}\frac{\hat{\xi}_s}{\mathcal{V}}\left[3 - 2\omega_S(U,S)\right] \\ F^U = K^{U_i\overline{U}_j}\frac{\omega_{\overline{U}_j}(U,S)}{2s^2\omega'_S(U,S)}F^S \end{array}\right\} F^{\chi_i} \neq 0 \text{ but } o(1/\mathcal{V}^2) \tag{3.51}$$

The F^S and F^U terms are suppressed respect to the Φ_i class' F-terms so the Mirror Mediation condition is again respected only at the leading order.

The Mirror Mediation analysis shows that the Mirror Mediation conditions are always matched at the leading order (except for what concern the gauge kinetic function) but they are broken once one considers also the subleading contributions so a cancellation at the leading order could make emerge the non Mirror Mediation terms which would produce FCNCs.
As shown in this Section the Mirror Mediation conditions are not exactly respected and this could lead to non flavour-universal masses (as what happens to the Ultra-local scenario), it is then fundamental to understand if the experimental bound (2.20) is respected or not by the two scalar masses. Once verified the experimental compatibility of the scalar masses we considered the UFB and CCB conditions (2.27).

3.4.1 Ultra-local Scenario

In the Ultra-local scenario the scalar masses are flavour non universal with the form:

$$m_\alpha^2 = M_{1/2}^2\Delta_{s,u}\log h_\alpha(U,S) \tag{3.52}$$

with $\Delta_{s,u} = -s^2\left(\partial_s^2 + \beta^{U_i}\partial_{u_i}\partial_s + \beta^{U_i}\beta^{\overline{U}_j}\partial_{u_i}\partial_{\overline{u}_j}\right)$. All the phenomenological implications came out from the exact functional form of $h_\alpha(U,S)$. We could try to consider an hypothetical parametrization that could preserve the flavour universality:

$$h_\alpha = g(U,S)g_\alpha(U,S). \tag{3.53}$$

In this parametrization the scalar masses could be divided into a universal term and a non-universal term:

$$\begin{array}{l} m^2 = M_{1/2}^2\Delta_{s,u}\log g(U,S) \\ \Delta m_\alpha^2 = M_{1/2}^2\Delta_{s,u}\log g_\alpha(U,S). \end{array} \tag{3.54}$$

Now looking at the experimental bound (2.20) the bound on the non universality of $h_\alpha(U,S)$ will be (considering $m \sim 10^4$ TeV):

$$\Delta_{s,u}\log g(U,S) \gtrsim 10^7\Delta_{s,u}\log g_\alpha(U,S). \tag{3.55}$$

The resulting suppression is huge and it is the result of a non supported hypothesis on the matter metric. The resulting mass in the Ultra-local scenario should make us agree on the fact that it is not compatible with the FCNCs constraints of the SM.

3.4.2 Local Scenario

For the local scenario the flavour universal mass term is dominant but there is also a flavour non universal term which is suppressed in the volume expansion but has to be analysed:

$$\begin{aligned} m_0^2 &= \left(c_s - \frac{1}{3}\right) \frac{5}{\omega_S'(U,S)} m_{3/2} M_{1/2} \\ \Delta m_\alpha^2 &= M_{1/2}^2 \Delta_{s,u} \log h_\alpha(U,S). \end{aligned} \tag{3.56}$$

The resulting lower bound for the volume is:

$$\mathcal{V} \gtrsim 10^{-2} \frac{3}{4} \frac{\hat{\xi} \omega_S'^2(U,S) \Delta_{s,u} \log h_\alpha(U,S)}{(2s)^{1/2} \left(c_s - \frac{1}{3}\right)}. \tag{3.57}$$

All the terms in the RHS are of the unity order providing that the bound is simply $\mathcal{V} \gtrsim 10^{-2}$. The resulting bound is compatible with the volume proposed in [39] $\mathcal{V} \sim 10^7$.

3.4.3 Literature Results

Comparing this result to the one obtained by de Alwis in [42] we obtained a low difference but a significant discrepancy in how the result is obtained. The result obtained in [42] is a lower bound for the volume of $\mathcal{V} \gtrsim 10^{-3}$ always considering the mass of the scalar as 10^4 TeV.
The discrepancy is generated by the matter metric considered by [42] which has been computed in [43].

