

# Counter-intuition and Scalar Masses

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## **Abstract**

The Bousso-Polchinski mechanism for discretely fine-tuning the cosmological constant favors a large bare negative cosmological constant. I argue (using generalizations of results of Klebanov and Tseytlin) that a similar mechanism for fine-tuning the scalar masses to small values favors a large bare negative (i.e. tachyonic) scalar mass. I comment briefly on the related issue of the role of low energy supersymmetry in string theory.

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## 1. The c.c. and $m_T$

In [1], Bousso and Polchinski proposed that in the presence of multiple flux quantum numbers arising from form fields on compactification geometries, the cosmological constant can be finely tuned to a value near zero (see also [2]). They modelled this mechanism in a very simple way, via a cancellation between a bare negative cosmological constant and positive contributions from form field kinetic terms integrated over a compactification:

$$\Lambda_{BP} = -|\lambda_{bare}| + \sum_i c_i Q_i^2. \quad (1.1)$$

Here  $i$  labels cycles on the compactification geometry, which carry flux quantum numbers  $Q_i$  and contribute to the cosmological term via the form field kinetic terms.  $\lambda_{bare}$  contains the rest of the contributions to the cosmological constant, and the mechanism applies when this is negative. In [1], the coefficients  $c_i$  (with incommensurate values of  $\sqrt{c_i}$ ) were taken to be constants sufficiently smaller than one to have control over the  $\alpha'$  expansion.<sup>1</sup> As explained in [1], a lattice of fluxes of dimensionality  $b \ll Q$  yields of order  $Q^{b-1}$  vacua in a shell of radius  $Q$  and unit thickness in flux space. As a result, with a multiplicity of fluxes one can discretely tune the total cosmological constant  $\Lambda_{BP}$  to a small value of order  $Q^{-b}$  in Planck units. It is then evident that within the framework of this simple model (1.1), there are more vacua with small cosmological constant if one starts with a larger magnitude of  $\lambda_{bare}$ , since the radius of the corresponding shell in flux space is larger.

In this note, I will point out a parallel argument for the distribution of values of scalar masses. After summarizing this in the next subsection, I will move on to discuss caveats to the simple picture (1.1) in §1.2. Finally in §2 I will explain a slightly more elaborate setup in which the mechanism for tuning scalar masses can be combined with the Higgs mechanism. Along the way I will describe the generalization of stringy tachyon couplings [5] to more general examples [6].

### 1.1. Scalar Masses

Now let us move from the cosmological constant to scalar masses. In [5], it was argued in the context of the AdS/CFT correspondence that RR fluxes couple to closed string

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<sup>1</sup> though in geometrical compactifications realizing the idea in [3,4] these coefficients are themselves nonlinear functions of the fluxes.

tachyons (specifically the twisted tachyon  $T$  of type II orbifolded by spacetime fermion number, otherwise known as “type 0”) via the interaction

$$\mathcal{L}_{TF} \sim -|T|^2 \left( \sum_I |F_I|^2 \right), \quad (1.2)$$

where  $I$  indexes different types of RR flux. This goes in the direction of stabilizing the tachyon (decreasing the magnitude of its negative mass squared in a nontrivial flux background). In §2 I will argue that a similar effect exists for more general twisted tachyons [6].<sup>2</sup>

The combination of the bare mass and flux contribution, applied to a compactification  $X$ , yields a mass formula of the form

$$m_T^2 = -\eta m_s^2 + \sum_I \tilde{c}_I |Q_I|^2 + \dots \equiv m_{bare}^2 + \sum_I \tilde{c}_I |Q_I|^2 \quad (1.3)$$

where the  $Q_I$  represent the independent flux quantum numbers (indexed by  $I$ ) of the RR fields coupling to the tachyon as in (1.2), and as in (1.1) the moduli-dependent coefficients  $\tilde{c}_I$  arise from the dimensional reduction on  $X$ . Here  $m_s$  is the string mass scale, and the term  $-\eta m_s^2$  represents the tree level tachyon mass squared (which in perturbative string theory ranges down to  $\eta$  of order one); the  $\dots$  represent other contributions to the tachyon mass which may arise depending on the background considered. We have included these other contributions with the  $-\eta m_s^2$  into a “bare” mass term; we will be interested in regions in the landscape in which  $m_{bare}^2 < 0$  so that the positive flux contribution can tune the total  $m_T^2$  close to zero.

