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A D-moduli problem?

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Abstract

We point out a generic problem in string-inspired supergravity models with an anomalous $U(1)_X$. A large number of matter multiplets charged under $U(1)_X$ remain massless above the supersymmetry-breaking scale because of degeneracy of vacua solving the D-flatness conditions. A toy model is analyzed as an illustration of the mechanism; we find the surprising result that many scalars remain massless after supersymmetry-breaking in a hidden sector.

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Disclaimer

In this Letter we consider a simple toy model which illustrates a generic problem in string-inspired supergravity models with an anomalous $U(1)_X$. A large number of matter multiplets charged under $U(1)_X$ remain massless above the supersymmetry-breaking scale because of degeneracy of vacua solving the D-flatness conditions. We refer to these multiplets as "D-moduli." For example, in the model described in Section 4.2 of [1], that we will refer to here as the FIQS model, there are 26 massless chiral multiplets associated with this degeneracy. In the toy model considered here, we find that the degeneracy is partially broken by the introduction of a generic supersymmetry-breaking term in such a way that the overall vacuum energy vanishes. However many scalar fields remain massless even after supersymmetry breaking. Specifically, in our toy model with the convention[1] that $\mathrm{Tr}Q_X > 0$, where Q_X is the generator of $U(1)_X$, the only remaining flat directions are those for which a linear combination of fields with the lowest value q_0 of the $U(1)_X$ charge acquires a *vev*. The corresponding linear combination of chiral multiplets forms a massive vector multiplet with the $U(1)_X$ gauge fields, while the orthogonal combinations are massless. In the FIQS model, for example, the lowest $U(1)_X$ charge is -8 with a 15-fold degeneracy, so 14 complex scalars (as well as 26 chiral fermions) would be massless if no other symmetries were broken. The scalar fields with $q_X > q_0$ would acquire soft masses of the same order as generically expected for squarks and sleptons. However in string-derived models such as this one, the D-moduli are charged under other $U(1)$'s, which partially lifts the degeneracy of the vacuum before supersymmetry breaking.

We first consider our toy model in the context of standard supergravity with scalars and their superpartners in chiral multiplets. Then we appeal more explicitly to the linear multiplet formulation of gaugino condensation [2] as the mechanism of supersymmetry breaking, which we refer to as the BGW model. Neither of our models is realistic. For example, we neglect the dependence of the matter Kähler metric on T-moduli (*i.e.* breathing modes). While including this would considerably complicate the analysis, it seems unlikely to provide a mechanism (at least at tree level) for lifting the degeneracy of the vacuum. We also neglect additional, nonanomalous $U(1)$ couplings of the fields with large *vev*'s; as discussed below, we do not expect them to lift the vacuum degeneracy completely. We illustrate this point using the FIQS model, which itself should probably be considered a toy model for

[1]Our charge normalization is such that $\mathrm{Tr}T_a^2 = \frac{1}{2}$ if T_a is a generator of $SU(N)$.

reasons discussed below. Finally we suggest a possible mechanism for lifting the remaining degeneracy after supersymmetry breaking, and comment on implications for cosmology.

For our toy chiral supergravity model, we assume a Kähler potential

$$K = \sum_i |A_i|^2 + \sum_i |B_i|^2 + \sum_i |\Phi_i|^2 + K(M, \bar{M}), \tag{1}$$

a superpotential

$$W = \lambda_{ijk} A_i B_j \Phi_k + W(M), \tag{2}$$

and a gauge group $G_{\text{gauge}} = SU(3)_c \times U(1)_X$. We denote $U(1)_X$ charges by:

$$Q_X A_i = n_i A_i, \quad Q_X B_i = p_i B_i, \quad Q_X \Phi_i = q_i \Phi_i, \quad Q_X M_a = 0. \tag{3}$$