$$\tilde{K}_{\alpha\bar{\beta}} = \frac{3\mu}{\mathcal{V} + \hat{\xi}/2} \left(i\omega^b_{\alpha\bar{\beta}} \tau_b^{1/2} - i\omega^s_{\alpha\bar{\beta}} \tau_s^{1/2} \right) \tag{3.58}$$

with ω^i the harmonic (1,1)-forms on the Calabi Yau orientifold evaluated at the position of the D3-breane. The computation is made considering the D3-brane in a fixed point and not in a singularity and it is improperly extended to the singularity case. In fact the non-universal mass term in [42] is:

$$\Delta m_\alpha^{2\beta} = m_{3/2}^2 \frac{3}{4} \sqrt{\frac{\tau_s}{\tau_b}} K'^{\ \beta}_{\alpha}. \tag{3.59}$$

The $\sqrt{\tau_s/\tau_b}$ term came out from the form of the matter metric in fact we have compared the de Alwis result for the matter metric with the result expressed in Chapter 1 and we have got that (assuming for simplicity $3\mu = i\omega^b_{\alpha\bar{\beta}} = i\omega^s_{\alpha\bar{\beta}} = 1$):

$$e^{K/2} \frac{1}{\tilde{K}^{3/2}} = 1 + \frac{3}{2} \frac{\tau_s^{1/2}}{\mathcal{V}^{1/3}} + o\left(\frac{1}{\mathcal{V}^{2/3}}\right). \tag{3.60}$$

Instead of having only 1 up to the $\hat{\xi}$ corrections the RHS has also a $\mathcal{O}\left(\sqrt{\tau_s/\tau_b}\right)$ term which is the origin of the discrepancy between our result and the one in [42].

3.4.4 CCB and UFB Constraints

Once we verified that the FCNC constraints are always respected by the local scenario we had to consider the CCB and UFB constraints that we mentioned in Section 2.2.3 and 2.2.4 which can be essentially summarised as:

$$A^2_{\alpha\beta\gamma} < m_0^2. \tag{3.61}$$

Considering now the the value of the trilinear terms and the scalar masses we have seen that the local scenario respects the condition in fact: $A^2_{\alpha\beta\gamma} \sim M^2_{1/2} \ll m_0^2 \sim m_{3/2} M_{1/2}$. The local scenario is acceptable also under this constraints.
The Ultra-local scenario which is already problematic from the mass-universality does not respect the CCB and UFB constraints because the terms are very similar in fact this would require a very peculiar fine tuning because the trilinear terms and the scalar masses are both proportional to the gaugino mass. The condition (3.61) becomes:

$$[\partial_{s,u} \log\left(h_\alpha h_\beta h_\gamma\right)]^2 < -\left(\partial_s^2 + \beta^{U_i}\partial_s\partial_{u_i} + \beta^{U_i}\beta^{\overline{U}_j}\partial_{u_i}\partial_{u_j}\right)\log(h_\alpha). \tag{3.62}$$

The solution of the inequality can be found only knowing the exact form of $h_\alpha(U,S)$ also by the fact the the β terms are $\mathcal{O}(1)$ so the most realistic hypothesis that we can make is that the condition is not respected, producing another phenomenological problem against the Ultra-local scenario.

Conclusions

String compactifications can provide many useful supersymmetric models with different phenomenological implications. Sequestered models provide a good solution to the hierarchy problem and to the cosmological moduli problem. However models with low-energy masses for the scalars can be affected by flavour changing neutral currents which one want to avoid in order to preserve the Standard Model predictions.
The model that we have analysed provides two different phenomenological scenarios for the MSSM: the ultra-local and the local one. It is important to notice that the two scenarios are distinguished only by different higher-order corrections to the Kähler metric for MSSM matter fields.

- The Ultra-Local scenario features a standard supersymmetry breaking spectrum where all soft-terms are of the same order of magnitude $M_{1/2} \sim m_\alpha \sim A_{\alpha\beta\gamma} \sim B\mu \sim 1$ TeV $\ll m_{3/2}$ with a low energy supersymmetry soft breaking as in models with matter originated by D7-branes, but introduces flavour non-universal scalar masses due to the functional dependence of the matter metric. In fact, the scalar masses are generated by the α' correction. Combining the definition of the F-term potential $\overline{F}^{\overline{m}} F^n K_{\overline{m}n} = V_0 + 3m_{3/2}^2$ with the form of the matter metric which is $\tilde{K}_\alpha = e^{K/3} h_\alpha$, the scalar masses become:
$$m_\alpha^2 = \frac{2}{3} V_0 - \frac{1}{3} \overline{F}^{\overline{m}} F^n \partial_{\overline{m}} \partial_n \log h_\alpha. \tag{3.63}$$
So by the fact that h_α depends only on S and U, the only contribution is the one generated by the shift of the minimum for S and U induced by α' effects. This means that the normalized scalar mass depends on the flavour structure of the matter metric. This implies that in general there is no way to assume a flavour universality. Once we have moved to the analysis of the trilinear terms we have seen that also in this case the Ultra-local scenario does not respect clearly the constraints. In fact, the exact flavour structure of the matter metric became important also for the trilinears and by our analysis, we can conclude that this scenario does not respect also the UFB and CCB constraints.
We can conclude that this scenario is phenomenological unacceptable and this could make us argue that the matter metric should be affected by the α' correction in a different way with respect to the Kähler potential and the definition from [16] should be exact only in the first order of the volume expansion, at least for D3-brane models.

