



# The string landscape, black holes and gravity as the weakest force

## Citation

Arkani-Hamed, Nima, Luboš Motl, Alberto Nicolis, and Cumrun Vafa. 2007. "The String Landscape, Black Holes and Gravity as the Weakest Force." *Journal of High Energy Physics* 2007 (6): 060–060. <https://doi.org/10.1088/1126-6708/2007/06/060>.

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# The String Landscape, Black Holes and Gravity as the Weakest Force

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## Abstract

We conjecture a general upper bound on the strength of gravity relative to gauge forces in quantum gravity. This implies, in particular, that in a four-dimensional theory with gravity and a  $U(1)$  gauge field with gauge coupling  $g$ , there is a new ultraviolet scale  $\Lambda = gM_{\text{Pl}}$ , invisible to the low-energy effective field theorist, which sets a cutoff on the validity of the effective theory. Moreover, there is some light charged particle with mass smaller than or equal to  $\Lambda$ . The bound is motivated by arguments involving holography and absence of remnants, the (in) stability of black holes as well as the non-existence of global symmetries in string theory. A sharp form of the conjecture is that there are always light “elementary” electric and magnetic objects with a mass/charge ratio smaller than the corresponding ratio for macroscopic extremal black holes, allowing extremal black holes to decay. This conjecture is supported by a number of non-trivial examples in string theory. It implies the necessary presence of new physics beneath the Planck scale, not far from the GUT scale, and explains why some apparently natural models of inflation resist an embedding in string theory.

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# 1 Introduction

By now it is clear that consistent theories of quantum gravity can be constructed in the context of string theory. This can also be done in diverse dimensions by considering suitable compactifications. This diversity, impressive as it may be for a consistent theory to possess, poses a dilemma: The theory appears to be more permissive than desired! However it was recently suggested [1] that the landscape of consistent theories of gravity one obtains in string theory is by far smaller than would have been anticipated by considerations of semiclassical consistency of the theory. The space of consistent low-energy effective theories which cannot be completed to a full theory was dubbed the ‘swampland’. Certain criteria were studied in [1] to distinguish the string landscape from the swampland. For example, one such criterion was the finiteness of the number of massless fields (see also [2] for a discussion of this point).

In this paper we propose a new criterion which distinguishes parts of the swampland from the string landscape. This involves the simple observation, clearly true in our own world, that “gravity is the weakest force”. We promote this to a principle and in fact find that surprisingly it is demanded by all consistent string theory compactifications! Roughly speaking this is the statement that there exist two “elementary” charged objects for which the repulsive gauge force exceeds the attractive force from gravity. More precisely, the conjecture we make is that this is true for a stable charged particle which minimizes the ratio  $|M/Q|$ . In other words one of our main conjectures in this paper is that (in suitable units) the minimum of  $|M/Q|$  is less than 1.

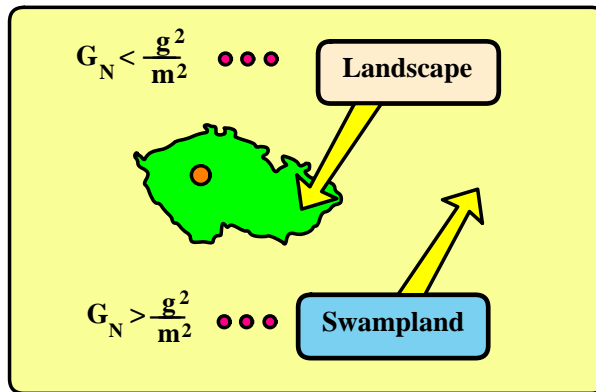
We motivate this conjecture from various viewpoints. In particular we show how this conjecture follows from the assumption of finiteness of the number of stable particles which are not protected by a symmetry principle. This finiteness criterion nicely extends some of the finiteness criteria discussed in [1, 2] in the context of bounds on the number of massless particles in a gravitational theory. We show why if our conjecture were not true, there would be an infinite tower of stable charged particles not protected by any symmetry principle. The conjecture also ties in nicely with the absence of global symmetries in a consistent theory of gravity.

Our conjecture, if true, has a number of consequences: For extremal (not necessarily supersymmetric) black holes the bound is  $M/|Q| = 1$ . This suggests that there are corrections to this formula for smaller charges, which makes it (in the generic case) an inequality  $M/|Q| < 1$ . Another aspect of our conjecture is that it naturally suggests that the  $U(1)$  effective gauge theory breaks down at a scale  $\Lambda$  well below the Planck scale  $\Lambda \sim gM_{\text{Pl}}$  (more precisely  $\Lambda \sim \sqrt{\alpha/G_N}$ ), where  $g$  is the  $U(1)$  gauge coupling constant. These restrictions

of low cutoff scales and forced presence of light charged particles are very surprising to the effective field theorist, who would not suspect the existence of the new UV scale  $\Lambda$ . As long as the Landau pole of the  $U(1)$  is above the Planck scale, the low-energy theorist would think that the cutoff of the effective theory should be near  $M_{\text{Pl}}$ , and if anything, smaller  $g$  seems to imply that the theory is getting even more weakly coupled.

Our conjecture, if true, has a number of consequences. If  $g$  is chosen to be one of the Standard Model gauge couplings near the unification scale, the scale  $\Lambda$  is necessarily beneath the Planck scale, close to the familiar heterotic string scale  $\sim 10^{17}$  GeV. Furthermore, the observation of any tiny gauge coupling, for instance in sub-millimeter tests of gravity, would necessitate a low-scale  $\Lambda$  far beneath the Planck/GUT scales. Finally our conjecture also explains why certain classes of apparently natural effective theories for inflation, involving periodic scalars or axions with parametrically large, super-Planckian decay constants, have resisted an embedding in string theory. [3].

