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Decoupling Dark Energy from Matter

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Abstract: We examine the embedding of dark energy in high energy models based upon supergravity and extend the usual phenomenological setting comprising an observable sector and a hidden supersymmetry breaking sector by including a third sector leading to the acceleration of the expansion of the universe. We find that gravitational constraints on the non-existence of a fifth force naturally imply that the dark energy sector must possess an approximate shift symmetry. When exact, the shift symmetry provides an example of a dark energy sector with a runaway potential and a nearly massless dark energy field whose coupling to matter is very weak, contrary to the usual lore that dark energy fields must couple strongly to matter and lead to gravitational inconsistencies. Moreover, the shape of the potential is stable under one-loop radiative corrections. When the shift symmetry is slightly broken by higher order terms in the Kähler potential, the coupling to matter remains small. However, the cosmological dynamics are largely affected by the shift symmetry breaking operators leading to the appearance of a minimum of the scalar potential such that dark energy behaves like an effective cosmological constant from very early on in the history of the universe.

Contents

1. Introduction

Cosmological observations [1–10] show evidence for an accelerated expansion, usually attributed to a new form of energy, dubbed dark energy. The simplest model for dark energy is the cosmological constant. From the point of view of particle physics, the cosmological constant is interpreted as the energy of the vacuum of the universe, which must be, according to observations, 120 orders of magnitude smaller than its "natural" value, the Planck energy scale. So far, the existence of a pure cosmological constant has been the most economical way of interpreting the observational data. Yet such a tiny cosmological constant is drastically at odds with particle physics [11] and therefore calls for a deeper explanation. Of course it could well be that the acceleration of the expansion of the universe is not due to dark energy but to a large scale modification of gravity. This possibility is under intense scrutiny (see e.g. [12, 13]). On the other hand, if correct, the observation of a tiny vacuum energy is all the more puzzling as the physics of phenomena at a such a low energy scale (below the electron-Volt) is well known and has been tested in laboratory experiments for at least 150 years. In particular, it would seem reasonable to treat dark energy using methods which have been so successful in describing particle physics from the atomic scale to the weak scale. This is realised in a large class of models where dark energy is attributed to a slowly rolling scalar field whose potential is of the runaway type [14–17]. The field value now is of the order of the Planck mass. Within the realm of high energy physics, this has prompted the use of supergravity where large field values can be handled [18–21].

In supergravity models of particle physics, two sectors are envisaged generically [22]. The so-called hidden sector breaks supersymmetry leading to a splitting of masses between the super-partners in the observable sector. In this setting the observable sector can be taken to be the Minimal Supersymmetric Standard Model (MSSM) whose phenomenology has been thoroughly studied and may be discovered at the Large Hadron Collider (LHC). Moreover supergravity models may play the role of low energy theory for a putative unified theory like string theory [23]. Dark energy must be included in this setting and is required to belong to a separate sector. This is to prevent the direct couplings between dark energy and baryons, which would lead to large deviations in tests of Newton's law [24–27] and the large discrepancy between the supersymmetry breaking scale and the vacuum energy scale.

One of the first problems in this approach is the need to have a good understanding of the theory close to its ultra-violet cut-off. Indeed, operators suppressed by the Planck scale may affect the dark energy predictions when the dark energy field reaches its present value. Another issue concerns the role of radiative corrections. Indeed the dark energy potential must be extremely flat and radiative corrections may lift the potential altogether. The flatness problem is particularly acute when dark energy couples to matter and the radiative corrections induce large field-dependent corrections. The non-field dependent corrections are also potent although they can be absorbed in the unknown cancellation of the bare cosmological constant. Another guise of the same problem is the potential presence of large deviations from Newton's law when the dark energy field is sufficiently light to mediate a fifth force between massive bodies. Two mechanisms can be envisaged in this case. Either a conspiracy leads to a small coupling of dark energy to matter at the present field value, the Damour-Polyakov effect [28], or the chameleon effect where a density dependence of the dark energy mass leads to a screening of the fifth force [29–31]. In the following we will address these issues in models where dark energy, the supersymmetry breaking sector and the MSSM interact only gravitationally. In these models, the coupling of dark energy to gravity is generically too large. We advocate that the use of a shift symmetry akin to the one solving the inflationary η problem alleviates this problem [32]. In inflationary models, the shift symmetry is often interpreted either as as axionic periodicity or as the result of the translational invariance of brane systems in extra-dimensional models [33]. We will not try to find the origin of the shift symmetry in the dark energy sector, leaving it for future work. We also study the effect of shift symmetry breaking operators and the resulting cosmological evolution. We find that models either lead to dark energy with a tiny coupling to matter thanks to the shift symmetry or the shift symmetry is broken and dark energy effectively behaves like a cosmological constant.

In the next section, we present the supergravity models of dark energy. In section 3, we consider a dark energy sector which exhibits a shift symmetry and analyse its effects on supersymmetry breaking, its coupling to the MSSM and study the deviations from Newton's law. We find that the presence of a shift symmetry is strongly motivated and gives a direct link between the dark energy potential and the superpotential of the model. In section 4, we break this shift symmetry by non-renormalisable operators in the Kähler potential and find that the resulting cosmology is akin to a pure cosmological constant. In section 5, we discuss briefly a different ansatz for the superpotential. Our findings are summarised in the last section.

2. Dark Energy and Shift Symmetry

We are interested in the coupling between matter, either baryonic or dark, and dark energy in the context of particle physics models connected to high energy physics. More precisely, we shall be concerned with supersymmetric models beyond the standard model of particle physics. These models provide a convenient framework within which one may study some aspects of string theory (the low energy effective action) or describe possible observable consequences at the LHC.

Supersymmetry has not yet been observed hence it must be broken at a scale which may be a few TeV in order to address the hierarchy problem between the Planck scale and the electro-weak scale. Supersymmetry is usually assumed to be broken in a hidden sector. In the following, we will concentrate on the possibility that the breaking of supersymmetry is transmitted to the standard model via gravitational interactions. Other mechanisms such as gauge mediation are popular and require a special treatment which is left for future work. Therefore, the usual structure of the standard model, the MSSM, is as follows. There are two sectors, the observable sector and the hidden sector and the assumption that the two sectors interact only gravitationally is expressed by [22]

$$K = K_{\rm MSSM}\left(\phi^a, \phi^{\bar{a}\dagger}\right) + K_{\rm h}\left(z, z^\dagger\right), \quad W = W_{\rm MSSM}(\phi^a) + W_{\rm h}(z), \tag{2.1}$$

where we restrict ourselves to a single field z in the hidden sector. Explicitly, the MSSM sector is defined by

$$K_{\rm MSSM} = \phi^a \phi^{\bar{a}\dagger}, \quad W_{\rm MSSM} = \frac{1}{3}\lambda_{abc}\phi^a\phi^b\phi^c + \frac{1}{2}\mu_{ab}\phi^a\phi^b, \tag{2.2}$$

where λ_{abc} are the Yukawa couplings and μ_{ab} has a single non-zero entry for the two Higgs bosons $\mu_{H_{\rm u}H_{\rm d}} \equiv \mu$. In this framework, the breaking of supersymmetry is parameterised by the F-term

$$F_z = {\rm e}^{\kappa_4^2 K_{\rm h}/2} D_z W_{\rm h} \tag{2.3}$$

at the minimum of the hidden sector potential, $\langle z \rangle = z_0$. In the above equation, we have used the definition $D_z W = \partial_z W + \kappa_4^2 K_z W$, where $\kappa_4^2 = 1/M_{\rm Pl}^2 = 8\pi G_{\rm N}$ and $K_z = \partial_z K$. In fact, the F_z term is defined in terms of the covariant derivative of the total superpotential. At high energy, the fields of the observable sector vanish and $W \simeq W_{\rm h}$. In the absence of dark energy, the bare cosmological constant is set to zero at tree level by requiring the cancellation at the minimum

$$|F_z|_{z_0} = \sqrt{3} m_{3/2}|_{z_0}, \tag{2.4}$$

where the gravitino mass is defined to be

$$m_{3/2} \equiv \kappa_4 W_{\rm h} {\rm e}^{\kappa_4^2 K_{\rm h}/2}. \tag{2.5}$$

Once supersymmetry is broken, the superpartners of the standard model fields acquire a mass. In the standard scenario, the running of the masses in the Higgs sector is such that the electroweak symmetry is broken at a low scale compared to the large scale, say the Grand Unified Theory (GUT) scale where supersymmetry breaking takes place. At the electroweak scale, the two Higgs fields pick up a vacuum expectation value (vev, which we denote by $v_{\rm u}$ and $v_{\rm d}$) such that the masses of the fermions become

$$m_{\rm u,d} = \lambda e^{\kappa_4^2 K/2} v_{\rm u,d}\,, \tag{2.6}$$

where λ is the appropriate Yukawa coupling for a particle of type u or d.