- The Local scenario provides a different phenomenological result named in literature as split supersymmetry breaking. In fact, there is a splitting between the soft terms leading to $M_{1/2} \sim A_{\alpha\beta\gamma} \sim 1$ TeV and $m_0 \sim B\mu \sim 10^4$ TeV, both smaller than the gravitino mass $m_{3/2} \sim 10^{11}$ TeV. It is important to notice that this scenario is acceptable if and only if the coefficient of the Kähler matter metric $c_s > 1/3$ in order to provide positively defined scalar masses otherwise this would introduce an unstable vacuum for the MSSM which can break the symmetry in the MSSM. The split of the supersymmetry breaking does not spoil the naturalness of the model since for $\tan\beta = 1$ the mass spectrum gives the right renormalization group running of the Higgs mass as demonstrated by [40]. Contrary to the Ultra-local, this scenario respects the FCNC constraints on the scalar masses, as we have demonstrated in Section 3.4.2, and respects the UFB and CCB constraints as demonstrated in Section 3.4.4. As in typical split supersymmetry breaking scenario, the lightest supersymmetric particles are the gluinos which can be the object of experimental studies. In particular, this model predicts a long-lived gluino with lifetime $c\tau_{\tilde{g}} \simeq 13$ m.

Both of the two phenomenological scenarios are characterized by the presence of the uplifting mechanism via the new Kähler moduli that can provide a fine-tuning for the cosmological constant value.
In the end the result of the phenomenological unacceptability of the ultra-local model could make us infer that the matter metric should be affected by different α' corrections with respect to the Kähler potential. In particular, c_s should be larger than 1/3 to provide positively defined mass terms. This means that the physical Yukawa couplings will depend on the volume via higher-order corrections, even if one would expect the opposite due to the fact that the Yukawa couplings are generated by the local interactions of the open string degrees of freedom. An interesting follow-up research line could be the study of the effect of the α' corrections to the matter metric and the understanding of the implications of the volume dependence of the physical Yukawa couplings.

Appendix A

Single-Hole Swiss Cheese Model

The Kähler potential is:

$$K = -2\log\left(\mathcal{V} + \frac{\hat{\xi}}{2}\right) \mathcal{V} = \left(\frac{T_b + \overline{T_b}}{2}\right)^{3/2} - \left(\frac{T_s + \overline{T_s}}{2}\right)^{3/2} \tag{A.1}$$

the superpotential is:

$$W = W_0 + A_s e^{-a_s T_s} \tag{A.2}$$

the supergravity F-term scalar potential is:

$$V_F = e^K \left(K^{i\bar{j}} D_i W D_{\bar{j}} \overline{W} - 3|W|^2\right).(\text{A.3})$$

The derivative of the Kähler potential are:

$$K_{T_b} = -\frac{3}{2}\frac{\tau_b^{1/2}}{(\mathcal{V} + \hat{\xi}/2)} \quad K_{T_s} = \frac{3}{2}\frac{\tau_s^{1/2}}{(\mathcal{V} + \hat{\xi}/2)} \tag{A.4}$$

with $\tau_i = Re(T_i)$, the inverse metric is:

$$K^{i\bar{j}} = \frac{8}{3}\frac{[\mathcal{V} + \hat{\xi}/2]}{[2\mathcal{V} - \hat{\xi}/2]} \begin{pmatrix} 3\tau_s^{3/2}\tau_b^{1/2} + \tau_b^{1/2}[\mathcal{V} + \hat{\xi}/2] & 3\tau_s\tau_b \\ 3\tau_s\tau_b & 3\tau_b^{3/2}\tau_s^{1/2} - \tau_s^{1/2}[\mathcal{V} + \hat{\xi}/2] \end{pmatrix} \tag{A.5}$$