This formula (1.3) is of the same form as the formula (1.1) for the structure of the cosmological term [1]. In particular, if one wishes to tune away the scalar mass to a small value, it is advantageous to start with a scalar such as  $T$  with a large bare negative mass squared. That is, as in the original Bousso-Polchinski argument (1.1), there are more choices of flux which tune away a larger bare negative value than ones which tune away a smaller bare negative value.

Tuning a scalar mass to a small value can also have the effect of tuning its VEV to a small value. For example, consider the cases with a stabilizing positive quartic potential

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<sup>2</sup> In the AdS/CFT context, one can analyze the vacuum structure at weak coupling to see if this effect actually stabilizes  $T$  in the limit of small radius and large flux density. One finds [7] that the instability has become milder at weak field theory coupling (though a corresponding instability evident in the one loop Coleman Weinberg potential remains).

term  $g_4|T|^4$  for  $T$ . Independently tuning the quadratic term (1.3) to a small negative value  $|m_T^2| \ll m_s^2 g_4$  produces a VEV of order

$$\langle T \rangle \sim |m_T|/\sqrt{g_4} \ll m_s \quad (1.4)$$

It would be interesting to apply this observation to scalar fields arising in models of the real world. One application might be to cosmological models requiring light scalar fields. The other obvious target for this mechanism is the Higgs particle of the standard model. In order to apply this observation to the Higgs, one requires a more realistic scalar field than the type 0 tachyon. Therefore in §2, I will turn to a slightly more elaborate setting in which this mechanism applies to gauge-charged scalar fields.

### *1.2. Further developments and important caveats*

Before proceeding to that, let me briefly review some further aspects of this line of development, including important qualifications applying to the simple models (1.1)(1.3) considered above.

The Bousso-Polchinski mechanism, while a compelling proposal especially given the presence of the basic ingredients (flux and compactification geometry) in string theory, suffered at the time from the lack of examples exhibiting sufficient compensating forces to plausibly fix all the moduli. That problem has been addressed in various special limits of string theory, specifically supercritical type II limits [8] and Calabi-Yau flux compactifications of the critical type IIB limit [3,4] (see also [9] for important work toward fixing all the moduli in various limits of the theory).

These models realize the spirit of the Bousso Polchinski mechanism, albeit with some important differences from the simple model (1.1). In the supercritical examples of [8], an explicit asymmetric orientifold fixes at tree level all of the runaway moduli except the dilaton. The dilaton is fixed by the contributions to the dilaton potential from supercriticality [10], orientifolds, and RR fluxes. This potential in the specific asymmetric orientifold models of [8] tunes the cosmological constant to small values by using a large  $D$  expansion rather than directly the Bousso Polchinski flux lattice, since in these particular models the flux lattice is too regular to provide by itself a finely spaced discretuum. In the examples of [4], the moduli-dependent coefficients  $c_i$  are themselves complicated functions of the fluxes since the fluxes help determine where the moduli get stabilized. As a result, the form (1.1) is at best schematic. Similarly, in (1.3) in general the coefficients  $\tilde{c}_I$  will be

nonlinearly related to the fluxes. However, as emphasized in the recent works [11], the issue is not so much the precise form of the positive terms in (1.1) (and now also (1.3)) but their independent distribution.

Another important caveat to simple arguments such as those reviewed above, as explained in [11], is that there may also be branches of solutions with vanishing coefficients of the classical contributions in (1.1)(1.3), such as cases in which the cosmological term and scalar masses are completely generated by non-perturbative effects. The question then becomes whether these cases are more numerous than those considered above. Similarly, in the case of the scalar masses in general we must consider off diagonal elements in the mass matrix, and their couplings to fluxes, which may yield a different accounting of preferred bare values.<sup>3</sup>

Finally, having stated these caveats, it is worth emphasizing that the Bousso-Polchinski argument naturally predicts a *nonzero* value for the cosmological constant among vacua without exact extended supersymmetry, since the discretuum produced by the quantized flux quantum numbers leads to the positive terms in (1.1) being discretely but not continuously tunable. As a result, the data [12] fits in well qualitatively with their mechanism in general and string theory in particular.