Then, the main issue is to understand how dark energy can be implemented into the above framework. Dark energy cannot belong to the observable sector as this would lead to a strong fifth force signal unless the coupling constants are artificially tuned to be small. We discard this possibility. Dark matter and dark energy could belong to the same sector. Couplings between dark matter and dark energy have been studied in the past (see e.g. [34–36]). However, in the MSSM, dark matter belongs to the observable sector. As a result, dark matter and dark energy will have only gravitational interactions. Finally, dark energy could belong to the supersymmetry breaking sector. In this work, we will assume that dark energy and the breaking of supersymmetry occur in separate sectors. This is motivated by the fact that supersymmetry breaking happens at a very large scale compared to the dark energy scale.

Therefore, one has to assume that there is a separate dark energy sector characterised by its own Kähler and superpotentials, $K_{\rm DE}$ and $W_{\rm DE}$. In the absence of the hidden sector and of the MSSM, the dark energy sector would be governed by a scalar potential

$$V_{\rm DE} = e^{\kappa_4^2 K_{\rm DE}} \left(|D_Q W_{\rm DE}|^2 - 3\kappa_4^2 |W_{\rm DE}|^2 \right), \tag{2.7}$$

which is a typical quintessence potential of the runaway type. However, as already mentioned, the dark energy sector cannot be considered as completely isolated since it always interacts gravitationally with the rest of the world.

All in all, our starting point will be the following separated Kähler- and superpotentials [37–39]

$$K = K_{\rm MSSM}\left(\phi^a, \phi^{\bar{a}\dagger}\right) + K_{\rm h}\left(z, z^\dagger\right) + K_{\rm DE}\left(Q, Q^\dagger\right), \tag{2.8}$$

$$W = W_{\rm MSSM}(\phi^a) + W_{\rm h}(z) + W_{\rm DE}(Q)\,, \tag{2.9}$$

where we restrict ourselves to a single field Q in the dark energy sector for convenience.

Let us now investigate what follows from the above equations. At high energy, the hidden sector field (still assumed to be stabilised) picks up a vev which is perturbed by the coupling to the dark energy sector. The perturbation is small enough (the dark energy scale is tiny compared to the supersymmetry breaking scale) to guarantee the existence of the perturbed minimum

$$\langle z \rangle = z_0(Q, Q^\dagger)\,. \tag{2.10}$$

Roughly speaking, this has two types of consequences. Firstly, the shape of the potential controlling the evolution of dark energy which, if the dark sector were isolated, would be given by Eq. (2.7) can be changed by supersymmetry breaking terms. Secondly, the gravitational physics at low energy is affected. The interaction between the hidden and dark sectors implies that all the soft breaking terms acquire a Q-dependent form. After renormalisation down to lower scales, the Higgs potential becomes Q-dependent, implying that both $v_{\rm u,d}$ become Q-dependent too. This has a drastic effect on the gravitational physics at low energy. Indeed the masses of the standard model fermions becomes Q-dependent

$$m_{\rm u,d}(Q) = \lambda e^{\kappa_4^2 K(Q,Q^\dagger)/2} v_{\rm u,d}(Q) \, , \tag{2.11}$$

where we have explicitly assumed that the Yukawa couplings are dark energy independent. Another important effect is the Q-dependence of the Quantum Chromo-Dynamics (QCD) scale Λ. Indeed the superpartners of the gauge bosons, the gauginos, acquire a mass only if the gauge coupling function f such that $\Re f = 1/g^2$ depends on z, see Eq. (2.24) in Ref. [37].

Let us combine all these results. If the dark energy potential is modified such that the mass of the quintessence field is now larger than 10^{-3}eV, then the problem with gravitational physics are evaded since the range of the corresponding force is too small to lead to effects that can be seen experimentally. However, in this case, the potential is not of the runaway form and the model is generally not interesting from the cosmological point of view. On the contrary, if the hidden sector is such that the dark energy potential preserves its shape, and, therefore, such that the model provides an interesting alternative to the cosmological constant, one encounters new problems. Indeed, in most runaway dark energy models, the dark energy field has a large value now, of the order of the Planck mass. Moreover the mass of the dark energy field is tiny and of the order of the Hubble rate now $H_0 \sim 10^{-43}$GeV $\ll 10^{-3}$eV. Such a low mass for a particle coupled to matter leads to gravitational problems and the existence of a detectable fifth force. This can be prevented if the coupling to gravity is small as imposed by the Cassini experiment [26]. Let us analyse the magnitude of the gravitational coupling of dark energy to matter. The gravitational coupling of a particle is simply given by [40]

$$\kappa_4 \alpha = \frac{{\rm d} \ln m}{{\rm d} Q_{\rm n}} \, , \tag{2.12}$$

where $Q_{\rm n}$ is the normalised dark energy field such that $\partial_Q \partial_{\bar{Q}} K_{\rm DE} (\partial Q)^2 = (\partial Q_{\rm n})^2/2$ and we have chosen $Q = Q^\dagger$ to be real.

The previous considerations would imply many low energy gravitational effects. In particular, one expects that the weak equivalence principle is violated at the microscopic level. Indeed, the masses of the u and d fermions have a different coupling $\alpha_{\rm u} \neq \alpha_{\rm d}$ implying that, in a gedanken experiment, particles of type u and d would fall at a different speed in a constant gravitational field. Of course, gravitational experiments are not carried out on microscopic particles but on macroscopic objects composed of many atoms. The mass of a particular atom can be decomposed as [40]

$$m_{\rm ATOM} \simeq M \Lambda_{\rm QCD} + \sigma' (N + Z) + \delta' (N - Z) + a_3 \alpha_{\rm QED} E_{\rm A} \Lambda_{\rm QCD} \, , \tag{2.13}$$

where $\Lambda_{\rm QCD} \simeq 180$MeV is the QCD scale, N the number of neutrons and Z the number of protons. The quantity M can be written $M = (N + Z) + E_{\rm QCD}/\Lambda$ where $E_{\rm QCD}$ is the strong interaction contribution to the nucleus binding energy. The number $E_{\rm A}$ is given by $E_{\rm A} = Z(Z-1)/(N+Z)^{1/3}$ and the quantity $a_3\alpha_{_{\rm QED}}\Lambda_{_{\rm QCD}}E_A$ represents the Coulomb interaction of the nucleus where $a_3\alpha_{_{\rm QED}} \simeq 0.77\times 10^{-3}$. Finally the coefficients δ' and σ' depend on the constituent masses and can be expressed as

$$\sigma' = \frac{1}{2}(m_{\rm u}+m_{\rm d})(b_{\rm u}+b_{\rm d}) + \frac{\alpha_{\rm QED}}{2}(C_{\rm n}+C_{\rm p}) + \frac{1}{2}m_{\rm e}\,, \tag{2.14}$$

$$\delta' = \frac{1}{2}(m_{\rm u}-m_{\rm d})(b_{\rm u}-b_{\rm d}) + \frac{\alpha_{\rm QED}}{2}(C_{\rm n}-C_{\rm p}) - \frac{1}{2}m_{\rm e}\,, \tag{2.15}$$

where $m_{\rm u} \sim 5$ MeV, $m_{\rm d} \sim 10$ MeV. The constants appearing in Eqs. (2.14) and (2.15) are given by: $b_{\rm u}+b_{\rm d} \simeq 6$, $b_{\rm u}-b_{\rm d} \sim 0.5$, $C_{\rm p}\alpha_{_{\rm QED}} \simeq 0.63$ MeV, $C_{\rm n}\alpha_{_{\rm QED}} \sim -0.13$ MeV. This implies that $\sigma'/\Lambda_{_{\rm QCD}} \simeq 3.8\times 10^{-2}$ and $\delta'/\Lambda_{_{\rm QCD}} \simeq 4.2\times 10^{-4}$.