The scalar potential will be:

$$\begin{aligned} V_F =& e^K \Big[(K^{T_b\overline{T_b}} K_{T_b} K_{\overline{T_b}} + K^{T_s\overline{T_s}} K_{T_s} K_{\overline{T_s}} + K^{T_s\overline{T_b}} K_{T_s} K_{\overline{T_b}} + K^{T_b\overline{T_s}} K_{T_b} K_{\overline{T_s}} - 3)|W|^2 + \\ & + (K^{T_s\overline{T_s}} K_{\overline{T_s}} + K^{T_s\overline{T_b}} K_{\overline{T_b}}) W_{T_s} \overline{W} + \\ & + (K^{T_s\overline{T_s}} K_{T_s} + K^{T_s\overline{T_b}} K_{T_b}) \overline{W}_{\overline{T_s}} W + \\ & + K^{T_s\overline{T_s}} \overline{W}_{\overline{T_s}} W_{T_s}\,]. \end{aligned} \tag{A.6}$$

The first bracket can be ease evaluated by the breaking of no scale structure:

$$K^{T_i\overline{T}_j} K_{T_i} K_{\overline{T}_j} = \frac{3}{\left[1 - \frac{\hat{\xi}}{4\mathcal{V}}\right]} = 3 + \frac{3}{4}\frac{\hat{\xi}}{\mathcal{V}} \tag{A.7}$$

and the second equal to:

$$\frac{-4\tau_s[\mathcal{V}+\hat{\xi}/2]}{[2\mathcal{V}-\hat{\xi}/2]} \sim -2\tau_s. \tag{A.8}$$

so in the large volume limit $\mathcal{V} \gg 1$ we will get:

$$V_F = \frac{3}{4}\frac{\hat{\xi}|W|^2}{\mathcal{V}^3} - \frac{2\tau_s}{\mathcal{V}^2}[W_{T_s}\overline{W} + \overline{W}_{\overline{T_s}}W] + \frac{8}{3}\frac{\tau_s^{1/2}}{\mathcal{V}}\overline{W}_{\overline{T_s}}W_{T_s}] \tag{A.9}$$

the dominant terms from the superpotential are:

$$\begin{aligned} &|W|^2 \sim W_o^2 \\ &[W_{T_s}\overline{W} + \overline{W}_{\overline{T_s}}W] \sim -2W_0 a_s A_s e^{-a_s T_s} \\ &\overline{W}_{\overline{T_s}}W_{T_s} \sim A_s^2 a_s^2 e^{-2a_s\tau_s} \end{aligned} \tag{A.10}$$

stabilizing on the vacuum the imaginary part of T_s such that $e^{-a_s Im(T_s)} = -1$ we get:

$$V_F = \frac{3}{4}\frac{\hat{\xi}W_0^2}{\mathcal{V}^3} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^2} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}} \tag{A.11}$$

looking for the minimum we get:

$$\begin{aligned} &\frac{\partial V_f}{\partial \mathcal{V}} = -\frac{9}{4}\frac{\hat{\xi}W_0^2}{\mathcal{V}^4} + 8\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^3} - \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}^2} = 0 \\ &\frac{9}{8}\hat{\xi}W_0^2 - 4W_0 A_s a_s e^{-a_s\tau_s}\tau_s\mathcal{V} + \frac{4}{3}A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}\mathcal{V}^2 = 0 \\ &\mathcal{V} = \frac{3}{2}\frac{W_0\tau_s^{1/2}}{A_s a_s e^{-a_s\tau_s}}\left(1 \pm \sqrt{1-\frac{3\hat{\xi}}{8\tau_s^{3/2}}}\right) \\ &\frac{\partial V_f}{\partial \tau_s} = 4\frac{W_0 a_s A_s}{\mathcal{V}^2}[e^{-a_s\tau_s} - a_s\tau_s e^{-a_s\tau_s}] - \frac{8}{3}\frac{A_s^2 a_s^2}{\mathcal{V}}\left[\frac{1}{2}\tau_s^{-1/2}e^{-2a_s\tau_s} - 2a_s\tau_s^{1/2}e^{-2a_s\tau_s}\right] = 0 \\ &\epsilon_s = 1/(4a_s\tau_s) \ll 1 \\ &2\left(1 \pm \sqrt{1-\frac{3\hat{\xi}}{8\tau_s^{3/2}}}\right) = \frac{1-4\epsilon_s}{1-\epsilon_s} \\ &\langle\mathcal{V}\rangle = \frac{3}{4}\frac{W_0\tau_s^{1/2}}{a_s A_s}e^{a_s\tau_s}\frac{1-4\epsilon_s}{1-\epsilon_s} \sim W_0 e^{a_s\tau_s} \\ &\langle\tau_s\rangle^{3/2} = \frac{1}{2}\hat{\xi}\frac{(1-\epsilon_s)^2}{1-4\epsilon_s} \end{aligned} \tag{A.12}$$

Looking now at the gravitino mass:

$$m_{3/2}^2 = e^K|W|^2 \sim \frac{W_0^2}{\langle\mathcal{V}\rangle^2} \sim e^{-2a_s\tau_s} \tag{A.13}$$