This restriction has a number of phenomenological consequences. It implies that, extrapolating the Standard Model to high energies, there *must* be new scale  $\Lambda$  beneath the Planck scale with  $\Lambda \sim \sqrt{\alpha_{\text{GUT}}/G_{\text{N}}} \sim 10^{17}$  GeV, and that any experimental observation of extremely weak gauge coupling must be accompanied by new ultraviolet physics at scales far beneath the Planck/GUT scales.



**Figure 1.** Consistent theories of quantum gravity (the landscape) represent a small portion of the effective field theories (the swampland) where additional conditions such as the weakness of gravity are satisfied.

The organization of this paper is as follows: In section 2 we present some loose conjectures which motivate our more precise conjectures discussed in section 3. In section 4 we present evidence for our conjecture drawn from string theory. We conclude with a discussion of

certain additional points in section 5.

## 2 Loose conjectures

### 2.1 Weak electric and magnetic gauge couplings

Consider a 4-dimensional theory with gravity and a  $U(1)$  gauge field with gauge coupling  $g$ . Naively, the effective theory breaks down near the scale  $M_{\text{Pl}}$  where gravity becomes strongly coupled, and certainly nothing seems to prevent taking  $g$  as small as we wish; if anything the effective theory seems to become even more weakly coupled and consistent.

Nonetheless, we claim that for small  $g$  there is a hidden new ultraviolet scale far beneath the Planck scale. For instance, we claim that there must be a light charged particle with a small mass

$$m_{\text{el}} \lesssim g_{\text{el}} M_{\text{Pl}} . \quad (1)$$

This statement should also hold for magnetic monopoles,

$$m_{\text{mag}} \lesssim g_{\text{mag}} M_{\text{Pl}} \sim \frac{1}{g_{\text{el}}} M_{\text{Pl}} . \quad (2)$$

Note that the monopole masses are a probe of the ultraviolet cutoff of a  $U(1)$  gauge theory. The monopole has a mass at least of order the energy stored in the magnetic field it generates; this is linearly divergent, and if the theory has cutoff  $\Lambda$ , this is of order

$$m_{\text{mag}} \sim \frac{\Lambda}{g_{\text{el}}} \quad (3)$$

which is indeed parametrically correct in all familiar examples—for the  $U(1)$  arising in the Higgsing of an  $SU(2)$ , this is the correct expression for the monopole mass with  $\Lambda \rightarrow m_W$ , while on a lattice with spacing  $a$ , this is the monopole mass with  $\Lambda \rightarrow 1/a$ . Therefore, the above constraint on monopole masses tells us that for small  $g$ , the effective theory must break down at a prematurely low scale

$$\Lambda \lesssim g M_{\text{Pl}} . \quad (4)$$

This conjecture can be rephrased as the plausible statement that “gravity is the weakest force”; if charged particles have  $m > g M_{\text{Pl}}$ , then the gauge repulsive force between them is overwhelmed by the gravitational attraction—while if there are states with  $m \lesssim g M_{\text{Pl}}$ , then gravity is subdominant. Of course, in highly supersymmetric systems, all velocity-independent forces vanish in any case, but as we will see, our constraints still apply.

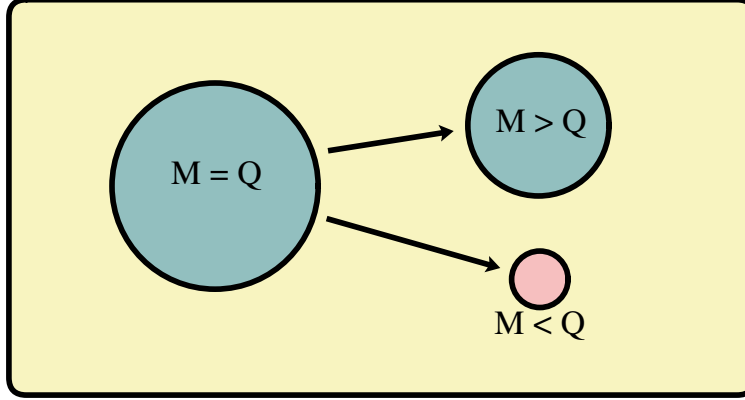
Finally, we are familiar with restrictions on the relative magnitudes of masses and charges from BPS bounds, but these have the opposite sign, telling us that  $M \geq Q$  in some units. Our bound might appear to be an “*anti*-BPS” bound. This is not quite accurate. First of all the BPS case is a limiting case of our bound, saturating it. Secondly we are not conjecturing that for a given charge sector all the masses are less than the charge, but that there exist *some* such state.

We have phrased our constraint as one on the strength of  $U(1)$  gauge couplings. This is a running coupling so we should ask about the scale at which it should be evaluated. For the statement that there exists a light charged particle with mass  $m < gM_{Pl}$ , it is natural to use the asymptotic value of  $g$ , the running coupling evaluated at the mass of the lightest charged particle. But for the statement about the existence of a cutoff  $\Lambda < gM_{Pl}$ , it is clearly most natural to consider the running coupling near the scale  $\Lambda$ . Clearly our bound will also apply to non-Abelian gauge theories, that can be Higgsed to  $U(1)$ ’s. In this case, the mass of the  $W$ ’s is  $m_W \sim g_{el}v$  where  $v$  is an appropriate vev, and these particles will satisfy our bound as long as the vevs don’t exceed the Planck scale, a statement not independent of conjecture about the finiteness of the volume of moduli spaces in [1].

## 2.2 Black holes and global symmetries

Why should such a conjecture be true? One motivation has to do with the well-known argument against the existence of global symmetries in quantum gravity. Gauge symmetries are of course legal, but as we take the limit  $g \rightarrow 0$ , the symmetry becomes physically indistinguishable from a global symmetry. Something should stop this from happening, and our conjecture provides an answer. As the gauge coupling goes to zero  $g \rightarrow 0$ , the cutoff on the effective theory  $\Lambda \rightarrow 0$  as well, so that the limit  $g \rightarrow 0$  can not be taken smoothly.