The fact that $m_{\rm u}$ and $m_{\rm d}$ are dark energy-dependent quantities imply that the coefficients α' and δ', and hence $m_{_{\rm ATOM}}$, are now Q-dependent quantities. However, this is not the only source of Q-dependence. Indeed, the low energy gauge couplings are given by

$$\frac{1}{\alpha_i(m)} = 4\pi f_i - \frac{b_i}{2\pi}\ln\left(\frac{m_{\rm GUT}}{m}\right)\,, \tag{2.16}$$

where $i = 1,\cdots,3$ for $U(1)_{\rm Y}$, $SU(2)_{\rm L}$ and $SU(3)$ respectively with $b_i = (-33/5,-1,3)$. The quantity f_i is the gauge coupling function already discussed before. This implies that the QCD scale is related to the gauge coupling function as

$$\Lambda_{_{\rm QCD}} = m_{_{\rm GUT}}{\rm e}^{-8\pi^2 f_3/b_3}, \tag{2.17}$$

where we have assumed gauge coupling unification at $m_{_{\rm GUT}}$. Since f_3 is a function of Q, $\Lambda_{_{\rm QCD}}$ is also a Q-dependent quantity. The same reasoning is true for $\alpha_{_{\rm QED}}$ since

$$\alpha_{_{\rm QED}} = \frac{\alpha_2^2}{\alpha_1+\alpha_2}\,, \tag{2.18}$$

α_1 and α_2 being related to f_1 and f_2.

We are now in a position where one can estimate the typical gravitational coupling of an atom. From the previous considerations, one obtains

$$\begin{aligned}\kappa_4\alpha_{_{\rm ATOM}} = & -\frac{8\pi^2}{b_3}\frac{\partial f_3}{\partial Q_{\rm n}} + \frac{N+Z}{M}\frac{\partial}{\partial Q_{\rm n}}\left(\frac{\sigma'}{\Lambda_{_{\rm QCD}}}\right) + \frac{N-Z}{M}\frac{\partial}{\partial Q_{\rm n}}\left(\frac{\delta'}{\Lambda_{_{\rm QCD}}}\right) \\ & + a_3\frac{E_A}{M}\frac{\partial\alpha_{_{\rm QED}}}{\partial Q_{\rm n}}\,,\end{aligned} \tag{2.19}$$

where the variations of σ' and δ' read

$$\begin{aligned}\frac{\partial}{\partial Q_{\rm n}}\left(\frac{\sigma'}{\Lambda_{_{\rm QCD}}}\right) = & \frac{\kappa_4}{2\Lambda_{_{\rm QCD}}}(b_{\rm u}+b_{\rm d})\left(\alpha_{\rm u}m_{\rm u}+\alpha_{\rm d}m_{\rm d}\right) + \frac{\kappa_4}{2\Lambda_{_{\rm QCD}}}\alpha_{\rm d}m_{\rm e} \\ & + \frac{8\pi^2}{b_3}\frac{\sigma'}{\Lambda_{_{\rm QCD}}}\frac{\partial f}{\partial Q_{\rm n}} + \frac{C_{\rm n}+C_{\rm p}}{2\Lambda_{_{\rm QCD}}}\frac{\partial\alpha_{_{\rm QED}}}{\partial Q_{\rm n}}\,, \\ \frac{\partial}{\partial Q_{\rm n}}\left(\frac{\delta'}{\Lambda_{_{\rm QCD}}}\right) = & -\frac{\kappa_4}{2\Lambda_{_{\rm QCD}}}\left(b_{\rm u}-b_{\rm d}\right)\left(\alpha_{\rm u}m_{\rm u}-\alpha_{\rm d}m_{\rm d}\right) - \frac{\kappa_4}{2\Lambda_{_{\rm QCD}}}\alpha_{\rm d}m_{\rm e} \\ & + \frac{8\pi^2}{b_3}\frac{\sigma'}{\Lambda_{_{\rm QCD}}}\frac{\partial f}{\partial Q_{\rm n}} + \frac{C_{\rm n}-C_{\rm p}}{2\Lambda_{_{\rm QCD}}}\frac{\partial\alpha_{_{\rm QED}}}{\partial Q_{\rm n}}\,,\end{aligned} \tag{2.20}$$

with the coefficients $\alpha_{\rm u}$ and $\alpha_{\rm d}$ given by

$$\alpha_{\rm u} = \frac{\kappa_4}{2}\partial_{Q_{\rm n}} K_{\rm DE} + \frac{\kappa_4}{2}\partial_{Q_{\rm n}} K_{\rm h} + \frac{\kappa_4}{v_{\rm u}}\frac{{\rm d}v_{\rm u}}{{\rm d}Q_{\rm n}} \ , \tag{2.21}$$

$$\alpha_{\rm d} = \frac{\kappa_4}{2}\partial_{Q_{\rm n}} K_{\rm DE} + \frac{\kappa_4}{2}\partial_{Q_{\rm n}} K_{\rm h} + \frac{\kappa_4}{v_{\rm d}}\frac{{\rm d}v_{\rm d}}{{\rm d}Q_{\rm n}} \ , \tag{2.22}$$

and similarly the fine structure constant has a variation induced by the Q-dependence of the hidden sector vev $z_0\left(Q_{\rm n}\right)$

$$\frac{\partial \alpha_{\rm QED}}{\partial Q_{\rm n}} = 4\pi\left[\frac{\alpha_1^2\alpha_2^2 - (2\alpha_1+\alpha_2)\alpha_2^3}{(\alpha_1+\alpha_2)^2}\right]\frac{\partial f}{\partial Q_{\rm n}} \ . \tag{2.23}$$

Although complicated, the above expressions allow us to compute $\alpha_{\rm ATOM}$ exactly for the MSSM model. In fact, the analysis can be simplified if one notices that the leading effect in the gravitational coupling $\alpha_{\rm ATOM}$ comes from the bare dependence on the dark energy Kähler potential

$$\alpha_{\rm ATOM} \simeq \left[\frac{N+Z}{M}\frac{m_{\rm u}+m_{\rm d}}{4\Lambda_{\rm QCD}}\left(b_{\rm u}+b_{\rm d}\right) - \frac{N-Z}{M}\frac{m_{\rm u}-m_{\rm d}}{4\Lambda_{\rm QCD}}\left(b_{\rm u}-b_{\rm d}\right)\right] \times\kappa_4\partial_{Q_{\rm n}}K_{\rm DE} + \cdots \tag{2.24}$$

Despite the smallness of the quark masses compared to the QCD scale, the prefactor of $\kappa_4\partial_{Q_{\rm n}}K_{\rm DE}$ is no less than 10%. Let us now consider simple examples where the main trend can be grasped. For a canonically normalised field, we find that $K_{\rm DE} = QQ^\dagger$, leading to $Q_{\rm n} = \sqrt{2}Q$ and

$$\partial_{Q_{\rm n}}K_{\rm DE} = Q_{\rm n}/2 \ . \tag{2.25}$$

The main constraint on the presence of a fifth force comes from the Cassini probe. The Cassini experiment leads to a bound on $|\alpha_{\rm ATOM}| \lesssim 10^{-3}$ [26]. As we can see from the previous expressions, this has drastic consequences for dark energy. Indeed, this implies that the value of $Q_{\rm n}$ now must be less than $10^{-2}m_{\rm Pl}$ to satisfy the Cassini bound. In all dark energy models based on runaway potentials, the value of the quintessence field is of the order of the Planck mass now. The above expression shows that this leads to a strong violation of the Cassini bound. Another interesting case is provided by the dilaton ($n=1$) or moduli fields ($n=3$), with $K_{\rm DE} = -n\ln\left[\kappa_4\left(Q+Q^\dagger\right)\right]$ leading to $Q_{\rm n} = \sqrt{n/2}\ln\left(\kappa_4 Q\right)$ and

$$\partial_{Q_{\rm n}}K_{\rm DE} = \sqrt{2n} \ . \tag{2.26}$$

This implies again a large violation of the Cassini bound.

From the previous considerations, it is now clear that the fact that the gravitational coupling is large has a similar origin to the so called η problem in supergravity inflation, where a supergravity correction depending on the Kähler potential of the inflaton leads to a large mass for the inflaton. In the dark energy context, the gravitational coupling problem springs from the supergravity correction to the fermion mass which appears as an exponential of the Kähler potential. A solution to the η problem is provided by a shift symmetry implying that the Kähler potential of the inflaton vanishes along the inflationary direction. Similarly a solution to the gravitational coupling constant problem can be obtained provided the first derivative of the dark energy Kähler potential vanishes identically along the dark energy direction. In practice this implies that

$$K_{\rm DE}\left(Q,Q^\dagger\right) = K_{\rm DE}\left(Q-Q^\dagger\right) \ , \tag{2.27}$$

where we expand $K_{\rm DE}(x) = -x^2/2 + \cdots$ in powers of x. Notice that $Q \to Q + c$ with c real is a shift symmetry of the Kähler potential. In this case, the main contribution to the gravitation coupling vanishes. As a result the gravitational coupling depends only on the hidden sector dynamics and its coupling to the dark energy sector.