## 2. Applications

One application for the scalar mass analogue of Bousso-Polchinski is to models of inflation which require a light scalar field. Another is to the Higgs particle in the Standard Model. For both these applications, particularly the latter, it is of interest to generalize the mechanism in §1.2 beyond the Type 0 setting in which the couplings (1.2) were motivated [5].

### 2.1. More general tachyons and their $RR$ couplings

A set of generalizations of the Type 0 tachyon is twisted tachyons in  $Z_k$  orbifolds of flat space [6]. These are closely related to the type 0 tachyon as can be seen from the large  $k$  limit [6,13] and the “twisted circle” examples where the  $Z_k$  rotation is accompanied by a shift on another dimension [14]. They also arise in nonsupersymmetric orbifolds of  $AdS/CFT$  [15,7], where as in [5,7] one finds that the small radius large flux regime has

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<sup>3</sup> I thank S. Dimopoulos for discussions on this issue.

a milder instability than that arising at large radius. As a result, we expect that these localized tachyons also have couplings to RR fluxes of the form (1.2). This could be checked via an analogous worldsheet computation to (1.2)[5], with the same ambiguities regarding the off shell nature of the zero momentum tachyon; these couplings are not prohibited by any selection rules from the worldsheet symmetries. It could in principle be checked off shell in string field theory, along the lines of [16], though the superstring case is still somewhat out of reach.

In fact there is a simpler argument that the coupling (1.2) persists to the case of tachyons localized by geometrical orbifolds, at least in a regime accessible to a general relativistic analysis (which occurs far enough along the flow in [6] and can be enforced also by resolving the tip of the cone independently with an extra shift [14]). This follows from considering the basic form of the spacetime deformation induced by the localized tachyon condensation [6], for which we will then consider the effects of adding RR flux. The tachyon condensation produces a smoothing of the tip of the cone inside a shell of dilaton/graviton (outside of which is the remaining part of the original solution); see figure 1 of [6]. This deformation has the property that the region inside the shell is of smaller volume after the tachyon condensation than the corresponding region before the condensation (see the discussion below equation (4.9) in [6]).

In order to understand the analogue of the coupling (1.2) in the localized case, we then need to determine the effect of bulk RR flux on this expanding shell configuration. For realistic application, this tachyonic cone must be embedded in a compact manifold, and the RR fluxes integrated over cycles in this compactification manifold are quantized. Since the integral of a bulk RR flux over the cycles it threads is quantized, the integral of this flux over the region of space inside the shell is constant (since the region outside the shell is identical to the original orbifold background). Since, as we just reviewed, the tachyon condensation lowers the volume of space inside the shell, the energy density contained in the flux increases upon tachyon condensation. Thus the flux pushes the tachyon back toward  $T = 0$ . Because we are considering bulk RR fluxes, which come from the untwisted sector of the orbifold, the quantum symmetry of the orbifold prohibits any linear coupling of  $T$  to these fluxes, but permits a coupling of the form (1.2).

I should emphasize that this coupling of the form (1.2) in the case of the localized tachyons will only occur for those fluxes which have support where the tachyon is localized. This reduces the number of independent fluxes contributing to (1.2) relative to the Type 0 case, but can still permit large numbers of such fluxes.

## 2.2. Bousso-Polchinski for charged scalars

In order to apply this tuning to the Higgs mechanism, we require a scalar field charged under a non-abelian gauge group. This is not the case even for the more general tachyons [6] near the orbifold point, but these tachyons can implement this mechanism indirectly through their couplings to charged scalar fields on D-branes.

In the presence of D-branes at the (blown up) orbifold, the tachyon VEVs  $T$  determine the VEVs of D-brane scalar fields  $\phi$ . That is, at leading order in  $T$  (where  $T = 0$  is the orbifold limit) they contribute couplings of the form  $-|T|\phi^2$  for D-brane probes [6]. Hence a small value for the VEV of  $T$  leads to a small scale of symmetry breaking for the D-brane Higgs fields  $\phi$ .