To see this more concretely, imagine that we have  $g \sim 10^{-100}$ , and consider a black hole with mass  $\sim 10M_{Pl}$ . Since the gauge coupling is so tiny, the black hole can have any charge between 0 and  $\sim 10^{100}$ , and still be consistent with the bound for having a black hole solution ( $M \geq QM_{Pl}$ ). But if there are no very light charged particles, none of this charge can be radiated away as the black hole Hawking evaporates down to the Planck scale. But then we will have a Planckian black hole labeled by a charge anywhere from 0 to  $10^{100}$ . This leads to  $10^{100}$  Planck scale remnants suffering from the same problems that lead us to conclude quantum gravity shouldn’t have global symmetries (see e.g. [4]).



**Figure 2.** An extremal black hole can decay only if there exist particles whose charge exceeds their mass.

The difficulties involving remnants are avoided if macroscopic black holes can evaporate all their charge away, and so these states would not be stable. Since extremal black holes have  $M = Q M_{\text{Pl}}$ , in order for them to be able to decay into elementary particles, these particles should have  $m < q M_{\text{Pl}}$ . Our conjecture also naturally follows from Gell-Mann’s totalitarian principle (“everything that is not forbidden is compulsory”) because there should not exist a large number of exactly stable objects (extremal black holes) whose stability is not protected by any symmetries.

Another heuristic argument leading to same limit on  $\Lambda$  is the following. Consider the minimally charged monopole solution in the theory. With a cutoff  $\Lambda$ , its mass is of order  $M_{\text{mon}} \sim \Lambda/g^2$  and its size is of order  $R_{\text{mon}} \sim 1/\Lambda$ . It would be surprising for the *minimally* charged monopole to already be a black hole because the values of all charges carried by a black hole should be macroscopic (and effectively continuous); after all, a black hole is a classical concept. Demanding that this monopole is not black yields

$$\frac{M_{\text{mon}}}{M_{\text{Pl}}^2 R_{\text{mon}}} \lesssim 1 \quad \Rightarrow \quad \Lambda \lesssim g M_{\text{Pl}} \quad (5)$$

### 2.3 Simple parametric checks

It is easy to check the conjecture in a few familiar examples. For  $U(1)$ ’s coming from closed heterotic strings compactified to four dimensions, for instance, we have

$$g M_{\text{Pl}} \sim M_s, \quad (6)$$



and there are indeed light charged particles beneath the string scale  $M_s$ , which also sets the cutoff of the effective theory. In Kaluza-Klein theory, we have

$$gM_{\text{Pl}} \sim \frac{1}{R}; \quad (7)$$

again there are charged particles at this scale and it also acts as the cutoff of the low-energy 4D effective theory. By  $T$ -duality, it is easy to see that the same thing will be true for winding number gauge symmetries, too.

Next let's consider  $U(1)$ 's in type I string theories compactified to 4D. Here

$$gM_{\text{Pl}} \sim \frac{M_s}{\sqrt{g_s}} \quad (8)$$

and indeed at weak string coupling this is even larger than the string scale. At large coupling, we revert to a heterotic dual. Similarly, we can consider a  $U(1)$  living on a stack of  $Dp$ -branes. Then

$$gM_{\text{Pl}} \sim \frac{M_s}{\sqrt{g_s(RM_s)^{(9-p)}}} \quad (9)$$

where  $R$  is a typical radius of the space transverse to the brane. Again, here we could try and violate the bound for  $(RM_s) \ll 1$  but then we revert to a  $T$ -dual description.

If we move a single D-brane far away from others, the objects charged under its  $U(1)$  become long strings and can be made heavier and heavier. For a non-compact space, it is clear that they can be made arbitrarily heavy—but to have 4D gravity, the space must be compact and this limits the mass of these charged particles. For a brane of codimension bigger than two in the large dimensions, which we'll take to have comparable sizes  $R$ , the back-reaction of the brane on the geometry is small and we can limit  $m_W \lesssim M_s R$ . Then, for instance for D3-branes moving in a spacetime with  $n$  large dimensions we have

$$\frac{m_W}{gM_{\text{Pl}}} \sim \sqrt{\frac{g_s}{(M_s R)^{(n-2)}}} \quad (10)$$

so again for  $n > 2$ ,  $g_s < 1$ , and  $R > l_s$ , the ratio is smaller than one. The magnetic monopole is a long D-string attached to the brane, so the analogous ratio depends on  $1/g$  instead and

$$\frac{m_{\text{mon}}}{(1/g)M_{\text{Pl}}} \sim \sqrt{\frac{1}{(M_s R)^{(n-2)}}} \quad (11)$$

is again satisfied.

Things are parametrically marginal for  $n = 2$ , but look like they might be violated for  $n = 1$ . As an example, suppose we have an  $AdS_5 \times X$  space, with  $L_{AdS}$  parametrically

larger than the string scale. By the introduction of a “Planck brane” we have a warped compactification down to 4D. Suppose also that we can have some D3-branes in the space with a  $U(1)$  gauge field living on them. To an effective field theorist, there is no obstacle to imagining that the internal space  $X$  is small, of order the string scale. But this is in conflict with our conjecture. The  $W$ ’s charged under the  $U(1)$  correspond to a long string ending on the D3-brane and hence have a mass scaling as  $L_{AdS}M_s^2$ , the  $U(1)$  gauge coupling is  $g_{YM}^2 = g_s$  while  $M_{\text{Pl}}^2 = M_s^8 V_X L_{AdS} / g_s^2$ . The requirement that