The energy on the minimum will be:

$$\langle V_F\rangle = -\frac{32}{3}\frac{\left(a_s A_s e^{-a_s\tau_s}\right)^3\epsilon_s}{W_0} + o(\epsilon_s^2) \sim -m_{3/2}^3 M_p. \tag{A.14}$$

In order to compute the mass of the modulus first we have to canonical normalize the fields:

$$\tau_b^c \sim \frac{\tau_b}{\langle\tau_b\rangle}\langle\tau_b\rangle \sim \langle\mathcal{V}\rangle^{2/3}\tau_s^c \sim \frac{\tau_s}{\langle\tau_b\rangle^{3/4}\langle\tau_s\rangle^{1/4}} \tag{A.15}$$

the scalar potential become:

$$V_F = \frac{3}{4}\frac{\hat{\xi}W_0^2}{\langle\tau_b\rangle^{9/2}\tau_b^{c9/2}} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s^c}\tau_s^c}{\langle\tau_b\rangle^3\tau_b^{c3}} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s^c}\tau_s^{c1/2}}{\langle\tau_b\rangle^{3/2}\tau_b^{c3/2}} \tag{A.16}$$

Asymptotically:

$$m_{ij}^2 \sim V_{ij} \Rightarrow m_{\tau_b^c}^2 \sim \langle\frac{\partial^2 V_F}{\partial\tau_b^{c2}}\rangle \sim \frac{W_0^2}{\langle\mathcal{V}\rangle^3} \sim m_{3/2}^3\frac{1}{M_p}. \tag{A.17}$$

Doing the same things for the mass of τ_s at the leading order we get

$$m_{\tau_s^c}^2 \sim \langle\frac{\partial^2 V_F}{\partial\tau_s^{c2}}\rangle \sim \frac{W_0}{\langle\mathcal{V}\rangle^2} \sim m_{3/2}^2. \tag{A.18}$$

Let us consider now the F-terms:

$$\begin{aligned} F^i &= e^{K/2}K^{i\bar{j}}D_{\bar{j}}\overline{W} \\ F^{T_b} &= \frac{-2\tau_b W_0}{\mathcal{V}} \sim m_{3/2}^{1/3}M_p^{5/3} \\ F^{T_s} &\simeq -\frac{2\tau_s W_0}{\mathcal{V}} + \frac{2\tau_s W_0}{\mathcal{V}}\frac{1-4\epsilon_s}{1-\epsilon_s} \simeq -\frac{6\tau_s W_0\epsilon_s}{\mathcal{V}} \sim m_{3/2}M_p \end{aligned} \tag{A.19}$$

Appendix B

Sequestered de Sitter String Model

The Kähler potential is:

$$K = -2\log\left(\mathcal{V} + \xi(S+\overline{S})^{3/2}\right) - \log(S+\overline{S}) + K_\alpha C^\alpha \overline{C}^{\overline{\alpha}} + K_{cs}(U) + \lambda_{SM}\frac{\left(\frac{T_{SM}+\overline{T}_{SM}}{2}\right)^2}{\mathcal{V}} + \\ + \lambda_b \frac{\left(\frac{G+\overline{G}}{2}\right)^2}{\mathcal{V}} + \lambda_{dS}\frac{\left(\frac{T_{dS}+\overline{T}_{dS}}{2}\right)^2}{\mathcal{V}} \tag{B.1}$$

with:

$$\mathcal{V} = \left(\frac{T_b+\overline{T_b}}{2}\right)^{3/2} - \left(\frac{T_s+\overline{T_s}}{2}\right)^{3/2} \tag{B.2}$$

and:

$$K_\alpha = \frac{f_\alpha(U,S)}{\mathcal{V}^{2/3}}\left(1 - c_s\frac{\xi(S+\overline{S})^{3/2}}{\mathcal{V}} + c_{SM}\left(\frac{T_{SM}+\overline{T}_{SM}}{2}\right)^p + c_b\left(\frac{G+\overline{G}}{2}\right)^p + K_{dS}\right) \tag{B.3}$$

the superpotential is:

$$W = W_0(U,S) + A_s(U,S)e^{-a_s T_s} + A_{dS}(U,S)e^{-a_{dS}(S+k_{dS}T_{dS})}. \tag{B.4}$$