Let us also notice that another motivation for a shift symmetric Kähler potential is the presence of large corrections to the dark energy potential coming from the coupling to the hidden sector

$$\delta V_{\rm DE} = m_{3/2}^2 K_{\rm DE}^{QQ^\dagger}\partial_Q K_{\rm DE}\partial_{Q^\dagger}K_{\rm DE} \ . \tag{2.28}$$

For a canonical Kähler potential, this leads to a mass equal to the gravitino mass for the dark energy field. Such a large mass implies that the dark energy potential develops a minimum where the dark energy field is stuck very early on in the history of the universe, hence acting as an effective cosmological constant. In the moduli (dilaton) case, the correction to the potential is exponential with a large prefactor resulting in a large contribution to the energy density when the dark energy field is of the order of the Planck scale; such a large value needs to be compensated by one further tuning on top of the bare cosmological constant fine-tuning. When the Kähler potential is shift symmetric, the correction to the dark energy potential vanishes identically. Smaller corrections exist though. We will address the question of their origin in the following.

3. Implementing the Shift Symmetry

In this section, our goal is to recompute the coefficient $\alpha_{\rm ATOM}$ in the case where the dark energy sector is shift symmetric and to show that, in this case, the Cassini bound can be easily satisfied. Therefore, we take

$$K_{\rm DE}(Q) = -\frac{1}{2}\left(Q-Q^\dagger\right)^2 \ , \quad W_{\rm DE}(Q) = w(Q) \ , \tag{3.1}$$

where $w(Q)$ is, for the moment, an arbitrary function. The next step is to solve for the vev of the hidden sector field $\langle z\rangle$. To be completely explicit, let us focus on the coupling of dark energy to a Polonyi model with [41]

$$K_{\rm h}\left(z,z^\dagger\right) = |z|^2\,, \quad W_{\rm h}(z) = m^2(z+\beta)\,, \tag{3.2}$$

where $m^2 \sim m_{3/2}m_{\rm Pl}$. We focus on the real direction $Q = Q^\dagger$ as the imaginary direction is massive with a mass of order $m_{3/2}$. Along this direction, the scalar potential reads

$$\begin{aligned} V(Q,z) \,=\, & {\rm e}^{\kappa_4^2|z|^2}\left[|w'|^2 + |m^2\left(1+\kappa_4^2|z|^2\right) + \kappa_4 z^\dagger\left(\kappa_4 m^2\beta + \kappa_4 w\right)|^2\right. \\ & \left. -\ 3|m^2\kappa_4 z + \kappa_4 m^2\beta + \kappa_4 w|^2\right] \ , \end{aligned} \tag{3.3}$$

where the prime denotes a derivative with respect to Q. We are looking for the minimum of the scalar potential along the hidden sector direction z which is stabilised at

$$\kappa_4 z_0 = 1 - \kappa_4 \beta\,, \quad \kappa_4 z_0 = \sqrt{3} - 1\,, \tag{3.4}$$

in the absence of dark energy. When dark energy is present, the minimum is perturbed and becomes $z_{\rm min}(Q) = z_0 + \delta z(Q)$ where

$$\delta z(Q) = \left(\sqrt{3} - 1\right) \frac{w}{m^2}\,, \tag{3.5}$$

where we have neglected higher order terms in w and w' (see section 4). This perturbation is very small due to the discrepancy between the dark energy scale and the supersymmetry breaking scale. The potential at the minimum becomes a sole function of Q and can be expressed as

$$V_{\rm DE}(Q) = -2\sqrt{3}{\rm e}^{\kappa_4^2 |z_0|^2} \kappa_4 m^2 w(Q)\,. \tag{3.6}$$

A simple method to obtain this equation is to remark that $V = V(z_0) + V'(z_0)\delta z + V''(z_0)/2(\delta z)^2 + \cdots \simeq V(z_0)$ since $V'(z_0) = 0$ by definition and we work at first order in δz. In order to guarantee the positivity of the dark energy potential, one must impose $w < 0$. As a result, for any negative and runaway superpotential, we have found that there is a corresponding dark energy model with a potential energy proportional to the superpotential. Let us evaluate the order of magnitude of $w_{\rm now}$. We find

$$|w_{\rm now}| \sim \frac{\rho_{\rm cri}}{m_{3/2}}\,, \tag{3.7}$$

where $\rho_{\rm cri}$ is the present day critical energy density. Therefore, one can check that $\delta z/z_0 \simeq \rho_{\rm cri}/(m_{3/2}^2 m_{\rm Pl}^2) \ll 1$ (we have used the fact that the unperturbed $z_0 \simeq m_{\rm Pl}$) and this justifies the approximation made above. In fact, the previous calculation is an estimate of the value of the small dimensionless parameter used in order to perform the perturbative expansion. Indeed, the small parameter is then given by $\kappa_4 w/m^2 \simeq \rho_{\rm cri}/m_{\rm Pl}^4 (m_{\rm Pl}/m_{3/2})^2 \simeq 10^{-88}(100{\rm GeV}/m_{3/2})^2$. Notice that the dimensionless parameter w'/m^2 also appears in the calculation. It should be considered of the same order as $\kappa_4 w/m^2$ since $(\kappa_4 w/m^2)/(w'/m^2) \simeq \kappa_4 Q \simeq 1$ now for a runaway potential.

Let us give a simple example where the potential can be evaluated. The gaugino condensation superpotential for N_f flavours of quarks in the fundamental representation of the $SU(N_c)$ gauge group can be expressed as (the sign of the superpotential depends on a choice of the phase of the meson matrix) [42]

$$w(Q) = -(N_{\rm c} - N_{\rm f}) \frac{\Lambda^{(3N_{\rm c} - N_{\rm f})/(N_{\rm c} - N_{\rm f})}}{Q^{2N_{\rm f}/(N_{\rm c} - N_{\rm f})}}\,, \tag{3.8}$$

where Λ is a strong interaction scale and Q is the field along the diagonal meson direction. In this case, the dark energy potential is a Ratra-Peebles potential $V(Q) = M^{4+n}/Q^n$ where $n = 2N_{\rm f}/(N_{\rm c} - N_{\rm f})$ and $M^{4+n} = 2\sqrt{3}{\rm e}^{\kappa_4^2|z_0|^2}\kappa_4 m^2 \Lambda^{(3N_{\rm c}-N_{\rm f})/(N_{\rm c}-N_{\rm f})}$. The dark energy scale is completely specified by the supersymmetry breaking scale and the strong interaction scale.

It is also interesting to notice that, usually, when supergravity is used to construct models of dark energy, one does not obtain the Ratra-Peebles potential but the SUGRA potential [18], $V(Q) = \mathrm{e}^{\kappa_4^2 Q^2} M^{4+n}/Q^n$. Since the exponential correction directly originates from a supergravity term of the form $\mathrm{e}^{\kappa_4^2 K_{\mathrm{DE}}}$, it is clear that, in the presence of a shift symmetry, this term is not recovered. This is why, here, one obtains the Ratra-Peebles potential exactly. In fact, this is rather unfortunate since it is known that the dark energy equation of state of the SUGRA potential, thanks to the exponential factor, is much closer to -1 (more precisely $w_{\mathrm{DE}} \simeq -0.86$, see Ref. [18]) and, hence, more compatible with the present day observations, than the equation of state of the Ratra-Peebles potential. This last one is indeed too far from -1 to be compatible with the constraints on w_{DE} unless one considers very small values of n which seems pretty contrived. Therefore, the fact that, in the presence of a shift symmetry, one loses the exponential correction in the dark energy potential should be considered as a drawback. In other words, although the shift symmetry has solved the η-problem, contrary to the case of inflation this does not lead to desired features for dark energy. However, one can also consider other forms for $w(Q)$ such as [43]

$$w(Q) = \Lambda^3 \left[A + (\kappa_4 Q - B)^{\alpha}\right] \mathrm{e}^{-\lambda \kappa_4 Q}\,, \tag{3.9}$$

where Λ is an energy scale and A, B and α free parameters. This would lead to a potential of the Albrecht-Skordis type with interesting phenomenological properties such as a low value for the equation of state. Of course, almost any shape for the superpotential can be invoked as long as the resulting equation of state is low enough. Eventually, one would like to have an intrinsic justification for a given superpotential coming from a more fundamental theory.

