PROCEEDINGS OF SCIENCE

On nonsupersymmetric Pati–Salam string models

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We report recent progress in the construction of heterotic string compactifications with spontaneously broken supersymmetry and Pati-Salam gauge symmetry. We study one-loop radiative corrections to the string effective potential and comment on its structure, along with the conditions allowing for an exponentially suppressed value of the cosmological constant.

Corfu Summer Institute 2021 "School and Workshops on Elementary Particle Physics and Gravity"
29 August - 9 October 2021
Corfu, Greece

*Speaker

PoS(CORFU2021)061

1. Introduction

An important open problem in string phenomenology concerns the mechanism of spontaneous supersymmetry breaking and its physical consequences which has attracted considerable interest in recent literature[1–18]. To this end, a tree-level analysis is insufficient for a quantitative comparison with low energy data, and the incorporation of quantum corrections to terms in the string effective action becomes a necessary yet daunting task. It is known that several stability issues need to be addressed. On the one hand, the absence of supersymmetry is accompanied by potential Hagedorn-like instabilities associated with the potential excitation of tachyonic modes. On the other hand, terms in the string effective action are no longer super-protected and, in particular, the effective potential receives corrections already at one-loop, triggering a sizeable back-reaction to the tree level geometry. Furthermore, the one-loop effective potential typically becomes a function of all tree-level moduli and the analysis of the potential with respect to all of them, as well as the identification of a moduli stabilisation mechanism, remain difficult open problems.

It is then clear that, in order to retain string effects, the spontaneous breaking of supersymmetry must be carried out at the full string level, in terms of an exactly solvable worldsheet CFT. In the context of heterotic theories, this is the case in coordinate-dependent compactifications [19–22] which essentially provide the stringy realisation of the Scherk-Schwarz mechanism [23, 24]. In the simplest case, the breaking can be formulated in terms of freely-acting $\mathbb{Z}_2$ orbifolds coupling the spacetime fermion parity $(-1)^F$ to an order-2 shift δ along a non-trivial cycle of the compactification manifold. In the case where the latter is identified with a circle S^1 of radius R, the Scherk-Schwarz mechanism introduces a mass gap for states charged under F and results in the spontaneous breaking of supersymmetry, with the breaking scale essentially controlled by the KK scale, $m_{3/2} \sim M_{\rm KK} \sim 1/R$.

In this report, we discuss a class of four-dimensional $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold compactifications with Pati–Salam gauge symmetry [25, 26], where $\mathcal{N} = 1$ supersymmetry is spontaneously broken à la Scherk-Schwarz, and which contain a number of semi-realistic features, such as the presence of chiral matter. The models are initially constructed using the free-fermionic formulation [27–29] and are then subsequently mapped to equivalent orbifold compactifications, where the dependence on the moduli is manifest (*c.f.* [30] and references therein).

The primary scope of our analysis is to study the one-loop effective potential and to identify its universal features. Clearly, this is a function of all moduli and the analysis of its shape with respect to all of them lies outside the scope of this study. Furthermore, although the theories we construct are tachyon-free at the fermionic point, small deformations along generic directions around it are certain to excite tachyonic modes, provided the Scherk-Schwarz radius R is sufficiently close to the self-dual (fermionic) point. However, as soon as R grows sufficiently larger than the string scale, the theory becomes protected against tachyonic instabilities, regardless of the values of the remaining moduli.

This report is organised as follows. In Section 2 we describe the construction of the models at a special point of moduli space, where all internal coordinates can be consistently fermionised. We then proceed in Section 3 to recast these theories as orbifold compactifications at generic points of the perturbative moduli space and investigate the shape of the one-loop effective potential as a function of the Scherk–Schwarz volume modulus. We discuss the classification and universality features of the models with respect to the super no-scale condition and the resulting features of the one-loop effective potential.

2. Nonsupersymmetric Pati–Salam models

The starting point of the model analysis is a class of heterotic Pati–Salam string models defined in the framework of the free fermionic formulation [27–29] by a set of $n = 10$ basis vectors $\{v_1, \ldots, v_n\}$, where