$$m_W^2 \lesssim g_{YM}^2 M_{\text{Pl}}^2 \tag{12}$$

then implies that

$$(V_X M_s^5) \gtrsim g_s (L_{AdS} M_s). \tag{13}$$

In other words, the volume of the internal 5D space *must* be as large as  $g_s L_{AdS}$  in string units. This is certainly true for the familiar  $AdS_5 \times S^5$  compactifications, where the volume of the  $S^5$  scales like  $L_{AdS}^5$ . One might try to orbifold the internal  $S^5$  to smaller volumes, but there are only 3  $U(1)$ ’s one might orbifold by, and since the size of the “slices” in the  $S^5$  can’t get much smaller than  $l_s$ , this still assures that the volume of  $X$  is larger than  $L_{AdS}^2$  in Planck units and safely satisfies our bound. We aren’t aware of any examples where  $X$  can be kept at the string scale with parametrically large  $L_{AdS}$ .

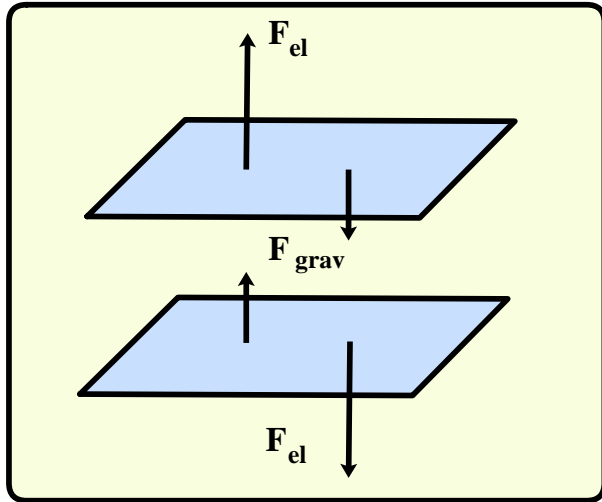
Finally, let us consider another possible way in which our bound might be parametrically violated. Consider a theory with some cutoff  $\Lambda$  but a large  $N$  number of  $U(1)$ ’s, perhaps associated with wrapped brane charges along a large number of cycles in some compactification. If we could somehow Higgs these  $U(1)$ ’s down to the diagonal subgroup, the low-energy coupling  $g_{diag}$  would be suppressed by  $\sim \frac{1}{\sqrt{N}}$ , and we can make the coupling very weak. However, at large  $N$ , there is also a “species problem”— $N$  can’t be made parametrically large without making gravity weak [4]. For instance, naively the cycles would have to each occupy a volume of order the string scale, so that the internal volume and hence  $M_{\text{Pl}}^2$  also grows as  $N$ , and hence  $N$  cancels out of the combination  $g_{diag} M_{\text{Pl}}$ . Similar issues were discussed in the “ $N$ -flation” proposal of [5].

## 2.4 Generalizations

It is natural to generalize our loose conjecture: consider a  $p$ -form Abelian gauge field in any number of dimensions  $D$ ; then there are electrically and magnetically charged  $p - 1$  and  $D - p - 1$  dimensional objects with tensions

$$T_{\text{el}} \lesssim \left(\frac{g^2}{G_N}\right)^{1/2}, \quad T_{\text{mag}} \lesssim \left(\frac{1}{g^2 G_N}\right)^{1/2}, \tag{14}$$

where the coupling  $g$  (the charge density) has a dimension of  $mass^{p+1-D/2}$ . We have discussed the case  $D = 4$ ,  $p = 1$  above. Whenever the objects carry central charges only, the inequalities above are saturated and coincide with the BPS bounds. However, the inequality becomes strict for other types of charges. This generalization can be used to rule out effective field theories that have been constructed in the literature for a variety of purposes.



**Figure 3.** For any kind of an electric field, there should exist “self-non-attractive” objects for whom the electric repulsive force exceeds the strength of gravity.

For instance, in [6], it was argued that a natural candidate for an inflaton could be found in 5D gauge theories compactified on a circle. Consider a  $U(1)$  gauge theory with gauge coupling  $g_5$  and Planck scale  $M_5^3$ . Compactifying on a circle, we have 4D gravity as well as the 4D periodic scalar  $\theta = A_5 R$  associated with the Wilson line around the circle. The effective action for  $\theta$  and gravity is simply

$$\int d^4x \sqrt{-g} \left[ F^2 (\partial\theta)^2 + M_{Pl}^2 R \right], \quad (15)$$

where

$$F^2 = \frac{1}{g_5^2 R}, \quad M_{Pl}^2 = M_5^3 R. \quad (16)$$

At one-loop level, light charged scalars can generate a potential  $V(\theta) \sim (1/R^4) \cos(\theta)$ . It is easy to see that  $\theta$  can have slow-roll inflation as long as  $F^2 \gg M_{Pl}^2$ , or

$$g_5^2 M_5^3 R^2 \ll 1. \quad (17)$$

This is perfectly consistent in effective field theory, but it runs afoul of our general constraint for  $D = 5$ ,  $p = 1$ . Indeed, straightforward attempts to embed this model into string theory fail. For instance, for the  $U(1)$  coming from closed strings,  $g_5^2 M_5^3 = M_s^2$  and so we must have  $R \ll M_s^{-1}$ , where this description breaks down and the  $T$ -dual is appropriate.