As discussed above, the Q dependence of the atomic masses appears via the Q-dependence of the Higgs vev. In the supersymmetric context, the electroweak symmetric breaking is radiatively induced as the Higgs masses evolve from the GUT scale to the weak scale [44]. One of the Higgs masses becomes negative, triggering the symmetry breaking. Therefore, one needs to compute v_{u} and v_{d} in the shift symmetric case. At the GUT scale, where SUSY is broken, the observable potential is corrected by the soft supersymmetry breaking terms as follows

$$\begin{aligned} V_{\mathrm{MSSM}} &= \mathrm{e}^{\kappa_4^2 K} V_{\mathrm{susy}} + A_{abc}\left(\phi_a\phi_b\phi_c + \phi_a^\dagger\phi_b^\dagger\phi_c^\dagger\right) + B_{ab}\left(\phi_a\phi_b + \phi_a^\dagger\phi_b^\dagger\right) \\ &\quad + m^2_{a\bar{b}}\phi^a\phi^{\bar{b}\dagger}\,, \end{aligned} \tag{3.10}$$

where V_{susy} is the potential in the absence of supersymmetry breaking. In the shift symmetric case, the soft terms are explicitly given by

$$\begin{aligned} A_{abc} &= \frac{\lambda_{abc}}{3}\mathrm{e}^{\kappa_4^2|z_0(Q)|^2}\left[\left(M_{\mathrm{S}} + \kappa_4^2 w^\dagger\right)\kappa_4^2|z_0(Q)|^2 + \kappa_4^2 m^2 z_0(Q)\right]\,, \\ B_{ab} &= \frac{\mu_{ab}}{2}\mathrm{e}^{\kappa_4^2|z_0(Q)|^2}\left[\left(M_{\mathrm{S}} + \kappa_4^2 w^\dagger\right)\left(\kappa_4^2|z_0(Q)|^2 - 1\right) + \kappa_4^2 m^2 z_0(Q)\right]\,, \\ m^2_{a\bar{b}} &= m^2_{3/2}(Q)\delta_{a\bar{b}}\,, \end{aligned}$$

where the gravitino mass can be expressed as

$$m_{3/2}(Q) = \mathrm{e}^{\kappa_4^2|z_0(Q)|^2/2}|M_{\mathrm{S}} + \kappa_4^2 w|\,, \tag{3.11}$$

and we have defined

$$M_{\rm S} = \kappa_4^2 \left\langle W_{\rm h}\right\rangle (Q)\,. \tag{3.12}$$

A crucial point is already apparent here. The soft terms become Q-dependent but only through the Q-dependence of z_0 and $w(Q)$. In particular, in the non shift symmetric case, the soft terms are all proportional to $e^{\kappa_4^2 K_{\rm DE}} \simeq e^{\kappa_4^2 Q^2}$, see Eqs. (2.22), (2.23) and (2.24) of Ref. [37], which is responsible for the large violation of the Cassini bound. Here, thanks to the shift symmetry, this factor is absent.

Let us now specialise the above formula for the observable potential to the Higgs sector. One obtains

$$\begin{aligned} V_{\rm Higgs} =& \left(|\mu|^2 e^{\kappa_4^2|z|^2} + m_{H_{\rm u}}^2\right) v_{\rm u}^2 + \left(|\mu|^2 e^{\kappa_4^2|z|^2} + m_{H_{\rm d}}^2\right) v_{\rm d}^2 \\ &- 2\mu B v_{\rm u} v_{\rm d} + \frac{1}{8}\left(g_1^2 + g_2^2\right)\left(v_{\rm u}^2 - v_{\rm d}^2\right)^2 \,, \end{aligned} \tag{3.13}$$

where we have used that $m_{1\bar{1}}^2 = m_{H_{\rm u}}^2$, $m_{2\bar{2}}^2 = m_{H_{\rm d}}^2$ with $m_{H_{\rm u}} = m_{H_{\rm d}} = m_{3/2}$ at the GUT scales and $B_{ab} = \mu B \epsilon_{ab}$. The soft terms A_{abc} are hidden in the above formula and appear in the two loop expression for the renormalised Higgs masses, see Eqs. (3.14) and (3.15) of Ref. [37]. As already mentioned, all the soft terms depend on Q and, as they are renormalised to low energy, they keep an intricate Q-dependence. Then, one has to minimise the above potential in order to find the expression of $v_{\rm u}$ and $v_{\rm d}$. But, thanks to the shift symmetry, the Q-dependence of the minimum is determined through a complicated function of $z_0(Q)$ only. Therefore, at first order, the low energy minimum after electroweak symmetry breaking has to be given by

$$v_{\rm u,d}(Q) = v_{\rm u,d}^0 + C_{\rm u,d}\frac{w}{m^2}\,, \tag{3.14}$$

where $C_{\rm u,d}$ are coefficients of order one and $v_{\rm u,d}^0$ are the vevs in the absence of dark energy. This is one of the main results of this article: this explicitly determines the Q-dependence of the Higgs vevs in presence of dark energy. The coefficients $C_{\rm u,d}$ are complicated functions of the parameters of the model such as the gravitino mass, the gaugino mass etc .... Here, we do not need their explicit expressions.

As a result of the previous considerations, the masses of the atoms are

$$m_{\rm ATOM}(Q) = m_{\rm ATOM}^0 + C_{\rm ATOM}\frac{w}{m^2}\,, \tag{3.15}$$

where $C_{\rm ATOM}$ depends on the type of atom and is of order one while $m_{\rm A}^0$ is the atomic mass in the absence of dark energy. Again, such an expansion is entirely due to the expansion of $\langle z\rangle$. As a consequence, the gravitational coupling is

$$\alpha_{\rm ATOM} = C_{\rm ATOM}\frac{\partial_{\kappa_4 Q_{\rm n}} w}{m^2 m_{\rm ATOM}^0}\,. \tag{3.16}$$

For $Q_{\rm n} \sim m_{\rm Pl}$, this is negligible as it behaves like

$$\alpha_{\rm ATOM} \simeq \frac{w_{\rm now}}{m^2 m_{\rm ATOM}^0} \simeq 10^{-70}\left(\frac{m_{3/2}}{100\,{\rm GeV}}\right)^{-2}\left(\frac{m_{\rm ATOM}^0}{1\,{\rm GeV}}\right)^{-1}\,. \tag{3.17}$$

Hence, as announced, if the dark sector is shift symmetric, dark energy decouples from baryonic matter altogether.

Let us turn our attention to cold dark matter, which is composed of neutralinos. Using the same argument as before, the mass of the lightest supersymmetric particle can be expanded as

$$m_{\rm CDM} = m^0_{\rm CDM} + C_{\rm CDM}\frac{w}{m^2}\ , \tag{3.18}$$

where $C_{\rm CDM}$ is a dimensionless coefficient of order one. This implies that the effective potential due to the coupling between dark energy and dark matter can be written as

$$\begin{aligned} V_{\rm eff}(Q) &\equiv V_{\rm DE}(Q) + n_{\rm CDM} m_{\rm CDM} \\ &= n_{\rm CDM} m^0_{\rm CDM} + \left(C_{\rm CDM}\frac{n_{\rm CDM}}{m^2\kappa_4} - 2\sqrt{3}{\rm e}^{\kappa_4^2|z_0|^2}m^2\right)\kappa_4 w(Q)\ , \end{aligned} \tag{3.19}$$

where we have used the expression (3.6) of the dark energy potential. The number density of dark matter particles, $n_{\rm CDM}$, can be estimated as

$$n_{\rm CDM} = (1+z)^3\,\frac{\Omega_{\rm CDM}\rho_{\rm cri}}{m^0_{\rm CDM}}\ , \tag{3.20}$$

implying that the coupling between dark matter and dark energy plays a role for a redshift larger than

$$1+z \simeq \left(\frac{m^2_{3/2}m^0_{\rm CDM}}{m_{\rm Pl}H_0^2}\right)^{1/3}\ , \tag{3.21}$$

where $m \sim m^0_{\rm CDM} \sim 1$ TeV leads to $z \sim 10^{20}$, *i.e.* the coupling between dark energy and dark matter is always negligible.