$$\begin{aligned}
v_1 = \mathbb{1} &= \{\psi^\mu,\ \chi^{1,\ldots,6}, y^{1,\ldots,6}, \omega^{1,\ldots,6} | \bar{y}^{1,\ldots,6}, \bar{\omega}^{1,\ldots,6}, \bar{\eta}^{1,2,3}, \bar{\psi}^{1,\ldots,5}, \bar{\phi}^{1,\ldots,8}\}\\
v_2 = S &= \{\psi^\mu, \chi^{1,\ldots,6}\}\\
v_3 = T_1 &= \{y^{12}, \omega^{12} | \bar{y}^{12}, \bar{\omega}^{12}\}\\
v_4 = T_2 &= \{y^{34}, \omega^{34} | \bar{y}^{34}, \bar{\omega}^{34}\}\\
v_5 = T_3 &= \{y^{56}, \omega^{56} | \bar{y}^{56}, \bar{\omega}^{56}\}\\
v_6 = b_1 &= \{\chi^{34}, \chi^{56}, y^{34}, y^{56} | \bar{y}^{34}, \bar{y}^{56}, \bar{\psi}^{1,\ldots,5}, \bar{\eta}^{1}\}\\
v_7 = b_2 &= \{\chi^{12}, \chi^{56}, y^{12}, y^{56} | \bar{y}^{12}, \bar{y}^{56}, \bar{\psi}^{1,\ldots,5}, \bar{\eta}^{2}\}\\
v_8 = z_1 &= \{\bar{\phi}^{1,\ldots,4}\}\\
v_9 = z_2 &= \{\bar{\phi}^{5,\ldots,8}\}\\
v_{10} = \alpha &= \{\bar{\psi}^{4,5}, \bar{\phi}^{1,2}\},
\end{aligned} \tag{1}$$

along with a set of phases

$$c_{ij} = c\begin{bmatrix} v_i \\ v_j \end{bmatrix} = \pm 1\ ,\ i \geq j = 1, \ldots, 10\ , \tag{2}$$

associated with generalised GSO (GGSO) projections. These give rise to a huge number of $2^{\frac{10(10-1)}{2}+1} \sim 7\times 10^{13}$ in principle distinct configurations with Pati–Salam gauge group. The full gauge symmetry of these configurations, aside from special cases where gauge symmetry enhancement occurs, is:

$$\begin{aligned}
G &= SO(6) \times SO(4) \times U(1)^3 \times SO(4)^2 \times SO(8)\\
&= SU(4) \times SU(2)_L \times SU(2)_R \times U(1)^3 \times SU(2)^4 \times SO(8)\ .
\end{aligned} \tag{3}$$

We focus our analysis on a non superymmetric subset of these vacua obtained by setting $c[{S \atop T_1}] = +1$. These vacua are then subject to a set of constraints regarding their phenomenological characteristics which can be expressed purely in terms of the GGSO phases. By employing projection and representation operators, the entire string spectrum can be encoded in a model-independent form. The generalised projectors for a given sector a are:

$$\mathbb{P}_a^{\pm} = \prod_{\xi \in \Xi^{\pm}(a)} \frac{1}{2}\left(1 \pm c\begin{bmatrix} a \\ \xi \end{bmatrix}^*\right)\ , \tag{4}$$

for states constructed from the ground state of the sector a and

$$\mathbb{P}_a^{\varphi} = \prod_{\xi \in \Xi(a)} \frac{1}{2}\left(1 + \delta_a \delta_\xi^{\varphi} c\begin{bmatrix} a \\ \xi \end{bmatrix}^*\right)\ , \tag{5}$$

with

$$\delta_a^{\varphi} = \begin{cases} -1\ ,\ \varphi \in a \\ +1\ ,\ \varphi \notin a \end{cases} \quad , \tag{6}$$

for states containing fermion oscillators φ . Here $\Xi^{\pm}(a)$ is the set of basis vectors which act as projection operators on the sector a in question.

The first condition we impose concerns the appearance of tachyonic states in the string spectrum, which is no longer protected from such instabilities by supersymmetry. The elimination of all physical tachyonic states from the spectrum then amounts to demanding that all relevant projectors in the tachyonic sectors vanish:

$$\begin{aligned} &\mathbb{P}_a^+ = 0\ ,\ a \in \{z_1, z_2, \alpha, z_1+\alpha\} \cup \{T_i + z_1, T_i + z_2, T_i + p z_1 + \alpha\}\ , \\ &\qquad\qquad\qquad\qquad\qquad\qquad i = 1,2,3,\ p = 0,1\ , \\ &\mathbb{P}_{T_j}^{\bar\phi} = 0\ ,\ \bar\phi \in \mathbb{1}_R - T_{j_R}\ ,\ j = 1,2,3\ . \end{aligned} \tag{7}$$

Having dealt with the tachyons, we can now proceed with the analysis of the massless spectrum. We introduce additional conditions with the construction of semi-realistic Pati–Salam models in mind. First of all, we require that the sectors $\mathcal{S}^i_{pq} = S{+}b_i{+}pT_j{+}qT_k$, where $p,q = 0,1$, $(i,j,k) = (1,2,3)$, $i \neq j \neq k$ give rise to chiral fermionic matter in complete Pati–Salam generations.