We assume $q_i \neq 0$ and $|q_i| \sim \mathcal{O}(1)$. We further assume that the q_i are such that $< D_X >= 0$ has a solution with $< A_i >=< B_i >= 0$; in other words we assume that there are flat directions that allow supersymmetry to remain unbroken in the absence of the nonperturbatively induced superpotential $W(M)$, as was found in the string-derived models studied in [1]. Scalar components of the gauge-charged superfields are given by $a_i = A_i|$, $b_i = B_i|$, $\phi_i = \Phi_i|$. We take A_i and B_i to be charged under $SU(3)_c$, while Φ_i and M_a are $SU(3)$ singlets. *E.g.*, A_i is a $\mathbf{3}$ while B_i is a $\bar{\mathbf{3}}$. Couplings such as (2) occur in semi-realistic heterotic orbifold models [1]. When $\langle\phi_i\rangle \neq 0$, generally required by D-flatness because of the Fayet-Illiopoulos (FI) term associated with [3] the anomalous $U(1)_X$, some color triplets acquire large masses. We demand that $SU(3)_c$ remain unbroken at all scales. Therefore

$$\langle a_i \rangle = \langle b_i \rangle = 0. \tag{4}$$

We interpret the gauge singlet superfields M_a as moduli; the supersymmetry-breaking superpotential $W(M)$ in (2):

$$\langle W(M)| \rangle \equiv \delta, \tag{5}$$

is assumed to be generated by nonperturbative dynamics. Cancellation of the $U(1)_X$ anomaly by the Green-Schwarz (GS) mechanism induces an FI term ξ in the $U(1)_X$ D-term [3]:

$$D_X = \sum_i q_i |\phi_i|^2 + \sum_i n_i |a_i|^2 + \sum_i p_i |b_i|^2 + \xi, \quad \xi = \frac{g_X^2}{24\pi^2} \mathrm{Tr} Q_X^3. \tag{6}$$

Motivated by semi-realistic models of string-derived effective supergravity with dynamical supersymmetry-breaking, we assume (in units where $m_P = 1/\sqrt{8\pi G} = 1$):

$$|\delta|^2 \ll |\xi| \ll 1. \tag{7}$$

In addition to D_X, we also have the $SU(3)_c$ D-term $D_c^{(r)}$, $r = 1, \ldots, 8$, which does not play a role in the following analysis. Let capital indices I, J, *etc.*, refer collectively to fields a_i, b_i, ϕ_i. Then the scalar potential of the toy model is given by:

$$\begin{aligned} V &= \frac{g_X^2}{2} D_X^2 + \frac{g_c^2}{2} D_c^{(r)} D_{(r)}^c + e^K \Big[\delta^{I\bar{J}} (W_I + W K_I)(\bar{W}_{\bar{J}} + \bar{W} K_{\bar{J}}) \\ &\quad + K^{a\bar{b}} (W_a + W K_a)(\bar{W}_{\bar{b}} + \bar{W} K_{\bar{b}}) - 3 W \bar{W} \Big]. \end{aligned} \tag{8}$$

We parameterize the vacuum value of the moduli sector F-term as

$$\left\langle K^{a\bar{b}} (W_a + W K_a)(\bar{W}_{\bar{b}} + \bar{W} K_{\bar{b}}) \right\rangle = \alpha |\delta|^2, \quad \alpha \sim 1. \tag{9}$$

The requirement of a vanishing cosmological constant gives

$$\begin{aligned} < V > &= \frac{g_X^2}{2} \left\langle D_X^2 \right\rangle + e^{<K>} |\delta|^2 \left(v^2 + \alpha - 3 \right) = 0, \\ \langle D_X \rangle &= \sum_i q_i |v_i|^2 + \xi, \end{aligned} \tag{10}$$

where $v_i = \langle \phi_i \rangle$ and $v = \sqrt{\sum_i |v_i|^2}$. Since V is gauge neutral, $\langle \partial V / \partial a_i \rangle$ and $\langle \partial V / \partial b_i \rangle$ are $SU(3)$-charged and vanish when (4) holds. The minimization condition for ϕ^i gives

$$< V_i > \equiv \left\langle \frac{\partial V}{\partial \phi_i} \right\rangle = 0 = \bar{v}_i \left[g_X^2 q_i \langle D_X \rangle + |\delta|^2 e^{<K>} (v^2 + \alpha - 2) \right]. \tag{11}$$