The gauge kinetic function is defined:

$$f_a = \delta_a S + \kappa_a T_{SM}, \tag{B.5}$$

with δ_a are universal constants for $\mathbb{Z}_n$ singularities but can be non universal for general singularities. By the fact that $\langle C\rangle = \langle T_{SM}\rangle = \langle G\rangle = 0$ the only relevant terms for the Kähler derivatives are:

$$\begin{aligned} K_{T_b} &= -\frac{3}{2}\frac{\tau_b^{1/2}}{\mathcal{Y}} - \frac{3}{4}\frac{\lambda_{dS}\tau_{dS}^2\tau_b^{1/2}}{\mathcal{V}^2} \\ K_{T_s} &= \frac{3}{2}\frac{\tau_s^{1/2}}{\mathcal{Y}} + \frac{3}{4}\frac{\lambda_{dS}\tau_{dS}^2\tau_s^{1/2}}{\mathcal{V}^2} \\ K_{T_{dS}} &= \frac{\lambda_{dS}\tau_{dS}}{\mathcal{V}} \\ K_S &= -\frac{1}{2s}\left(1 + \frac{3}{2}\frac{\hat{\xi}}{\mathcal{Y}}\right) \end{aligned} \tag{B.6}$$

with $\tau_i = Re(T_i)$.
The inverse matrix will be:

$$
\begin{aligned}
K^{T_b\overline{T_b}} &= \frac{4}{3}\frac{\tau_b^{1/2}\mathcal{Y}^2}{(\mathcal{V}-\hat{\xi})}\left[1+\frac{2\tau_{dS}^2\lambda_{dS}\mathcal{Y}}{\mathcal{V}^2}\right]^{-1}\left[\left(1+\frac{2\tau_{dS}^2\lambda_{dS}\mathcal{Y}}{\mathcal{V}^2}\right)\left(1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}+\frac{9}{8}\frac{\hat{\xi}^2}{\mathcal{Y}^2}\right)+\right.\\
&\qquad\left.+\frac{3\tau_s^{3/2}}{\mathcal{Y}}\left(1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}\right)\right]\\
K^{T_s\overline{T_s}} &= -\frac{4}{3}\frac{\tau_s^{1/2}\mathcal{Y}^2}{(\mathcal{V}-\hat{\xi})}\left[1+\frac{2\tau_{dS}^2\lambda_{dS}\mathcal{Y}}{\mathcal{V}^2}\right]^{-1}\left[\left(1+\frac{2\tau_{dS}^2\lambda_{dS}\mathcal{Y}}{\mathcal{V}^2}\right)\left(1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}+\frac{9}{8}\frac{\hat{\xi}^2}{\mathcal{Y}^2}\right)+\right.\\
&\qquad\left.-\frac{3\tau_b^{3/2}}{\mathcal{Y}}\left(1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}\right)\right]\\
K^{T_s\overline{T_b}} &= \frac{4\tau_b\tau_s\mathcal{Y}}{(\mathcal{V}-\hat{\xi})}\left[1+\frac{2\tau_{dS}^2\lambda_{dS}\mathcal{Y}}{\mathcal{V}^2}\right]^{-1}\left[1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}\right]\\
K^{S\overline{S}} &= \frac{s^2}{(\mathcal{V}-\hat{\xi})}\left[(4\mathcal{V}-\hat{\xi})-\frac{3}{4}\frac{\lambda_{dS}\tau_{dS}}{\mathcal{V}}\mathcal{Y}\right]\\
K^{T_b\overline{S}} &= -\frac{3\hat{\xi}s\tau_b}{(\mathcal{V}-\hat{\xi})}\\
K^{T_s\overline{S}} &= -\frac{3\hat{\xi}s\tau_s}{(\mathcal{V}-\hat{\xi})}\\
K^{T_b\overline{T}_{dS}} &= \frac{\tau_{dS}\tau_b}{2\mathcal{V}(\mathcal{V}-\hat{\xi})}\left[1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}+\frac{9}{8}\frac{\hat{\xi}^2}{\mathcal{Y}^2}\right]\\
K^{T_s\overline{T}_{dS}} &= \frac{\tau_{dS}\tau_s}{2\mathcal{V}(\mathcal{V}-\hat{\xi})}\left[1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}+\frac{9}{8}\frac{\hat{\xi}^2}{\mathcal{Y}^2}\right]\\
K^{T_{dS}\overline{S}} &= -\frac{9\hat{\xi}s\tau_{dS}}{2(\mathcal{V}-\hat{\xi})}\\
K^{T_{dS}\overline{T}_{dS}} &= \frac{2\mathcal{V}}{\lambda_{dS}(\mathcal{V}-\hat{\xi})}\left[(\mathcal{V}-\hat{\xi})+\frac{\lambda_{dS}\tau_{dS}\mathcal{Y}^2}{2\mathcal{V}^2}\left(1-\frac{3}{4}\frac{\hat{\xi}}{\mathcal{Y}}+\frac{9}{8}\frac{\hat{\xi}^2}{\mathcal{Y}^2}\right)\right]
\end{aligned}
\tag{B.7}
$$

All the other terms can be neglected.