Finally, we analyse the corrections to the scalar potential induced by the coupling of the dark energy field to matter since, so far, we have only considered the model at the classical level. Let us examine the loop corrections to the dark energy potential induced by the MSSM masses (bosons and fermions) which we parameterise as

$$m_i = m_i^0 + C_i\frac{w}{m^2}\,. \tag{3.22}$$

where C_i are species dependent constants. The Coleman–Weinberg potential gives the one-loop correction to the scalar potential and reads

$$V_{\rm 1loop} = \frac{1}{32\pi^2}{\rm Str}\left(m_i^2\right)\Lambda_{\rm c}^2 + \frac{1}{64\pi^2}{\rm Str}\left[m_i^4\ln\left(\frac{m_i^2}{\Lambda_{\rm c}^2}\right)\right]\ , \tag{3.23}$$

where $\Lambda_{\rm c}$ is the cut-off of the theory and the symbol "Str" denotes the sum over all the bosons minus the sum over all the fermions (one should not confuse m_i, the MSSM masses, with m, the SUSY breaking mass). The main correction comes from the quadratic divergence and reads

$$\delta V_{\rm DE} = \frac{1}{16\pi^2}{\rm Str}\left(m_i^0 C_i\right)\frac{\Lambda_{\rm c}^2}{m^2}w\,. \tag{3.24}$$

This term renormalises the scale appearing in the dark energy potential, while preserving the functional form of the potential, see Eq. (3.5). In other words, under the effect of the radiative corrections we find that a superpotential of the form $w(Q) = M_0^3 f(\kappa_4 Q)$, where f is a dimensionless function, becomes

$$w(Q) = \left[1 - \frac{1}{32\sqrt{3}\pi^2} \mathrm{e}^{-\kappa_4^2|z_0|^2} \mathrm{Str}\left(m_i^0 C_i\right) \frac{\Lambda_\mathrm{c}^2}{\kappa_4 m^4}\right] M_0^3 f(\kappa_4 Q) \,. \tag{3.25}$$

The correction term can be large if the cut-off scale is of order of the GUT scale. The main point is that the functional form of the potential is not modified and one can absorb the radiative correction in a redefinition of the scale M_0

$$M_0^3 \to \left[1 - \frac{1}{32\sqrt{3}\pi^2} \mathrm{e}^{-\kappa_4^2|z_0|^2} \mathrm{Str}\left(m_i^0 C_i\right) \frac{\Lambda_\mathrm{c}^2}{\kappa_4 m^4}\right] M_0^3 \,. \tag{3.26}$$

As in the usual renormalisation programme, the physical scale is the one including the radiative corrections and not the bare one appearing in the original Lagrangian. We do not know if this property can be extended to all loops.

To conclude this section, let us recap our main findings. We have shown that requiring the existence of a shift symmetry in the dark sector allows us to design a model where, at the same time, the runaway shape of the dark energy potential is preserved and the coupling between quintessence and the observable sector (ordinary and dark matter) is made negligible and, hence, compatible with the local tests of gravity. We have also shown that the shape of the potential is not modified by the quantum corrections, at least at one loop.

4. Breaking the Shift Symmetry

We have just seen that the existence of an exact shift symmetry implies an effective decoupling between dark energy and matter. A problem springs from the sensitivity of the model to higher order operators in the Kähler potential which break the shift symmetry. Therefore, we reconsider the calculation of the previous section but now relax the assumption that the Kähler potential is shift symmetric. Then, the new minimum of the potential is given by

$$\begin{aligned}\kappa_4 \delta z = & -\frac{1}{2\sqrt{3}} \Bigg\{ \sqrt{3}\kappa_4^2 \left(K_{\mathrm{DE}}^{-1}\right)^{Q^\dagger Q} \left(\partial_Q K_{\mathrm{DE}}\right)^2 + \Bigg[2\sqrt{3}\left(1-\sqrt{3}\right) \\ & + \left(2\sqrt{3}-1\right) \kappa_4^2 \left(K_{\mathrm{DE}}^{-1}\right)^{Q^\dagger Q} \left(\partial_Q K_{\mathrm{DE}}\right)^2 \Bigg] \frac{\kappa_4 w}{m^2} + \left(2\sqrt{3}-1\right) \\ & \times \kappa_4 \left(K_{\mathrm{DE}}^{-1}\right)^{Q^\dagger Q} \partial_Q K_{\mathrm{DE}} \frac{w'}{m^2} \Bigg\} \left[1 + \kappa_4^2 \left(K_{\mathrm{DE}}^{-1}\right)^{Q^\dagger Q} \left(\partial_Q K_{\mathrm{DE}}\right)^2\right]^{-1} \,. \end{aligned} \tag{4.1}$$

Several comments are in order at this point. Firstly, in the shift symmetric case, one has $\partial_Q K_{\mathrm{DE}} = 0$ in the dark energy direction and the above formula reduces to Eq. (3.5) as expected. Secondly, Eq. (4.1) is in fact universal in the sense that nothing has been assumed about the dark sector. The previous expression of the minimum relies on an expansion in $\kappa_4 w/m^2$ only. This expansion is an extremely good approximation as we have already shown that $\kappa_4 w/m^2 \simeq 10^{-88}$, thanks to the hierarchy between the SUSY breaking scale and the dark energy scale (*i.e.* the cosmological constant scale). Therefore, it represents a general expression for the position of the minimum in the Polonyi model in presence of dark energy regardless of the precise form of the dark sector. Thirdly, one notices in Eq. (4.1) the presence of a "zeroth order term", *i.e.* a term which is not proportional to $\kappa_4 w/m^2$ or to w'/m^2, namely $\sqrt{3}\kappa_4^2 \left(K_{\rm DE}^{-1}\right)^{Q^\dagger Q} \left(\partial_Q K_{\rm DE}\right)^2$. This means that, even if $\kappa_4 w/m^2$ is extremely small, there is still a correction originating from the dark sector Kähler potential. Moreover, since $\kappa_4 \delta z \ll 1$, as it was perturbatively determined, one must require for consistency that $\kappa_4^2 \left(K_{\rm DE}^{-1}\right)^{Q^\dagger Q} \left(\partial_Q K_{\rm DE}\right)^2 \ll 1$. Otherwise, one should solve numerically the higher order algebraic equation which controls the position of the new minimum.

Let us now determine the shape of the dark energy potential. Straightforward calculations lead to

$$V_{\rm DE}(Q) = m^4 {\rm e}^{\kappa_4^2\left(z_0^2+K_{\rm DE}\right)}\left\{\kappa_4^2 \left(K_{\rm DE}^{-1}\right)^{Q^\dagger Q} \left(\partial_Q K_{\rm DE}\right)^2 + \left[2\kappa_4^2 \left(K_{\rm DE}^{-1}\right)^{Q^\dagger Q} \right.\right.$$
$$\left.\left. \times \left(\partial_Q K_{\rm DE}\right)^2 - 2\sqrt{3}\right]\frac{\kappa_4 w}{m^2} + 2\kappa_4 \left(K_{\rm DE}^{-1}\right)^{Q^\dagger Q} \partial_Q K_{\rm DE}\frac{w'}{m^2}\right\}. \tag{4.2}$$

Again, this formula represents the expression of the dark energy potential in the most general case. In the shift symmetric case, it reduces to Eq. (3.6).

Having established the above general results, let us now focus on the breaking of shift symmetry. For this purpose, we now specialise the Kähler potential and write

$$K_{\rm DE} = -\frac{1}{2}\left(Q - Q^\dagger\right)^2 + \delta K_{\rm DE}\left(Q, Q^\dagger\right). \tag{4.3}$$