This implies that $v_i \neq 0$ for only one value of $q_i \equiv -q$, which, as we shall see, must be negative under our assumption (7). Note that (11) require $\langle D_X \rangle \sim |\delta|^2 \lll |W| >$, so that supersymmetry breaking is dominated by the moduli sector under our assumptions. (11) and (10) together imply

$$v^2 + \alpha - 3 = -O(|\delta|^2), \quad v^2 + \alpha - 2 \approx 1, \tag{12}$$

so that $\langle D_X \rangle > 0$. Next we consider the spectrum of Φ^i. The fermion mass matrix takes the form

$$(M_f^2)^i_j = v^i \bar{v}^{\bar{j}} \left[2 g_X^2 q^2 + e^{<K>} |\delta|^2 (v^2 + \alpha) \right], \tag{13}$$

and the scalar mass matrix (in Landau gauge) takes the form

$$
\begin{aligned}
M_s^2 &= \begin{pmatrix} M^2 & N^2 \\ (N^\dagger)^2 & (M^\dagger)^2 \end{pmatrix}, \\
(M^2)^i_j &= (M_f^2)^i_j - v^i \bar{v}^{\bar{j}} g_X^2 q^2 + \delta^i_j \left[g_X^2 q_i \langle D_X \rangle + e^{<K>} |\delta|^2 \left(v^2 + \alpha - 2 \right) \right] \\
(N^2)^{\bar{i}}_j &= \bar{v}^{\bar{i}} \bar{v}^{\bar{j}} \left[g_X^2 q^2 + e^{<K>} |\delta|^2 (v^2 + \alpha - 1) \right]. \qquad (14)
\end{aligned}
$$

In the absence of supersymmetry breaking, $\langle D_X \rangle = \delta = 0$, the superfield

$$
\Pi = \frac{1}{v} \sum_i \bar{v}^i \Phi^i \qquad (15)
$$

is eaten by the $U(1)_X$ gauge supermultiplet to form a massive vector multiplet. The orthogonal combinations

$$
D_\alpha = \sum_i c^i_\alpha \Phi^i, \quad \sum_i v^i c^i_\alpha = 0, \quad \sum_i \bar{c}^i_\alpha c^i_\beta = \delta_{\alpha\beta}, \qquad (16)
$$

are the massless D-moduli. When $\delta \neq 0$, the moduli M_a mix with Π; in other words the "Goldstone" chiral multiplets associated with supersymmetry breaking and with $U(1)_X$ breaking mix at order $|\delta|^2/g_X^2 q^2 v^2$. There is no mixing between the D-moduli and the M-moduli, so the D-moduli masses can be read directly from (13) and (14) by setting $v_i = 0$; one obtains $M^2_{Df} = 0$, and, using the vacuum conditions (10) and (11), for the scalars $d_\alpha = D_\alpha|$:

$$
m_\alpha^2 = (q_\alpha + q) g_X^2 < D_X > = \frac{(q_\alpha + q)}{q} \left[m^2_{\tilde{G}} + O(|\delta|^4) \right], \qquad (17)
$$

where

$$
m_{\tilde{G}} = e^{<K>/2} \delta \qquad (18)
$$

is the gravitino mass. From (17) it is clear that $m_\alpha^2 \geq 0$ if and only if $q = -\min(q_1, \ldots, q_{N_\Phi})$. The d_α that are linear combinations of ϕ^i with $q_i = -q$ remain massless, while the others acquire masses of order of the gravitino mass.

We now turn to a more specific model for supersymmetry breaking *via* gaugino condensation, as realized in the linear multiplet formulation for the dilaton [2]. Our model is an approximation to the BGW model, in that we neglect the moduli-dependence of the Kähler metric for the gauge-charged matter fields that we consider. With this approximation, the scalar potential takes the form[2]

$$
\begin{aligned}
V &= \frac{1}{2}g_X^2 D_X^2 + \sum_I \left|(W_I + K_I W)e^{K/2} + \beta u K_I\right|^2 \\
&\quad + f(\ell)\left|u(1+b_a) - 4\ell W e^{K/2}\right|^2 - \frac{3}{16}\left|u b_a - 4W e^{K/2}\right|^2, \qquad (19)
\end{aligned}
$$

where $u = e^{K/2}\tilde{u}(\ell, t)$ is the gaugino condensate which is determined by the equations of motion [2] as a function of the dilaton ℓ and the T-moduli t (treated here as constants[3]), $W = W(\phi)$, b_a is the β-function coefficient for the condensing gauge group, and the function $f(\ell)$ depends on the Kähler potential for the dilaton. The terms in (19) are, in a one-to-one correspondence, the counterparts in this model of the terms in (8). The vacuum conditions are, assuming $< W(\phi) >=< W_I >= 0$,