For example, D-brane probes of  $\mathbf{C}/\mathbf{Z}_k$  leads to a quiver theory [17] including a gauge group  $U(N)^k$  with bifundamental scalar matter  $Z_{j,j+1}$  with  $j = 1, \dots, k$  indexing the factors in the gauge group. Including the leading effects of the twisted tachyons which in a convenient basis we will label  $T_j$  (satisfying  $\sum_j T_j = 0$ ), one obtains the following potential for the charged scalars  $Z$  in the presence of a fixed background  $T_j$  [6]

$$V(Z) = \sum_j (|Z_{j,j+1}|^2 - |Z_{j-1,j}|^2 + T_j)^2 + \mathcal{O}(T^2) \quad (2.1)$$

This leads generically to a Higgs mechanism breaking  $U(N)^k$  to  $U(N)$  (with special subspaces leading to  $U(N)^{k-2}$ ,  $U(N)^{k-4}$ , etc.). The value of the tachyon VEV determined by (1.3)(1.4) thus also governs the VEVs of the brane Higgs fields  $Z$ .

Thus, in simple situations with stringy tachyons, the tachyon VEV controls the VEVs of D-brane Higgs fields. Since the number of flux choices leading to a small VEV (1.4) for  $T$  is greater for greater bare negative tachyon mass squared, this generalization of the Bousso Polchinski mechanism to Higgs scalar masses favors string backgrounds with a highly tachyonic bare mass term.

## 3. Discussion

It is interesting that the mechanism explored in this paper in itself favors backgrounds which include those closely related to tachyonic non-supersymmetric perturbative string theories – starting points which are ordinarily discarded in considering application to the real world. Of course in order to reach definite conclusions about statistically favored vacua

of string theory, one needs much more comprehensive information about the full space of vacua, as well as ultimately dynamical information affecting vacuum selection.

There have been several very interesting forays recently into the statistics of supersymmetry breaking in various corners of M theory [11], related to studies of distributions of IIB flux vacua [18] and to new phenomenological ideas involving high scale supersymmetry breaking [19]. The studies [11] aim in part at determining whether string theory in itself (without phenomenological input) can be seen to predict in a statistical sense either low or high scale supersymmetry breaking. In other words, they aim to determine whether more tuning of the UV string theory is required to enforce low energy supersymmetry (having tuned the cosmological constant) than is required to tune the cosmological constant without low energy supersymmetry.

In this regard, it is important to note that there are many starting points with string scale or higher supersymmetry breaking, which were neglected in most of [11], whose contribution to the statistics must be included before a reliable accounting of the generic scale of supersymmetry breaking can be made. For example, the addition of spacefilling branes and antibranes, and non-critical dimensionality (dimension  $D \neq 10$  for the superstring), arise as consistent possibilities in perturbative string theory, with some such examples related by tachyon condensation to the better studied critical string backgrounds with low energy supersymmetry [10]. These a priori more generic starting points admit sufficient forces to fix their moduli and avoid tachyons in perturbation theory at weak bare string coupling [8].<sup>4</sup> These intrinsically string-scale SUSY breaking directions in discrete parameter space threaten to dwarf the statistical contribution of the special case of critical low energy supersymmetric models; for example the number of flux choices grows exponentially with dimensionality  $D$ , as does the number of independent topological quantum numbers of compactification geometries.<sup>5</sup>

One argument often made for supersymmetry is its enhancement of theoretical control; indeed it has led to tour-de-force calculations of the two derivative effective action of some interesting theories. On the other hand, in the context of the problem of fixing the moduli

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<sup>4</sup> They may generically be subject to non-perturbative decays, as in [4], or in some special cases [20], but this would still allow them to persist for timescales exponentially long as a function of appropriate coupling constants.

<sup>5</sup> More simply: the statistics of string theory papers may be quite different from the statistics of string theory vacua. For example, ten and two dimensional target spaces have been by far the most popular, but the other dimensions are available.



it can happen that non-renormalization theorems simply postpone needed moduli-fixing contributions to highly subleading (e.g. non-perturbative) orders in an expansion about weak coupling. In fixing moduli, ultimately one balances different orders in this expansion off each other, and in some ways it is simpler to implement this in situations where the needed forces arise at early orders in perturbation theory (as in [8]). As in the analogous perturbative analysis of gauge theories [21], a perturbative balancing of contributions at different orders can produce well controlled non-supersymmetric backgrounds via the introduction of large discrete quantum numbers. In any case, theoretical control does not in itself provide an argument for an observational prediction.

Despite these difficulties, I find the statistical program [1,18,11] for seeking generic properties of string vacua extremely interesting, particularly in its prospects for refining our notions of naturalness. The mechanism discussed in this note provides a new twist on these issues in the regions of the landscape to which it applies.

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