In order to perturbatively break the shift symmetry, we consider $\delta K_{\rm DE}\left(Q, Q^\dagger\right)$ to be a small correction in comparison to the shift symmetric zeroth order term. Then, it is easy to show that the new minimum of the potential is given by

$$\delta z \simeq \left(\sqrt{3} - 1\right)\frac{w}{m^2} - \frac{\kappa_4}{2}\left(\partial_Q \delta K_{\rm DE}\right)^2, \tag{4.4}$$

where we have written $\left(K_{\rm DE}^{-1}\right)^{Q^\dagger Q} = 1$ as the non-shift symmetric additional terms would give higher order corrections. This equation should be compared to Eqs. (3.5) and (4.1). Of course, if $K_{\rm DE}$ is shift symmetric, then the second term vanishes and one recovers Eq. (3.5). As already noticed, in order for the calculation to be consistent, the second term should be small $\kappa_4^2\left(\partial_Q \delta K_{\rm DE}\right)^2 \ll 1$ although not necessarily of the same order as $\kappa_4 w/m^2$. However, this is sufficient to neglect "second order" terms of the form "$\kappa_4^2\left(\partial_Q \delta K_{\rm DE}\right)^2 \times \kappa_4 w/m^2$". Then, the potential $V_{\rm DE}(Q)$ now contains only two parts and can be expressed as:

$$V_{\rm DE}(Q) \simeq -2\sqrt{3}m^4 {\rm e}^{\kappa_4^2 z_0^2}\left(\frac{\kappa_4 w}{m^2}\right) + m_{3/2}^2\left(\partial_Q \delta K_{\rm DE}\right)^2, \tag{4.5}$$

where $\kappa_4 z_0 = \sqrt{3} - 1$ and we have used the fact that $m^4\kappa_4^2 \simeq m_{3/2}^2$. The first part depends only on the superpotential w, and is nothing but Eq. (3.6), whereas the second part depends only on the Kähler potential $\delta K_{\rm DE}$. The second term has its origin in the breaking of the shift symmetry. In the following we will discuss the cosmological dynamics of the Q–field. Let us now be slightly more specific and write the corrections to the Kähler potential in the following form:

$$\delta K_{\rm DE} = c_p \frac{(Q+Q^\dagger)^p}{m_{\rm Pl}^{p-2}}\,. \tag{4.6}$$

In this case, the dimensionless parameter $\kappa_4^2\left(\partial_Q \delta K_{\rm DE}\right)^2 \simeq c_p^2 (Q/m_{\rm Pl})^{2p-2}$ which means that a choice such that $c_p \lesssim 10^{-2}$ is enough to guarantee the validity of our approximation up to $\kappa_4 Q \simeq \mathcal{O}(1)$. Furthermore, we assume that the first term in the potential above has the form M^{4+n}/Q^n. In this case, the resulting full potential reads

$$V_{\rm DE}(Q) = \frac{M^{4+n}}{Q^n} + m_{3/2}^2 m_{\rm Pl}^2 c_p^2 p^2 2^{2p-2}\left(\frac{Q}{m_{\rm Pl}}\right)^{2p-2}\,. \tag{4.7}$$

The potential possesses two branches: it goes as Q^{-n} for very small vevs and as Q^p for large vevs. In between there is a minimum and the overall runaway shape is lost. If we assume that this minimum is located at a vev which is small in comparison to the Planck mass, then we can approximate the denominator in the above expression by one. Under this assumption, the scalar field value at the minimum is

$$\left(\frac{Q_{\rm min}}{m_{\rm Pl}}\right)^{n+2p-2} \simeq \frac{M^{4+n}}{m_{3/2}^2 m_{\rm Pl}^{2+n}}\,. \tag{4.8}$$

The mass scale M can be fixed by assuming that, at the minimum at the present time, $V(Q_{\rm min}) \simeq \rho_{\rm cri}$, where we recall that $\rho_{\rm cri}$ is the critical density now. This gives

$$M^{4+n} \simeq \left(\frac{\rho_{\rm cri}}{m_{\rm Pl}^2 m_{3/2}^2}\right)^{n/(2p-2)} \rho_{\rm cri} m_{\rm Pl}^n\,. \tag{4.9}$$

Using this result in Eq. (4.8) and ignoring again numbers of order one, we obtain

$$\frac{Q_{\rm min}}{m_{\rm Pl}} \simeq \left(\frac{\rho_{\rm cri}}{m_{\rm Pl}^2 m_{3/2}^2}\right)^{1/(2p-2)} \simeq \left(\frac{H_0}{m_{3/2}}\right)^{1/(p-1)}\,. \tag{4.10}$$

This equation implies that the value of $Q_{\rm min}$ is much smaller than the Planck mass. This also justifies *a posteriori* our assumption to approximate the denominator in Eq. (4.7) by one.

Having determined the parameters of the model in the case where the shift symmetry is broken, let us now see whether the results of the previous section are preserved. Firstly, we notice that the mass of the quintessence field at the minimum is given by

$$m_Q^2 \simeq m_{3/2}^{2/(p-1)} H_0^{(2p-4)/(p-1)} \tag{4.11}$$

This mass is very small in comparison to 10^{-3}eV and, therefore, may lead to a fifth force when the dark energy field sits at the minimum. As a consequence, one needs to recompute the gravitational coupling. Since the shift symmetry is broken, one should now use Eq. (2.24). This leads to

$$\alpha_{\rm ATOM} \simeq \kappa_4 \partial_Q \delta K_{\rm DE} \simeq c_p \frac{H_0}{m_{3/2}} \simeq 10^{-44} c_p \left(\frac{m_{3/2}}{100\,{\rm GeV}}\right)^{-1} \ll 1\,. \tag{4.12}$$

This implies that the models are safe gravitationally and no fifth force is present. This formula should be compared to Eq. (3.17). As expected, we see that $\alpha_{\rm ATOM}$ is larger in the non shift symmetry case but the remarkable thing is that it is still very small. Therefore, the gravitational tests are still satisfied even if the shift symmetry is broken. A last remark is that, in principle, one should have computed the quantity $\alpha_{\rm ATOM}$ using the normalised field $Q_{\rm n}$ which does not coincide with Q due to the shift symmetry breaking term in the Kähler potential. However, for small Q (as is the case at the minimum, see above), the corrections are negligible and the field can effectively be considered as canonically normalised. Therefore, this justifies the previous calculation.

In fact, the only major effect of the shift symmetry breaking is the modification of the cosmological dynamics. As we will see in the following, the cosmological consequences of the setup just described has a lot in common with the theory presented in [38]. Let us discuss the evolution of the quintessence field. Assuming that the initial field value of the quintessence field just after reheating is much smaller than the minimum value (4.10), the potential is well approximated by an inverse-power law potential. As is well known, in this case there is a particular attractor solution, given by

$$Q_{\rm att} = Q_{\rm p}\left(\frac{a}{a_{\rm p}}\right)^{3(1+w_{\rm B})/(n+2)}, \tag{4.13}$$

where it is assumed that the background is dominated by a fluid with equation of state $w_{\rm B}$ and the scale factor a is a function of conformal time η, given by

$$a(\eta) = a_{\rm p}\left(\frac{\eta}{\eta_{\rm p}}\right)^{2/(1+3w_{\rm B})}. \tag{4.14}$$

The constant $Q_{\rm p}$ is given by

$$Q_{\rm p}^{-n-2} = \frac{18}{n^2 a_{\rm p}^2 \eta_{\rm p}^2 M^{4+n}}\frac{1-w_{\rm Q}^2}{(1+3w_{\rm B})^2} \approx \frac{\Omega_{\rm r}^0 z_{\rm reh}^4 m_{3/2}^{n/(p-1)}}{H_0^{n/(p-1)} m_{\rm Pl}^{n+2}}, \tag{4.15}$$

where we have used the expression of M given by Eq. (4.9) and chosen the time $\eta_{\rm p}$ to be the reheating time. In the above formula, $z_{\rm reh}$ is the reheating redshift and $\Omega_{\rm r}^0$ represents the present day radiation energy density. Finally, $w_{\rm Q} = (-2+nw_{\rm B})/(n+2)$ is the equation of state of the quintessence field. This solution is valid as long as $Q_{\rm att} \ll Q_{\rm min}$. At a certain time, however, $Q_{\rm att}$ is comparable to $Q_{\rm min}$. This happens when the scale factor reaches the value

$$a_{\rm min} \simeq a_{\rm p}\left[\Omega_{\rm r}^0 z_{\rm reh}^4 \left(\frac{Q_{\rm min}}{m_{\rm Pl}}\right)^2\right]^{1/[3(1+w_{\rm B})]}. \tag{4.16}$$

Using $\Omega_{\rm r}^0 = 10^{-5}$ and $z_{\rm reh} = 10^{28}$, this corresponds to a redshift

$$z_{\rm min} \simeq 10\left(\frac{Q_{\rm min}}{m_{\rm Pl}}\right)^{-1/2}. \tag{4.17}$$

For $m_{3/2} \simeq 1$ TeV, the last equation gives $z_{\rm min} \simeq 10^{13}$, so the field will sit at the minimum before big bang nucleosynthesis. In this case, the quintessence field has no dynamics and the model becomes effectively equivalent to a cosmological constant.