The number of generations is given by:

$$n_g = n_L - \bar n_L = \bar n_R - n_R\ , \tag{8}$$

where n_L, $\bar n_R$, count the number of copies of the representations $F_L(\mathbf{4},\mathbf{2},\mathbf{1})$ and $\overline{F}_R(\bar{\mathbf{4}},\mathbf{1},\mathbf{2})$ containing fermion generations while $\bar n_L$, n_R count the corresponding anti-generations in $\overline{F}_L(\bar{\mathbf{4}},\mathbf{2},\mathbf{1})$ and $F_R(\mathbf{4},\mathbf{1},\mathbf{2})$,

$$\begin{aligned} n_L &= 4\sum_{i=1}^{3}\sum_{p,q=0}^{1} \mathbb{P}^-_{\mathcal{S}^i_{pq}} \frac{1}{2}\left(1 + X^{SU(4)}_{\mathcal{S}^i_{pq}}\right)\frac{1}{2}\left(1 + X^{SO(4)}_{\mathcal{S}^i_{pq}}\right)\ , \\ \bar n_R &= 4\sum_{i=1}^{3}\sum_{p,q=0}^{1} \mathbb{P}^-_{\mathcal{S}^i_{pq}} \frac{1}{2}\left(1 - X^{SU(4)}_{\mathcal{S}^i_{pq}}\right)\frac{1}{2}\left(1 - X^{SO(4)}_{\mathcal{S}^i_{pq}}\right)\ , \\ \bar n_L &= 4\sum_{i=1}^{3}\sum_{p,q=0}^{1} \mathbb{P}^-_{\mathcal{S}^i_{pq}} \frac{1}{2}\left(1 - X^{SU(4)}_{\mathcal{S}^i_{pq}}\right)\frac{1}{2}\left(1 + X^{SO(4)}_{\mathcal{S}^i_{pq}}\right)\ , \\ n_R &= 4\sum_{i=1}^{3}\sum_{p,q=0}^{1} \mathbb{P}^-_{\mathcal{S}^i_{pq}} \frac{1}{2}\left(1 + X^{SU(4)}_{\mathcal{S}^i_{pq}}\right)\frac{1}{2}\left(1 - X^{SO(4)}_{\mathcal{S}^i_{pq}}\right)\ . \end{aligned} \tag{9}$$

Here we have introduced the representation operators

$$X^{SU(4)}_{\mathcal{S}^i_{pq}} = \begin{cases} -c\begin{bmatrix} \mathcal{S}^i_{pq} \\ \mathcal{S}^j_{0,1-q}{+}\alpha \end{bmatrix}^*, & j \neq i = 1,2 \\ -c\begin{bmatrix} \mathcal{S}^i_{pq} \\ \mathcal{S}^1_{1-q,0}{+}\alpha \end{bmatrix}^*, & i = 3 \end{cases} \quad , \quad X^{SO(4)}_{\mathcal{S}^i_{pq}} = -c\begin{bmatrix} \mathcal{S}^i_{pq} \\ \alpha \end{bmatrix}^*, i = 1,2,3\ .$$

We note that due to the complexification of the internal fermions: $y^{12,34,56}$, $\omega^{12,34,56}$, $\bar y^{12,34,56}$, $\bar\omega^{12,34,56}$, physical states of these models exhibit a multiplicity of 4, preventing the construction of models with 3 generations within the present setup. This issue is expected to be resolved in constructions with real fermions.

Similar GGSO related conditions can be imposed which ensure the presence of the heavy and light Higgs bosons, $H\,(\mathbf{4},\mathbf{1},\mathbf{2})$ and $h\,(\mathbf{1},\mathbf{2},\mathbf{2}) = H_u\left(\mathbf{1},\mathbf{2},+\frac{1}{2}\right) + H_d\left(\mathbf{1},\mathbf{2},-\frac{1}{2}\right)$, responsible for the spontaneous breaking of the Pati–Salam and SM symmetries at the effective field theory level.