$$
\begin{aligned}
< V > &= \left\langle \frac{1}{2}g_X^2 D_X^2 + f(\ell)|u|^2(1+b_a)^2 - \frac{3}{16}|u|^2 b_a^2 + \beta^2|u|^2 v^2 \right\rangle \\
&\equiv \left\langle \frac{1}{2}g_X^2 D_X^2 \right\rangle + \hat{V} = 0, \\
< V_i > &= \bar{v}^{\bar{\imath}}\left(q_i \left\langle g_X^2 D_X \right\rangle + \hat{V} + \beta^2 < |u|^2 >\right). \qquad (20)
\end{aligned}
$$

In the spirit of the previous section, we assume (7) with $< |u| >\sim |\delta|$, so that $D_X \sim |\delta|^2$, $\hat{V} \sim |\delta|^4$. As in the previous example, one chiral supermultiplet with $q_i = -q$, the lowest $U(1)_X$ charge, forms a massive gauge supermultiplet with the $U(1)_X$ gauge superfield, while the remaining chiral superfields have massless fermions and scalar masses now given by

$$
m_\alpha^2 = (q_\alpha + q)g_X^2 < D_X >= \frac{(q_\alpha + q)}{q}\left[\left(\frac{4\beta}{b_a}\right)^2 m_{\tilde{G}}^2 + O(|\delta|^4)\right]. \qquad (21)
$$

When the GS term that cancels the T-duality anomaly is included, the parameter β is given by

$$
\beta = (p_\alpha - b_a)/4, \qquad (22)
$$

where p_α measures the coupling of the fields ϕ^i to the GS term. Here the situation is the same as for squarks and sleptons; if $p_\alpha = 0$, $m_\alpha \approx m_{\tilde{G}}$, while

[2]The full scalar potential for the BGW model is given in [4].

[3]When the dilaton is dynamical, $g_X \to g_X(\ell)$ and $g_X^2\xi \to 2\ell\xi$. (When nonperturbative string effects are neglected, $g_X^2(\ell) = 2\ell$ at the string scale.)

if the D-moduli couple to the GS term with the same strength b as the T-moduli, $m_\alpha \approx m_t/2 \approx |b/b_a - 1| m_{\tilde{G}}$. In the BGW model with $b = b_{E_8} \approx 10 b_a$ we get $m_\alpha \approx 10 m_{\tilde{G}}$. However, in the presence of Wilson lines that break the gauge group to a phenomenologically viable one, in general [5] $b < b_{E_8}$. In particular, in the FIQS model discussed below, with an $SO(10)$ condensing group, $b = b_a$, which implies that the moduli masses are much smaller than the gravitino mass, so the FIQS model is not viable in the context of the BGW supersymmetry breaking scenario.

As mentioned above, in more realistic models the D-moduli are charged under additional, nonanomalous gauge groups $U(1)_a$. Assuming affine level one, so that the gauge couplings are all equal at the string scale:[4] $g_a = g_X \equiv g$, the potential (8) or (19) takes the form

$$V = \frac{g^2}{2} \sum_a D_a^2 + \hat{V}. \tag{23}$$

where we have set to zero D-terms corresponding to nonabelian gauge groups such as $SU(3)$ in our toy model. The minimization conditions take the form

$$0 = < V_i > = g^2 \sum_a < D_a > q_i^a \bar{v}_i + < \hat{V}_i > . \tag{24}$$

In the models we are considering, we may write

$$< \hat{V}_i > = \bar{v}_i f, \tag{25}$$

where f is some function of the v^i and the moduli *vevs*'s. Then for any i such that $v^i \neq 0$, (24) implies:

$$0 = g^2 \sum_a D_a q_i^a + f = g^2 \sum_a q_i^a \sum_j |v^j|^2 + g^2 q_i^X \xi + f. \tag{26}$$

We are interested in models with F-flat and D-flat directions, *i.e.* in which the set of equations

$$D_a = 0 \tag{27}$$

has a solution along some F-flat direction. For example in the FIQS model, with $a = 1, \ldots 8$, (a continuous degeneracy of) solutions exist with nonvanishing *vev*'s for a set of 27 complex scalar fields[5]

$$\Phi^i \equiv S_\alpha^i,\ Y_A^i, \quad \alpha = 1, \cdots 5, \quad A = 1, \cdots 4, \quad i = 1, \cdots 3, \tag{28}$$

[4] Strictly speaking we should integrate out the heavy modes at the scale $v^i \sim .1$ and run the couplings their values to the condensation scale $< u^{\frac{1}{3}} > \sim 10^{-4}$; we neglect such renormalization effects here.