If the field is initially not on the attractor, there are two possibilities. The first one, called the undershoot case, is that initially $Q_{\rm ini} > Q_{\rm att}$. In this case, the field remains initially frozen at a value $Q = Q_{\rm ini}$ until $Q_{\rm ini} = Q_{\rm att}$. Using the expression for $Q_{\rm att}$, one obtains the redshift when this happens

$$z_{\rm u} \approx 10\left(\frac{Q_{\rm min}}{m_{_{\rm Pl}}}\right)^{n/4}\left(\frac{Q_{\rm ini}}{m_{_{\rm Pl}}}\right)^{-(n+2)/4} = 10\left(\frac{Q_{\rm min}}{m_{_{\rm Pl}}}\right)^{-1/2}\left(\frac{Q_{\rm min}}{Q_{\rm ini}}\right)^{(n+2)/4}. \tag{4.18}$$

Since $Q_{\rm min} > Q_{\rm ini}$ one has $z_{\rm u} > z_{\rm min}$ and, therefore, the attractor is always joined before the minimun is reached. This means that the above analysis is still valid in an undershoot situation.

The second case is the overshoot case, in which $Q_{\rm ini} < Q_{\rm att}$. In this case, the field is initially dominated by kinetic energy and evolves according to

$$Q = Q_{\rm ini} + m_{_{\rm Pl}}\sqrt{\frac{3\Omega_{Q_{\rm ini}}}{4\pi}}\left(1 - \frac{a_{\rm ini}}{a}\right). \tag{4.19}$$

As the kinetic energy is red-shifted away, at a certain point in time the potential energy becomes comparable to the kinetic energy. The field will then be frozen until it joins the attractor solution. The frozen redshift can estimated to be

$$z_{\rm froz} \simeq 10\left(\frac{Q_{\rm min}}{m_{_{\rm Pl}}}\right)^{n/6}\Omega_{Q_{\rm ini}}^{-(n+2)/12}. \tag{4.20}$$

However, the field could also never reach this regime because it "feels" the presence of the minimum before being frozen. When the field is dominated by kinetic energy, the minimum is felt at the following time

$$z_{\rm kin\rightarrow min} \simeq 10^{10}\Omega_{Q_{\rm ini}}^{-1/6}. \tag{4.21}$$

This time should be compared with $z_{\rm froz}$. One can show that $z_{\rm froz} > z_{\rm kin\rightarrow min}$ if $\Omega_{Q_{\rm ini}} < (H_0/m_{3/2})^{2/(p-1)}$. However, the point is that these two redshifts are large and that, in any case, the minimum is reached very early in the history of the Universe (*i.e.* before big bang nucleosynthesis). When the field is approaching the minimum, the potential is no longer of run-away form. One can show that the field oscillates around the minimum and that the oscillations are damped by the Hubble expansion [with a factor $\mathrm{e}^{-3(1-w_{\rm B})N_{\rm P}/2}$]. Indeed the situation is as described in [38], the only difference is the mass of the field.

To summarise this section, for generic initial conditions, the field will approach the minimum and settle there at very high redshift. The field will oscillate around the minimum and the model behaves essentially like the standard ΛCDM model, since the coupling to matter is suppressed in the theory presented so far. Indeed, very quickly, the oscillations are damped and the field stays at the minimum: the model is effectively equivalent to a cosmological constant.

5. A Different Ansatz

We have studied the coupling of dark energy to the standard model assuming that the three sectors- dark energy, supersymmetry breaking and the MSSM- are decoupled and only interact gravitationally, *i.e.* both the Kähler potential and the superpotential are

$$K = K_{\rm DE} + K_{\rm h} + K_{\rm MSSM}\,, \quad W = W_{\rm DE} + W_{\rm h} + W_{\rm MSSM}\,. \tag{5.1}$$

This implies that the standard model couplings become functions of the dark energy field measured in Planck units, *i.e.* if the gravitational interactions were turned off, no coupling between the sectors would exist. As a result, we have shown that the dark energy Kähler potential must be almost shift symmetric. The breaking of the shift symmetry can only occur via non-renormalisable interactions. In the shift symmetric case, no gravitational consequences of the existence of a nearly massless dark energy field can be detected and the cosmological evolution of the universe is of the uncoupled quintessence type. In the broken shift symmetric case, no gravitational effect can be detected either and the cosmological evolution is akin to the ΛCDM one. These results are intrinsically dependent on the initial ansatz for the Kähler potential and the superpotential. Introducing direct couplings between the dark energy sectors and the other two sectors would increase the gravitational effects unless the couplings were chosen to exactly cancel the gravitationally induced interactions. This would appear as an unnatural fine-tuning.

It happens that there is a hidden assumption in the way we have specified the model. Indeed, we have required that if we removed the dark energy sector, the models would reduce to the usual hidden sector symmetry breaking situation with two sectors coupled via gravitational interactions. If we were to dismiss this extra hypothesis, we could envisage a more general (and more drastic) type of decoupling. Indeed the supergravity Lagrangian is a function (forgetting the gauge sector) of $G = \kappa_4^2 K + \ln \kappa_4^6 |W|^2$. A possible decoupling consists in separating [45]

$$G = G_{\rm DE} + G_{\rm h} + G_{\rm MSSM} \tag{5.2}$$

This implies that

$$K = K_{\rm DE} + K_{\rm h} + K_{\rm MSSM}\,, \quad W = \kappa_4^6 W_{\rm DE} W_{\rm h} W_{\rm MSSM}\,. \tag{5.3}$$

From these expressions, one deduces that the scalar potential reads

$$V = {\rm e}^{\kappa_4^2 K} \left(V_{\rm DE} + V_{\rm h} + V_{\rm MSSM} - 3\kappa_4^2 |W|^2 \right)\,, \tag{5.4}$$

with

$$V_A = D_i W_A K_A^{i\bar{j}} \bar{D}_{\bar{j}} W_A^\dagger \tag{5.5}$$

where $D_i W_A = \partial_i W_A + \kappa_4^2 \partial_i K_A W$ for A running over the dark, hidden and observables sectors and fields ϕ_i in each sector respectively. Assuming a small value of the vacuum energy now such that $V_{\rm DE,now} \simeq \kappa_4^2 |W_{\rm DE}|^2$ for a dark energy field around the Planck scale leads to an upper bound on the gravitino mass

$$m_{3/2} \leq H_0, \tag{5.6}$$

obtained by saying that the superpotentials in the hidden and MSSM sectors must be less than the Planck scale cubed. This is a very low value for the gravitino mass which violates the Fayet bound on the gravitino mass $m_{3/2} \geq 10^{-5}$ eV [46]. Hence this approach does not seem to be promising. We can conclude that our original ansatz whereby the three different sectors are separated is the simplest setting to model dark energy in supergravity. Of course, more complex models could be built with particular couplings between the different sectors. It would be very interesting to construct such models explicitly when motivated by more fundamental theories such as string theory.

6. Conclusion

In this paper we have found a way of solving one of the outstanding problems associated with dark energy; namely how to incorporate it in a supersymmetric model of particle physics. We have taken a model of dark energy as a new sector on top of the usual observable and supersymmetry breaking sectors of particle physics phenomenology. We have found that the runaway shape of the dark energy potential is determined by the superpotential of the dark energy sector once a shift symmetry has been introduced. This prevents the existence of long range fifth forces. We have found a remarkable property: the shape of the dark energy potential is not modified by radiative corrections. The radiative corrections only modify the overall scale of the superpotential and one can absorb this modification in a redefinition of the bare superpotential. As a result, the runaway property of the dark energy potential is not destroyed by one-loop radiative corrections. If the shift symmetry is exact then the dark energy retains its properties when incorporated into a supersymmetric model. This is reminiscent of the solution to the hierarchy problem which prompted the emergence of supersymmetric models whereby the sensitivity of the Higgs mass to large scales disappears. This is due to the cancellation of bosonic loops by fermionic loops in the perturbative expansion. Of course, when supersymmetry is broken this cancellation is no longer exact. Similarly, when the shift symmetry is not exact, the presence of higher order corrections to the Kähler potential induce a drastic modification of the scalar potential. Indeed, it is not of the runaway type anymore but has a minimum implying that the model behaves essentially like a Λ-CDM model since very early in the history of the universe. Hence the existence of a shift symmetry is crucial for both the stability of dark energy to radiative corrections and the compliance with gravity tests. As shift symmetries can be motivated in string theory– they already appear as a solution to the η problem of inflation–, one may hope that dark energy models may be constructed using stringy ingredients. In this case, a proper understanding of the dark energy superpotential would be crucial: it would lead to a direct comparison with observations such as the measurement of the dark energy equation of state and its time dependence.

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