A common feature in the spectra of heterotic compactifications such as the ones presented here are exotic fractionally charged states [31–35]. Since no experimental evidence for such particles exists, we will demand that such states, always appear in vector-like pairs and can therefore acquire large masses. We impose this relaxed condition as opposed to the more strict demand that all exotic vanish, since the appearance of such states is inevitable in models with complex internal fermions.

Finally, we discard all models which exhibit enhancements of the observable Pati–Salam gauge symmetry, while allowing for the enhancement of the hidden sector gauge groups.

After examining all these constraints, we find that 11 GGSO phases are irrelevant to our analysis. In fact, $c\left[\begin{smallmatrix}\mathbb{1}\\ \mathbb{1}\end{smallmatrix}\right]$ and $c\left[\begin{smallmatrix}\mathbb{1}\\ S\end{smallmatrix}\right]$ amount to conventions, while $c\left[\begin{smallmatrix}S\\ b_1\end{smallmatrix}\right]$, $c\left[\begin{smallmatrix}S\\ b_2\end{smallmatrix}\right]$ and $c\left[\begin{smallmatrix}b_1\\ b_2\end{smallmatrix}\right]$ correspond to a flip in the overall chirality and can therefore be fixed without loss of generality. Additionally, the phases $c\left[\begin{smallmatrix}\mathbb{1}\\ b_{1,2}\end{smallmatrix}\right]$, $c\left[\begin{smallmatrix}\mathbb{1}\\ z_{1,2}\end{smallmatrix}\right]$, as well as $c\left[\begin{smallmatrix}T_1\\ b_2\end{smallmatrix}\right]$ and $c\left[\begin{smallmatrix}T_2\\ b_1\end{smallmatrix}\right]$ are not necessary for the implementation of the above constraints. In order to proceed with a detailed analysis of the model parameter space, we perform a two-stage computer-aided scan, taking advantage of the fact that some critical model properties can be determined at the $SO(10)$ level generated by the vectors v_i, $i = 1, \ldots, 9$ [36, 37]. For each $SO(10)$-level model which exhibits promising phenomenology we then re-introduce and iterate over the full set GGSO phases related to the $SO(10)$-breaking vector α and perform a comprehensive scan of all descendant Pati–Salam models in order to find those that meet all the aforementioned phenomenological criteria. Using this method we have collected a sample of approximately 5×10^4 Pati–Salam models which fulfill all the constraints.

3. Moluli dependence of one-loop effective potentials

While the fermionic formulation allows for a full model-independent analysis of the string spectrum, it can only offer insight on the dynamics at the “fermionic point” of moduli space, where all tori describing the compactified space are consistently fermionized. In order to investigate the properties of our models at more general points of the perturbative moduli space, we employ the orbifold formulation, in which the one-loop partition function corresponding to models generated by the basis (1) can be cast in the from:

$$
\begin{aligned}
Z = & \frac{1}{\eta^2\bar\eta^2}\frac{1}{2^4}\sum_{\substack{h_1,h_2,H,H'\\ g_1,g_2,G,G'}\in\mathbb{Z}_2}\frac{1}{2^3}\sum_{\substack{a,k,\rho\\ b,\ell,\sigma}\in\mathbb{Z}_2}\frac{1}{2^3}\sum_{\substack{H_1,H_2,H_3\\ G_1,G_2,G_3}\in\mathbb{Z}_2}(-1)^{a+b+HG+H'G'+\Phi}\\
&\times\frac{\vartheta[^a_b]}{\eta}\frac{\vartheta[^{a+h_1}_{b+g_1}]}{\eta}\frac{\vartheta[^{a+h_2}_{b+g_2}]}{\eta}\frac{\vartheta[^{a-h_1-h_2}_{b-g_1-g_2}]}{\eta}\\
&\times\frac{\bar\vartheta[^k_\ell]^3}{\bar\eta^3}\frac{\bar\vartheta[^{k+H'}_{\ell+G'}]}{\bar\eta}\frac{\bar\vartheta[^{k-H'}_{\ell-G'}]}{\bar\eta}\frac{\bar\vartheta[^{k+h_1}_{\ell+g_1}]}{\bar\eta}\frac{\bar\vartheta[^{k+h_2}_{\ell+g_2}]}{\bar\eta}\frac{\bar\vartheta[^{k-h_1-h_2}_{\ell-g_1-g_2}]}{\bar\eta}\\
&\times\frac{\bar\vartheta[^{\rho+H'}_{\sigma+G'}]}{\bar\eta}\frac{\bar\vartheta[^{\rho-H'}_{\sigma-G'}]}{\bar\eta}\frac{\bar\vartheta[^\rho_\sigma]^2}{\bar\eta^2}\frac{\bar\vartheta[^{\rho+H}_{\sigma+G}]^4}{\bar\eta^4}\\
&\times\frac{\Gamma^{(1)}_{2,2}[^{H_1}_{G_1}|^{h_1}_{g_1}](T^{(1)},U^{(1)})}{\eta^2\bar\eta^2}\frac{\Gamma^{(2)}_{2,2}[^{H_2}_{G_2}|^{h_2}_{g_2}](T^{(2)},U^{(2)})}{\eta^2\bar\eta^2}\frac{\Gamma^{(3)}_{2,2}[^{H_3}_{G_3}|^{h_1+h_2}_{g_1+g_2}](T^{(3)},U^{(3)})}{\eta^2\bar\eta^2}\,,
\end{aligned}
\tag{10}
$$