[5] Our notation differs slightly from that of [1].

which are charged only under the $U(1)_a$, with the charges q_a independent of the index i and

$$q_X^S = -8, \quad q_X^Y = 4, \quad q_a^{Y_1} = q_a^{Y_2}, \quad q_a^{Y_3} = q_a^{Y_4}. \tag{29}$$

The set of equations (27) have solutions that break 6 of the 8 $U(1)$'s, including $U(1)_X$, leaving unbroken the weak hypercharge of the Standard Model and one additional $U(1)$; in the effective low energy theory at scales $\mu \ll \xi$ there are no supermultiplets that carry both this latter $U(1)$ charge and Standard Model gauge charges [1]. Since 6 $U(1)$'s are broken, 6 of the supermultiplets in (28) are eaten by massive vector multiplets, and we are left with 21 D-moduli supermultiplets, instead of the 26 we would have in the absence of additional $U(1)_a$ charges for these fields.

Now consider the effect of the supersymmetry breaking term f in (24). The solution to (27) for $a \neq X$ requires that at least one field Y have a nonvanishing vacuum value. Therefore $D_{a \neq X} = 0$ is not a solution to (24) since the previous analysis without the additional $U(1)_a$ requires in this case that only $< S > \neq 0$ when $f \neq 0$. Hence we are led to solve[6] the set (26) of coupled equations for the $|v^i|^2$. We have analyzed these equations using the math package Maple and the $U(1)_a$ charge assignments of the FIQS model, and find that the minimum corresponds to

$$\begin{aligned} &Y_3^i = Y_4^i = 0, \quad \sum_i |S_\alpha^i|^2 = f_\alpha(\xi, g^2, f), \\ &\sum_i \left(|Y_1^i|^2 + |Y_2^i|^2\right) = f_Y(\xi, g^2, f), \end{aligned} \tag{30}$$

with an additional constraint of the form $f = f(g^2, \xi)$ to assure vanishing of the cosmological constant. Now Y_1^i and Y_2^i correspond to 12 real fields constrained by one equation to give 11 moduli, and each of the 5 choices of α in S_α^i correspond to 6 real fields subject to one constraint giving 5 moduli each. In this model there are 6 $U(1)$'s that get broken, so 6 moduli are eaten, leaving a total of

$$5 \times 5 + 11 - 6 = 30$$

D-moduli. Note that while there are fewer "light" D-moduli ($m \sim m_{\tilde{G}}$) than in the toy model with the same $U(1)_X$ charges but no additional $U(1)$'s, there

[6] Again we are oversimplifying; once supersymmetry is broken there is no reason to assume that the F-terms involving D-moduli couplings in the superpotential remain zero.

are actually more massless scalars (30 instead of 28) after supersymmetry breaking.

We remark that the first condition in (30) (which corresponds to $Y_3^{1i} = Y_3^{2i} = 0$ in the notation of [1]) has as a consequence that all of the down-type quarks are massless at tree level, and their masses must be generated by radiative corrections. Leaving aside the moduli problem alluded to above, this could be a phenomenological improvement of the model as compared with the solution of (27) in the absence of supersymmetry breaking. In that case there are both up- and down-type quark masses at tree level, but (unless unmotivated mixing is introduced in the Kähler potential) the CKM matrix is unrealistic: the heaviest up quark is not in the same $SU(2)_L$ gauge multiplet as the heaviest down quark. It has recently been shown that all down-type quark masses can be generated entirely from radiative corrections, subject to certain conditions on the high energy theory and supersymmetry breaking scenario [6]. Whether or not viable quark masses can be gotten by this mechanism in the FIQS model is under investigation.