Indeed, Banks, Dine, Fox, and Gorbatov [3] subsequently studied the more general question of whether in compactifications to 4D it is possible to get periodic scalars (“axions”) with decay constants  $F$  parametrically larger than the Planck scale. In all the examples they studied, they found that either  $F$  can’t be made larger than  $M_{\text{Pl}}$ , or that if it could, there was also an instanton of anomalously small action  $S_{\text{inst}} \sim M_{\text{Pl}}/F$ , so the instanton generated unsuppressed potential generated terms up to  $\cos(N\theta)$  with  $N \sim M_{\text{Pl}}/F$ , ruining the parametric flatness of the potential. This observation is subsumed in our generalized conjecture; for  $D = 4$ ,  $p = 0$ , the 0-form is an axion; the “tension” of the object charged under it is simply the action of an instanton coupling to the axion, while the axion gauge coupling is  $g \sim 1/F$  where  $F$  is the axion decay constant so our constraint gives precisely

$$S_{\text{inst}} \lesssim \frac{M_{\text{Pl}}}{F}. \quad (18)$$

### 3 Sharpening the claim

Working in  $M_{\text{Pl}} = 1$  units, we are making a conjecture about mass/charge ratios

$$(M/Q) \lesssim 1 \quad (19)$$

To find a sharper conjecture, we have to decide (a) what states should satisfy this bound and (b) what to mean by “1”. For the last point, it is natural to take the  $(M/Q)$  ratio that is equal to one for large extremal black holes. As for the states to consider, there are three natural possibilities:

- (I)  $\left(\frac{M_{q_{\min}}}{q_{\min}}\right) \leq 1$ , for the state of minimal charge;
- (II)  $\left(\frac{M_{\min}}{q_{M_{\min}}}\right) \leq 1$ , for the lightest charged particle;
- (III)  $\left(\frac{M}{q}\right)_{\min} \leq 1$ , for the state with smallest mass/charge ratio.

Of course, for these statements to have a sharp meaning, the state must be exactly stable for  $M$  to be meaningful. The particle of smallest charge is *not* guaranteed to be stable—for

instance, a heavy charged particle of charge  $+1$  can decay into two lighter charged particles of charge  $-2, +3$ . If the particles with charges  $-2$  and  $+3$  are light, they will form a Kepler/Coulomb bound state of charge  $+1$ . This state will be stable but its  $M/Q$  ratio may be larger than for the states with charges  $-2$  and  $+3$ . In particular, it may be larger than one.

Furthermore, there are easy counterexamples to the conjecture (I) in string theory, even when the minimally charged particles *are* exactly stable. For instance, in the weakly coupled  $SO(32)$  heterotic string, the spinor of  $SO(32)$  is exactly stable and has minimal half-integral charges under the  $U(1)$ 's inside the  $SO(32)$ , but is heavy and can violate our bounds. A generalization of these are the half-integrally charged winding strings considered by Wen and Witten [7], that have fractionally charges but are also heavy. So (I) can't be right.

Of course the *lightest* charged particle is exactly stable, as is the particle with smallest  $(M/Q)$  (as follows trivially from the triangle inequality), so both (II) and (III) are well-defined conjectures. Obviously (II) is the stronger of the two (and it clearly implies (III)) and *forces* the effective theory to contain a light charged particle. Conjecture (III) can in principle be satisfied by a heavy state with large  $Q$  which would reduce the impact of the inequality on physics at low energies. Even though we have no counterexamples for conjecture (II) most of our evidence only supports the weaker conjecture (III).

When there are several  $U(1)$ 's, the generalization of the conjecture is clear. In every direction in charge space, including electric and magnetic charges, at large values of the charges, we have extremal black hole solutions. The conjectures (II) and (III) then imply the existence of light charged particles with  $(M/Q) < (M/Q)_{\text{extremal}}$  in certain directions of charge space. More precisely, there should always exist a set of directions in the charge space that form a basis of the full space where the inequality is satisfied.

It is interesting to see what the spectrum of a theory which violates our conjectures looks like. Suppose for simplicity that there is only one “elementary” particle with minimal charge 1, but with  $M > Q$ . Since the net force between two of these particles is attractive, there is a Kepler bound state of two of these particles, with charge 2, but with a mass smaller than  $2M$ , so that the mass/charge ratio decreases. We can continue to add further particles to make further bound states, with  $(M/Q)$  continually decreasing. This proceeds till the bound state eventually turns into an extremal black holes, and asymptotically, we reach  $(M/Q) = 1$ . It is easy to see that *all* of these particles are exactly stable: since  $(M/Q)$  is a decreasing function of  $Q$ , none of these states can decay into a collection of particles with smaller charges.

On the other hand, if there are any states with  $(M/Q) < 1$ , then the macroscopic black holes can always decay, and the number of exactly stable particles will be finite. Suppose that, among the states with  $(M/Q) < 1$ , the one with smallest charge has charge  $Q_{\min}$ . Then, by the same argument as above, we expect that the lightest particles with charges smaller than  $Q_{\min}$  are exactly stable.

So our mass/charge ratio conjecture (III) can be seen to follow from a very simple general conjecture valid for both charged and uncharged particles: The number of exactly stable particles in a theory of quantum gravity in asymptotically flat space is finite. Actually this statement is not quite correct. Clearly we can have an infinite number of exactly stable BPS states, and many of these are safely bound; consider for instance dyons of electric/magnetic charge  $(n, 1)$  for large  $n$ . However, the number of exactly stable (and safely bound) states *in any given direction in charge space* is finite.

Even for neutral particles, this implies that the number of massless degrees of freedom is finite, and such a restriction is indeed suggested by the species problem associated with the Bekenstein bound. If there is a principle dictating the number of exactly stable particles to be finite, it is reasonable to expect that in all the vacua in the landscape, the number of exactly stable states is typically of order a few. In this case, the minimal charge  $Q_{\min}$  for which  $(M/Q) < 1$  should not be too large, since as we saw above the number of exactly stable states grows with  $Q_{\min}$ . This then substantiates our loose conjectures  $m \lesssim gM_{\text{Pl}}$ .