where Φ is a phase that encodes information related to the GGSO phases (2). In this notation, the Neveau–Schwarz and Ramond sectors are defined by the parameter $a = 0, 1$ respectively, with summation over b acting as the GSO projection. The twisted sectors of the $\mathbb{Z}_2\times\mathbb{Z}_2$ orbifold are determined by $h_{1,2}$ and the projections $g_{1,2}$. The complex right-moving fermions generating the $SO(16)\times SO(16)$ lattice are labeled by k and ρ and their corresponding projections ℓ and σ. The twists utilised in order to break the gauge symmetry and produce the observable Pati–Salam factor are then introduced via (H, G) and (H', G').

Finally, the moduli dependence is manifest through the two dimensional toroidal lattices $\Gamma_{2,2}$, on which the order two shifts corresponding to three additional freely-acting $\mathbb{Z}_2$ orbifolds labeled by (H_i, G_i) are introduced. At the fermionic point defined as $T = i$, $U = (1 + i)/2$, these lattices factorise into products of Jacobi theta functions:

$$
\Gamma_{2,2}[^{H_i}_{G_i}|^{h_i}_{g_i}] = \frac{1}{2}\sum_{\epsilon_i,\zeta_i\in\mathbb{Z}_2}|\vartheta[^{\epsilon_i}_{\zeta_i}]\vartheta[^{\epsilon_i+h_i}_{\zeta_i+g_i}]|^2(-1)^{H_i\zeta_i+G_i\epsilon_i+H_iG_i}\,,
\tag{11}
$$

recovering the fermionic formulation description.

All model dependent information is then carried by the orbifold phase Φ, containing bilinears of the boundary condition parameters, with specific choices corresponding to each possible model. The implementation of the Scherk–Schwarz mechanism [23, 24] imposes strict constraints on the form of this phase [26, 30].

This reformulation of the theory provides the tools necessary for the investigation of the one-loop effective potential

$$
V_{\text{one-loop}}(t_I) = -\frac{1}{2(2\pi)^4}\int_{\mathcal{F}}\frac{d^2\tau}{\tau_2^3}Z(\tau,\bar\tau;t_I),
\tag{12}
$$

as a function of the moduli of the tree two-tori $t_I = \left\{T_1^{(i)} + iT_2^{(i)}, U_1^{(i)} + iU_2^{(i)}, i = 1, 2, 3\right\}$. In the absence of supersymmetry, this potential receives non-vanishing contributions from the entire tower of bosonic and fermionic states in the spectrum. This potential typically leads to large, negative values for the cosmological constant and the introduction of a dilaton tadpole. In order to probe the behaviour of the potential away from the fermionic point, we consider deformations along the volume modulus T_2 of the Scherk-Schwarz torus while keeping the remaining moduli fixed at the fermionic point. It turns out that the large volume dynamics can be summarised as follows

$$V_{\text{one-loop}}(T_2 \gg 1) \sim -C\frac{n_b - n_f}{T_2^2} + \text{ exponentially suppressed terms }, \tag{13}$$

where C is a normalization constant, while $n_{b,f}$ stand for the number of the massless bosonic and fermionic degrees of freedom of the theory, with all other contributions being suppressed exponentially. The cosmological constant we obtain in that case is well above the observational upper bound even for values of T_2 that lie in the region of a few TeV, necessitating the elimination of the $1/T_2^2$ term by demanding that $n_b = n_f$. Models which exhibit this massless Bose–Fermi degeneracy, termed "super no-scale models" [1, 38–40] are capable of producing an exponentially suppressed cosmological constant[1].