F-term stabilisation

The F-term potential is given by:

$$
V_F = e^K\left(K^{i\bar{j}}D_iWD_{\bar{j}}\overline{W}-3|W|^2\right). \tag{B.8}
$$

Due to the spoil of the no scale structure induced by the ξ corrections we will get:

$$
K_iK^{i\bar{j}}K_{\bar{j}} = 3+\frac{3}{4}\frac{\hat{\xi}}{\mathcal{V}} \tag{B.9}
$$

The coefficient of the $|W|^2$ term is at the leading order:

$$\frac{3}{4}\frac{\hat{\xi}}{2s\mathcal{V}^3} \tag{B.10}$$

the coefficient of $\overline{W}_{\overline{T}_S}W$:

$$\frac{-2\tau_s}{2s\mathcal{V}^2}. \tag{B.11}$$

The scalar potential will be:

$$V_F = \frac{3}{4}\frac{\hat{\xi}|W|^2}{2s\mathcal{V}^3} - \frac{2\tau_s}{2s\mathcal{V}^2}[W_{T_s}\overline{W} + \overline{W}_{\overline{T}_s}] + \frac{8}{3}\frac{\tau_s^{1/2}}{2s\mathcal{V}} \tag{B.12}$$

the dominant terms from the superpotential are:

$$\begin{aligned}
&|W|^2 \sim W_o^2 \\
&[W_{T_s}\overline{W} + \overline{W}_{\overline{T_s}}W] \sim -2W_0 a_s A_s e^{-a_s T_s} \\
&\overline{W}_{\overline{T_s}}W_{T_s} \sim W_0^2 A_S^2 a_s^2 e^{-2a_s\tau_s}
\end{aligned} \tag{B.13}$$

stabilizing on the vacuum the imaginary part of T_s such that $e^{-a_s Im(T_s)} = -1$ we get:

$$V_F = \frac{1}{2s}\left[\frac{3\hat{\xi}W_0^2}{4\mathcal{V}^3} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^2} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}}\right] \tag{B.14}$$

Considering now also the T_{dS} modulus we get the F-temr potential (fixing $e^{-i2a_{dS}\varphi_{dS}} = -1$) we get:

$$V_F^{dS} = \frac{(\kappa_{dS}a_{dS}A_{dS})^2}{s\lambda_{dS}}\frac{e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})}}{\mathcal{V}} \tag{B.15}$$

The total potential will be:

$$\begin{aligned}
V^{Tot} =& \frac{1}{2s}\left[\frac{3\hat{\xi}W_0^2}{4\mathcal{V}^3} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^2} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}}\right] + \\
&+ \frac{(\kappa_{dS}a_{dS}A_{dS})^2}{s\lambda_{dS}}\frac{e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})}}{\mathcal{V}} + \frac{1}{2s}\left(\sum_i q_{h,i}|\phi_{h,i}|^2 + \frac{q_{dS}}{4\pi}\frac{\lambda_{dS}\tau_{dS}}{\mathcal{V}}\right)^2
\end{aligned} \tag{B.16}$$

Minimisation of the potential

Stabilising the D-term in the de Sitter potential we fix $\langle q_{h,i}|\varphi_{h,i}|^2\rangle = 0$ so the total potential will be:

$$\begin{aligned}
V^{Tot} =& \frac{1}{2s}\left[\frac{3\hat{\xi}W_0^2}{4\mathcal{V}^3} - 4\frac{W_0 A_s a_s e^{-a_s\tau_s}\tau_s}{\mathcal{V}^2} + \frac{8}{3}\frac{A_s^2 a_s^2 e^{-2a_s\tau_s}\tau_s^{1/2}}{\mathcal{V}}\right] + \\
&+ \frac{(\kappa_{dS}a_{dS}A_{dS})^2}{s\lambda_{dS}}\frac{e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})}}{\mathcal{V}} + \frac{1}{2s}\left(\frac{q_{dS}}{4\pi}\frac{\lambda_{dS}\tau_{dS}}{\mathcal{V}}\right)^2
\end{aligned} \tag{B.17}$$