## 4 Evidence for the conjecture

Our conjecture is now phrased sharply enough that we can look for non-trivial checks of it in known stringy backgrounds. Clearly in highly supersymmetric situations where  $U(1)$ 's are associated with central charges, there will be BPS states saturating our inequality. This will for instance be the case in theories with 32 supercharges. However, already with 16 supercharges non-trivial checks are possible, for instance in compactifications of the heterotic string on tori with generic Wilson lines, where most of the  $U(1)$ 's are not central charges.

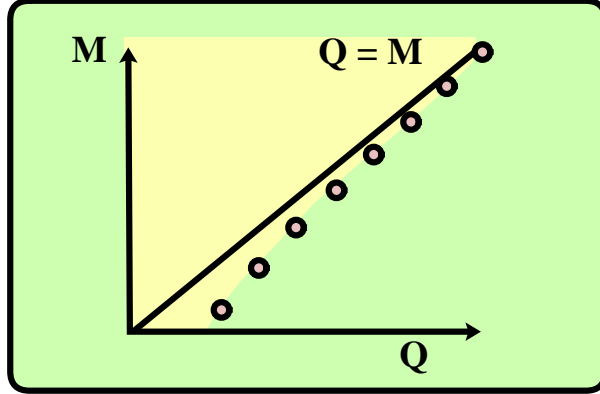
Consider for instance the  $SO(32)$  heterotic string compactified on  $T^6$ . At a generic point on moduli space, there is a  $U(1)^{28}$  gauge symmetry. We will check our conjecture for electric charges only; by  $S$ -duality, this check will carry over to magnetic charges as well. A general set of electric charges is a 28-dimensional vector

$$Q = \begin{pmatrix} Q_L \\ Q_R \end{pmatrix} \tag{20}$$

where  $Q_L$  is 22-dimensional vector and  $Q_R$  is 6-dimensional vector. The charges are quantized, lying on the 28-dimensional even self-dual lattice with

$$Q_L^2 - Q_R^2 \in 2\mathbb{Z} \quad (21)$$

Moving around in moduli space corresponds to making  $SO(22, 6)$  Lorentz transformations on the charges.



**Figure 4.** The charge  $M$  of the heterotic string states of charge  $Q$  approaches the  $M = Q$  line from below. The yellow area denotes the allowed region.

The extremal black hole solutions in this theory were constructed by Sen [8]. For  $Q_R^2 - Q_L^2 > 0$ , there are BPS black hole solutions with mass

$$M^2 = \frac{1}{2}Q_R^2 \quad (22)$$

where we work in units with  $M_{\text{Pl}} = 1$ . For  $Q_L^2 - Q_R^2 > 0$ , the black holes are *not* BPS; still, the extremal black holes have mass

$$M^2 = \frac{1}{2}Q_L^2. \quad (23)$$

We can compare this with the spectrum of perturbative heterotic string states, given by

$$M^2 = \frac{1}{2}Q_R^2 + N_R = \frac{1}{2}Q_L^2 + N_L - 1 \quad (24)$$

where  $N_{R,L}$  are the string oscillator contributions and where we chose units with  $\alpha' = 4$ . The  $-1$ , coming from the tachyon in the left-moving bosonic string, is crucial. Note that this spectrum nicely explains the BH spectrum of the theory, as the highly excited strings are progenitors of extremal black holes. Consider large  $Q_L, Q_R$ , with  $Q_R^2 > Q_L^2$ . Then,

the minimal  $M^2$  compatible with these charges will have  $N_R = 0$ ,  $N_L = \frac{1}{2}(Q_R^2 - Q_L^2) + 1$ , which are BPS, with  $M^2 = \frac{1}{2}Q_R^2$ . On the other hand, for  $Q_L^2 > Q_R^2$ , the minimal  $M^2$  is with  $N_L = 0$ , and  $N_R = \frac{1}{2}(Q_L^2 - Q_R^2) - 1$ . These are not BPS, but for large  $Q_L^2$ , they have  $M^2 = \frac{1}{2}Q_L^2$ .

But the string spectrum also guarantees that, as we go down to smaller charges along a basis of directions in charge space, we are guaranteed to find a state with a mass/charge ratio smaller than for extremal BH's. The inequality is saturated for the BPS states which have  $Q_R^2 > Q_L^2$ , but for  $Q_L^2 > Q_R^2$  the extremal black holes have  $M^2 = \frac{1}{2}Q_L^2$  while there is always a state with mass

$$M^2 = \frac{1}{2}Q_R^2 = \frac{1}{2}Q_L^2 - 1 \quad (25)$$

since there is a charge vector with  $Q_L^2 - Q_R^2 = 2$  on the charge lattice.

## 4.1 Gauge symmetries vs. global symmetries

It is possible to generalize this argument to *any* perturbative heterotic string compactification, including compactifications on  $K3$  and arbitrary Calabi-Yau threefolds, as a straightforward generalization of the familiar argument that all global symmetries in this theory are gauged. For any integral  $U(1)$  gauge symmetry coming from the left-movers of heterotic string, there is a worldsheet current  $J(z)$ . Because it is a  $(1, 0)$  primary field, one can construct the  $(1, 1)$  vertex operator  $J(z) \bar{\partial} X_\mu e^{ikX}(z, \bar{z})$  with  $k^2 = 0$  for a spacetime gauge field and prove that the corresponding symmetry is a gauge symmetry. (Analogously, a symmetry coming from the right-movers would be associated with a current  $\tilde{J}(\bar{z})$ , a  $(0, 1)$  primary field, and the vertex operator would be  $\tilde{J}(\bar{z}) \partial X_\mu e^{ikX}(z, \bar{z})$  with  $k^2 = 0$ .) We can take the worldsheet CFT to consist of the  $U(1)$  part together with the rest. We can bosonize the current with level  $k$  as

$$J(z) \sim \sqrt{k} \partial\phi(z). \quad (26)$$

Then, the operator

$$O(z) = :e^{i\sqrt{k}\phi(z)}: \quad (27)$$

is always in the CFT. There are a number of ways to see this. One way is to note that this operator has local OPE with all the operators in the theory. In fact using the integrality of the  $U(1)$  charge, the content of any operator will be of the form

$$V \sim : \exp(ip\phi/\sqrt{k}) : \cdot V'$$

where  $p$  is an integer-valued charge and where  $V'$  has no exponential parts in  $\phi$ . It is easy to see that  $O(z)$  will have local OPE with this operator. Therefore, the completeness of



the CFT spectrum (i.e. the statement that the operator content of the theory is maximal consistent with local OPE, as follows from modular invariance) forces us to have  $O(z)$  as an allowed operator in the theory.