The super no-scale condition $n_b = n_f$, along with the demand that supersymmetry breaking be consistent with the Scherk–Schwarz mechanism are then imposed as additional constraints on top of those presented in the previous section. By investigating models satisfying all of the above criteria, we notice that they can be classified according to their effective potential, although their actual spectra and interactions may differ.

Explicit calculations of the one-loop effective potential as a function of the T_2 modulus reveal the existence of four broad categories of models based on the general structure of the potential. Four typical examples are presented in Figure 1. The first category consists of models (see Fig. 1A) with a positive semi-definite potential exhibiting a global maximum at the fermionic point and an exponential suppression driving the theory towards a small positive cosmological constant as T_2 is varied. The second category (see Fig. 1B) also possesses a positive semi-definite potential, however, here the self-dual point corresponds to a local minimum, accompanied by a pair of maxima by its sides. At large volume this case also leads to small positive values for the cosmological constant. The two remaining categories (see Fig. 1C,D) correspond to potentials with a global minimum with large negative value at the fermionic point. In the third category, the theory exhibits two positive maxima, while the potential of fourth category is negative semi-definite.

While the cases with a negative potential may at first appear problematic, since they imply supersymmetry breaking at the string scale and the stabilization of the cosmological constant at large negative values, we note that this is not necessarily provide grounds for discarding them. The reason is that the super no-scale property and the shape of the effective potential are highly sensitive to all degrees of freedom in the string spectrum and are not preserved when modding out by additional orbifold group factors. This was explicitly checked by analyzing an $SO(10)$ parent-model and its Pati–Salam descendants and we indeed confirmed that, as expected, the introduction of additional $\mathbb{Z}_2$ orbifold factors drastically affects those properties. Thus, the negative-potential solutions constructed here may still be relevant at an intermediate stage, as the new parent models

[1]Note that, although the super no-scale condition $n_b = n_f$ appears universal, this is not necessarily the case. Actually, the form of the dominant power law suppression of the effective one-loop potential also depends on the particular embedding of the Scherk-Schwarz shift vector in the 2-torus. Different choices for this lattice vector typically lead to the T^{-2} term dressed by different automorphic functions of the U-modulus, with potentially different combinatorics factors. In principle, this implies that the super no-scale condition would then need to be suitably generalized in order to cancel the power law term. We defer further discussion of this point to future work.

that may produce vacua with appealing phenomenological features, upon introducing additional orbifolds factors to further reduce to the Standard Model gauge symmetry.

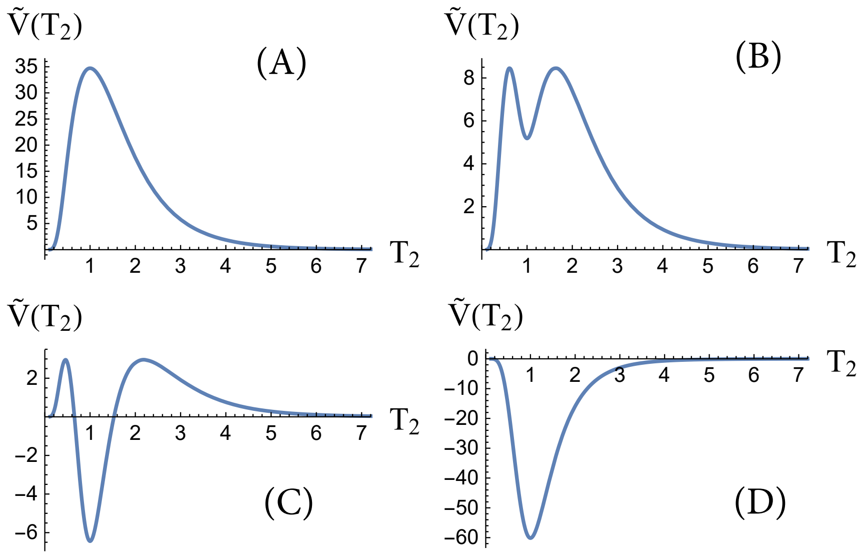

Figure 1: An example for each of the four distinct categories of models based on the overall shape of the rescaled potential $\tilde{V}(T_2) = 2(2\pi)^4 V(T_2)$.

Acknowledgments

The research of K.V. is co-financed by Greece and the European Union (European Social Fund - ESF) through the Operational Programme "Human Resources Development, Education and Lifelong Learning" in the context of the project "Strengthening Human Resources Research Potential via Doctorate Research" (MIS-5000432), implemented by the State Scholarships Foundation (IKY).

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