In the generic gaugino condensation model of [2], supersymmetry breaking arises from the Veneziano-Yankielowicz part of the superfield Lagrangian:

$$\mathcal{L}_{\rm VY} = \int d^4\theta \frac{E}{8R} U \left[b' \ln(e^{-K/2}U) + \sum_\alpha b_\alpha \ln \Pi_\alpha \right] + \text{h.c.}\,. \tag{31}$$

The values of b' and b_α are determined by anomaly matching and are related to the β-function coefficient by

$$b_a = b' + \sum_\alpha b_\alpha. \tag{32}$$

The (weight two) chiral field U is the gaugino condensate superfield: $U| = u$, while the (weight zero) chiral fields Π are matter condensates. Condensation occurs provided there is also a superpotential for the matter condensates:

$$W(\Pi, T) = \sum_\alpha c_\alpha(T) \Pi_\alpha, \tag{33}$$

where the moduli-dependence of the coefficient assures modular invariance. In [2] it was assumed that Π is a composite operator containing fields charged only under the condensing gauge group. However in many models – such as the FIQS model with a hidden sector $SO(10)$ and matter in 16's – there is no operator that can be constructed from these fields alone that is invariant under the $U(1)_a$, and the coefficient c_α must depend on the Φ^i. It is possible that these additional couplings of the D-moduli are sufficient to lift the remaining degeneracy of the vacuum – and they may also give $O(m_{\tilde{G}})$ masses to the D-moduli fermions. This is because (25) would no longer hold, so that we do not get the single condition (26), but rather several independent conditions from (24). An analysis of this case requires a careful treatment of renormalization effects. However, it appears likely that any masses generated by these additional couplings will be governed by the supersymmetry breaking scale.

A large number of light scalar fields is problematic for cosmology in realistic models. The D-moduli have no gauge couplings in the effective low energy theory since they are charged only under the $U(1)$'s that are broken near the string scale. Therefore, unless they have unsuppressed superpotential couplings to relatively light particles, they are subject to the same constraints as, *e.g.,* the T-moduli [7]–[9]. The problem is somewhat alleviated if they couple to the Green-Schwarz term with $p_\alpha = b$ in (22) and $b \gg b_a$. In this case, like the T-moduli in the BGW model, their masses can exceed the gravitino mass by an order of magnitude. In Z_3 and Z_7 compactifications, with no T-moduli-dependent string threshold corrections [10], $b \geq b_a$ where the inequality is saturated if there are no twisted sector fields that are charged under the gauge group $\mathcal{G}_a$. This is the case, in particular for the FIQS model that we used above to illustrate the case of many $U(1)$'s. As we noted previously, this is not a viable option in the BGW context, since it gives unacceptably small moduli masses. In addition, models with $b \gg b_a$ alleviate [2] problems associated with dilaton cosmology.

It is plausible that the fermions can be sufficiently diluted by inflation to be harmless. The decay and annihilation of the particles of the Minimal Supersymmetric Standard Model suppresses the energy density of a decoupled massless fermion relative to a neutrino only by a factor of about 20, so 20 such fermions would contribute the equivalent of one neutrino species to the energy density during Nucleosynthesis if there is no other suppression mechanism. Fermions with order TeV masses would vastly overclose the universe unless they are inflated away or are sufficiently short lived. With only gravitational strength couplings, dimensional analysis suggests decay rates $\Gamma_\alpha \sim m_\alpha^3/m_P^2$, which were shown [11] to be marginally acceptable for the T-moduli and dilaton fermions. Any remnant of the broken $U(1)$'s in the tree-level couplings of the D-moduli would tend to suppress the decay rate. For example $U(1)$ invariant couplings would give rates $\Gamma_\alpha \sim m_\alpha^5/m_P^4$

that are unacceptably small without sufficient dilution. An estimate of the D-moduli fermion lifetimes requires a detailed understanding of the effective theory below the scale where the $U(1)_a$ are broken.

The point that we wish to emphasize here is that there are, generically, many more light moduli than have been previously considered, which may imply much stronger constraints on their masses and/or couplings. A full analysis of the D-moduli spectrum in the context of the BGW model for supersymmetry breaking, including dynamical T-moduli and dilaton, will be given elsewhere.

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