Another way to understand the existence of the operator  $O(z)$  is to note that it corresponds to spectral flow by 1 unit in the  $U(1)$ . This simply corresponds to changing the boundary conditions on the circle by  $\exp(2\pi i \theta p)$  where  $p$  denotes the  $U(1)$  charge and  $\theta$  goes from 0 to 1.

We thus see that the state corresponding to  $O(z)$  exists in the spectrum of CFT. Since by assumption this is a left-mover state, this corresponds to  $N_L = 0$  and so  $M^2 = \frac{1}{2}Q_L^2 - 1$ , while asymptotically, the excited strings correspond to extremal black holes with  $M^2 = \frac{1}{2}Q_L^2$ , so our string state is indeed *sub*-extremal.

## 5 Possible relation to subluminal positivity constraints

It is natural to conjecture that since there must exist states for which  $(M/Q) < 1$  while the extremal black holes have  $(M/Q) = 1$ , the extremal limit for  $(M/Q)$  for black holes is approached from below, that is, that the leading corrections to the extremal black hole masses from higher-dimension operators should again *decrease* the mass. This implies some positivity constraint on some combination of higher-dimension operators.

It is interesting that similar positivity constraints have been discussed in [9], where it was found that certain higher-dimension operators must have positive coefficients in order to avoid the related diseases of superluminal signal propagation around configurations with a nonzero field strength and bad analytic properties of the  $S$ -matrix. For instance, consider the theory of a  $U(1)$  gauge field in four dimensions. The leading interactions are  $F^4$  terms, and the effective Lagrangian is of the form

$$- F_{\mu\nu}^2 + a(F^2)^2 + b(F\tilde{F})^2 + \dots \quad (28)$$

If the scale suppressing the dimension 8 operators is far beneath the Planck scale, we can ignore gravity, and the claim of [9] is that  $a, b$  must be positive to avoid superluminal propagation of signals around backgrounds with uniform electric or magnetic fields, and also to satisfy analyticity and dispersion relation for the photon-photon scattering amplitude.

Of course these higher dimension operators also change the mass/charge relation for extremal black holes. Indeed, there are many other operators which do this as well; at the leading order they include  $R^2$  and  $RF^2$  type terms as well. But we can imagine that the  $F^4$

terms dominate in the limit where the scale suppressing the  $F^4$  terms is far smaller than the Planck scale. Treating the  $a, b$  terms as perturbations, we can solve for the modified Black Hole background to first order in  $a, b$ , and find the new bound on  $M$  which has a horizon and no naked singularity. To first order in  $a, b$ , we find that for a black hole with electric and magnetic charges  $(Q_e, Q_m)$  and working with  $M_{\text{Pl}} = 1$

$$M_{\text{extr}}^2 = (Q_e^2 + Q_m^2) - \frac{2a(Q_e^2 - Q_m^2)^2}{5(Q_e^2 + Q_m^2)^2} - \frac{32b}{5} \frac{Q_e^2 Q_m^2}{(Q_e^2 + Q_m^2)^2} \quad (29)$$

So, for purely electric or magnetic black holes, we have

$$M_{\text{extr}}^2 = Q^2 - \frac{2a}{5} \quad (30)$$

which indeed decreases for the “right” sign of  $a > 0$ . The same statement holds for the dyonic black holes as long as  $b > 0$  which is also the “right” sign. The result (29) has, in fact, an  $SO(2)$  symmetry mixing  $Q_e$  and  $Q_m$  for  $a = 4b$ , much like the stress-energy tensor derived from (28) for the same values  $a = 4b$ .<sup>1</sup> The effect of other four-derivative terms on the extremal black hole masses will be studied elsewhere [10].

There is another hint of a connection between our work and [9]. The superluminality/analyticity constraints were shown to be violated by the Dvali-Gabadadze-Porrati [11] brane-world model for modifying gravity in the IR. Interestingly, this model represents another example of trying to make interactions in the theory much weaker than gravity: the model has a 5D bulk with Planck scale  $M_5$ , but with a large induced Einstein-Hilbert action  $\int d^4x \sqrt{-g_{\text{ind}}} M_{\text{Pl}}^2 R^{(4)}$  on the brane. With  $M_{\text{Pl}} \gg M_5$ , this (quasi)-localizes gravity on the 4D brane. Again naively, there is nothing wrong with taking  $M_{\text{Pl}}$  large, as it seems to make the theory more weakly coupled; in this way it is similar to taking the limit of tiny gauge couplings in our examples, but we can here prove that the theory leads to superluminality and acausality in the IR, and is inconsistent with the standard analyticity properties of the  $S$ -matrix.