$$\frac{\partial V^{Tot}}{\partial \tau_s} = 0 \rightarrow \frac{2a_s A_s e^{-a_s\tau_s}\mathcal{V}}{3W_0\tau_s^{1/2}} = \frac{1-a_s\tau_s}{\frac{1}{2}-2a_s\tau_s}$$
$$\frac{\partial V^{Tot}}{\partial \tau_{dS}} = 0 \rightarrow (\kappa_{dS}a_{dS}A_{dS})^2 e^{-2a_{dS}(s+\kappa_{dS}\tau_{dS})} = \frac{(q_{dS}\lambda_{dS})^2\tau_{dS}}{(4\pi)^2\mathcal{V}}$$
$$\frac{\partial V^{Tot}}{\partial \mathcal{V}} = 0 \rightarrow \mathcal{V}_\pm = \frac{3W_0\tau_s^{1/2}}{2a_sA_se^{-a_s\tau_s}}\left(1+\frac{q_{dS}^2\lambda_{dS}\tau_{dS}(1+\lambda_{dS}\tau_{dS})}{(4\pi)^2 4W_0a_sA_se^{-a_s\tau_s}\tau_s}\right)$$
$$\left[1\pm\sqrt{1-\frac{3\hat{\xi}}{8\tau_s^{3/2}\left(1+\frac{q_{dS}^2\lambda_{dS}\tau_{dS}(1+\lambda_{dS}\tau_{dS})}{(4\pi)^2 4W_0a_sA_se^{-a_s\tau_s}\tau_s}\right)}}\right] =$$
$$= \frac{3W_0\tau_s^{1/2}}{2a_sA_se^{-a_s\tau_s}}\left(1+\frac{\alpha}{a_s\tau_s}\right)\left[1\pm\sqrt{1-\frac{3\hat{\xi}}{8\tau_s^{3/2}\left(1+\frac{\alpha}{a_s\tau_s}\right)}}\right] \tag{B.18}$$

Defining $\epsilon_s = 1/4a_s\tau_s \ll 1$ we get:

$$\langle\tau_s\rangle^{3/2} = \frac{3}{2}\hat{\xi}\frac{(1-\epsilon_s)^2}{(1-4\epsilon_s)[(4-4\epsilon_s)(1+4\alpha\epsilon_s)-(1-4\epsilon_s)]} =$$
$$= \frac{\hat{\xi}}{2}\left[1+\left(\frac{3}{4}-\frac{4}{3}\alpha\right)4\epsilon_s+\left(\frac{3}{4}-\frac{2}{3}\alpha+\frac{16}{3}\alpha^2\right)16\epsilon_s^2+o(\epsilon_s^3)\right]$$
$$\langle\mathcal{V}\rangle = \frac{3W_0\langle\tau_s\rangle^{1/2}}{4a_sA_s}e^{a_s\langle\tau_s\rangle}\frac{1-4\epsilon_s}{1-\epsilon_s} = \frac{3W_0\langle\tau_s\rangle^{1/2}}{4a_sA_s}e^{a_s\langle\tau_s\rangle}+o(\epsilon_s)$$
$$\langle s\rangle = \frac{1}{g_s} \tag{B.19}$$

Shift of the Dilaton and Complex Structure Minimum

The α' correction introduce a shift in the dilaton minimum that is:

$$D_SW \simeq D_SW_{flux}|_{\hat{\xi}=0} + W_{n.p.}K_S + W_{n.p.S} + W_0K_S^{\hat{\xi}} =$$
$$= -\frac{3}{4}\frac{W_0\hat{\xi}}{s\mathcal{V}}\left[1-\epsilon_s+\frac{s\epsilon_s}{2}\partial_S\log A_s(U,S)+\right.$$
$$\left.+\frac{\lambda_{dS}^{1/2}\epsilon_s^{1/2}}{\kappa_{dS}\hat{\xi}^{1/2}}\left(\frac{\sqrt{2}s}{a_{dS}}\partial_S\log A_{dS}(U,S)-\frac{1}{\sqrt{2}a_{dS}}-\sqrt{2}s\right)\right. =$$
$$= -\frac{3}{4}\frac{W_0\hat{\xi}\omega_S(U,S)}{s\mathcal{V}} \tag{B.20}$$

The complex structure moduli are stabilised supersymmetrically for $\hat{\xi}$ but introducing the α' correction the minimum is shifted producing:

$$D_{U_i}W \simeq D_{U_i}W_{flux} + W_{n.p.}K_{c.s.U_i} + W_{n.p.U_i} =$$
$$= -\frac{3}{2}\frac{W_0\hat{\xi}\epsilon_s}{\mathcal{V}}\partial_{U_i}\left[K_{c.s.}+\log A_s(U,S)+\frac{\lambda_{dS}^{1/2}}{\kappa_{dS}a_{dS}\epsilon_s^{1/2}\hat{\xi}^{1/2}}\left[K_{c.s}+\log A_{dS}(U,S)\right]\right] =$$
$$= -\frac{3}{4}\frac{W_0\hat{\xi}\omega_{U_i}(U,S)}{s\mathcal{V}} \tag{B.21}$$

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