## 6 Discussion

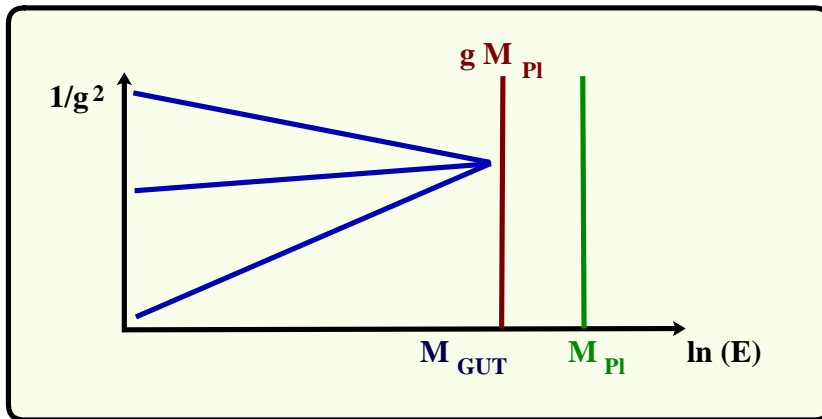
In this note, we have argued that there is a simple but powerful constraint on low-energy effective theories containing gravity and  $U(1)$  gauge fields. An effective field theorist would not see any problem with an arbitrarily weak gauge coupling  $g$ , but we have argued that

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<sup>1</sup>While our inequality  $M_{\text{extr}} < |Q|$  holds uniformly for  $a > 0, b > 0$ , we would also be able to satisfy our constraint and find a basis of directions where the inequality holds whenever at least one parameter ( $a$  or  $b$ ) is positive.

in fact there is a hidden ultraviolet scale  $\Lambda \sim gM_{\text{Pl}}$ , where the effective field theory breaks down, and that there are light charged particles with mass smaller than  $\Lambda$ . While this statement is completely unexpected to an effective field theorist, it resonates nicely with the impossibility of having global symmetries in quantum gravity, and the associated ability for large charged black holes to dissipate their charge in evaporating down to the Planck scale.

The specific forms of our conjecture are sharp, and if they are wrong it should be possible to find simple counter-examples in string theory, though we have not found any. The strongest form is that for the lightest charged particle along the direction of some basis vectors in charge space, the  $(M/Q)$  ratio is smaller than for extremal black holes. Such an assumption allows all extremal black holes to decay into these states. The weaker statement says that there should exist *some* state with mass/charge ratio smaller than for extremal black holes. In all the examples we have seen, this state has a “reasonably small” charge, so it is light; however, the weaker form allows the possibility that the smallest  $M/Q$  is realized for some large charge  $Q_*$  and objects that are “nearly” extremal black holes. While the number of exactly stable states would be finite in this case, it would still be extremely large. If this weaker form of the conjecture is true it is likely that there is some distribution of  $Q_*$  peaked for charges of order 1, but perhaps with sporadic exceptions at larger  $Q_*$ .



**Figure 5.** Because the gauge couplings at very high scales are smaller than one, our conjecture naturally predicts the existence of a new scale beneath the Planck scale.

If true, our conjecture shows that gravity and the other gauge forces can not be treated independently. In particular, any approach to quantum gravity that begins by treating pure gravity and is able to add arbitrary low-energy field content with any interactions is clearly excluded by our conjecture. Of course in string theory all the interactions are unified in a

way that makes treating them separately impossible. In particular, if we take the standard model gauge (augmented by SUSY or split SUSY or other particles leading to precision gauge coupling unification), we have perturbative gauge couplings at a very high energy scale, and our conjecture then implies that there *must* be new physics at a scale beneath the Planck scale, given by  $\Lambda \sim \sqrt{\alpha/G_N}$  which is close to the familiar heterotic string scale  $\sim 10^{17}$  GeV.

Our conjecture also offers a new experimental handle on ultraviolet physics, by searching for extremely weak new gauge forces. Indeed, if a new gauge force coupling to, say,  $B - L$  is discovered, with coupling  $g \sim 10^{-15}$ , in the current generation sub-millimeter force experiments, we would claim that there must be new physics at an ultraviolet scale  $\sim gM_{\text{Pl}} \sim \text{TeV}$ . Forces of this strength naturally arise in the context of large extra dimensions with fundamental scale near a TeV [12]; what is interesting is our claim that new physics *must* show up near the TeV scale.

It would be interesting to investigate whether there is an analogous conjecture in Anti-de-Sitter spaces, since here it can be translated into a statement about the spectrum of operators in the dual CFT that can perhaps be proved on general grounds.

Finally, it is interesting that the constraint implied by our conjecture seems to at least parametrically exclude apparently natural models for inflation based on periodic scalars with super-Planckian decay constants, which seem perfectly sensible from the point of view of a consistent effective theory. Of course, in the real world we don't need a parametrically large decay constant to get parametrically large numbers of  $e$ -foldings of inflation—60  $e$ -foldings will do. If the strong form of our conjecture is true, one might be tempted to conclude that there is a sharp obstacle to getting this sort of inflation in quantum gravity. If as is more likely the weaker form is true, then one might say that even though the low-energy theorists' notion of technical naturalness is misleading and such models are non-generic, there might be sporadic examples where they are possible. Clearly these are two very different pictures. The latter is more consistent with much of the philosophy of exploration in the landscape so far: things like a small cosmological constant are taken to be non-generic, tuned, but possible. But it is extremely interesting that phenomena of clear physical interest, like inflation with trans-Planckian excursions for the inflaton, which might even be forced on us experimentally by the discovery of primordial gravitational waves, seem to be pushing up against the limits of what quantum gravity seems to want to allow. Further exploration of the boundaries between the swampland and the landscape should shed more light on these issues.

## 7 Acknowledgements

We are grateful to Shamit Kachru, Megha Padi, and Joe Polchinski for discussions. We thank Jacques Distler for a discussion on the scale-dependence of our loose conjectures, and Juan Maldacena for clarifying discussion on the number of exactly stable states in quantum gravity. The work of NAH is supported by the DOE under contract DE-FG02-91ER40654. The work of LM is supported by a DOE OJI award. CV is supported in part by NSF grants PHY-0244821 and DMS-0244464.

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