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CHAPTER 1

Introduction

1.1 The standard model of cosmology

For a long time it has been realised that the observed darkness of the night contradicts an eternal infinite universe. The first known account of the dark night paradox, most commonly known as Olbers' paradox (Olbers, 1823), is credited to Digges (Digges, 1576) and Kepler (Kepler, 1610). Later accounts include work by Halley (Halley, 1720a,b) and Chesaux (Cheseaux, 1744). The first attempt on a resolution to this problem is by Lord Kelvin (Kelvin, 1901). For a historic account including translated reprints of the cited articles, see Harrison (1987).

In the following years, progress was much more swift. In 1915, Einstein published his papers concerning the theory of general relativity (Einstein, 1915, 1916), which point to a universe that is unstable against collapse, and thus must either be expanding at a declining rate or collapsing at an increasing rate. Einstein perceived this as a problem — he assumed the universe was stationary — and introduced a cosmological constant to solve this (Einstein, 1917). However, a cosmological constant still does not allow for a stable stationary universe, as shown by de Sitter in 1918. In two papers, de Sitter analysed the behaviour of empty universes and showed that any of those universes only has unstable fixed points (de Sitter, 1918a,b).

In the years following, Friedmann (1922), Lemaître (1927)[1], Robertson (1929) and Walker (1933) wrote down a maximally symmetric metric that does allow for

[1] English translation: Lemaître (1931)

expansion or collapse by an overall scale factor $a(t)$,

$$\mathrm{d}s^2 = \mathrm{d}t^2 + a(t)^2 \mathrm{d}\mathbf{x}^2 \ , \tag{1.1}$$

and used Einstein's field equations to solve for the evolution of $a(t)$,

$$H^2 = \left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G}{3}\rho(t) - \frac{kc^2}{a^2} + \frac{\Lambda c^2}{3} \ , \tag{1.2}$$

where c is the speed of light and k is a parameter that describes the spatial curvature of the universe, $k = 0, \pm 1$ for a flat, positive or negatively curved spatial hypersurface. Furthermore, Λ parametrises the cosmological constant Einstein introduced to allow stable cosmological solutions with matter and ρ describes the energy density. In this equation the Hubble parameter $H = \dot{a}/a$ is introduced, which is the rate of expansion of the universe. Its current best estimates include 70.6 ± 3.1 (km/s)/Mpc (Suyu et al., 2010) and $70.4^{+1.3}_{-1.4}$ (km/s)/Mpc (Jarosik et al., 2010). The Hubble parameter has the units of inverse time, H^{-1} is therefore a timescale, the Hubble time, which is approximately the age of the universe. Multiplying with the speed of light gives the Hubble radius cH^{-1}, which is the maximum distance anything can have travelled in a Hubble time.

The evolution of ρ can be found from the second Friedmann equation

$$\frac{\ddot{a}}{a} = -\frac{4\pi G}{3}\left(\rho + \frac{3p}{c^2}\right) + \frac{\Lambda c^2}{3} \tag{1.3}$$

together with an equation of state $p = f(\rho)$ that relates the energy density to a pressure p. We will take a linear relation $p = w\rho$, with w an arbitrary parameter, as this suffices to describe the cases of interest in this chapter. In fact, with this linear relation also curvature and the cosmological constant can be defined in terms of an effective energy density parameter ρ and effective pressure p, see table 1.1. Then all the contributions ρ_{matter}, $\rho_{\text{radiation}}$, ρ_k and ρ_Λ have to sum to $3H^2/(3\pi G) = \rho_c$, which allows one to write the Friedmann equation (eq. 1.2) as

$$1 = \sum_i \frac{\rho_i}{\rho_c} \tag{1.4}$$

where i labels the different components. The first experimental evidence of a non-stationary universe came in 1929, when Hubble showed that almost all galaxies are redshifted and thus move away from us, and that the redshift is positively correlated with the distance of the galaxy. Much more convincing evidence came in 1965 when Penzias and Wilson found the Cosmic Microwave Background. This background was already predicted by Gamow (1948a,b) and Alpher and Herman (1948) and is created in the early universe, about 380,000 years after the big bang. In an expanding universe, temperatures are higher in the past, and before 380,000 years after the big bang the temperatures were too high for neutral hydrogen to exist. As photons are very effectively scattered by free electrons, the mean free path of a photon was much shorter than the Hubble radius, cH^{-1}. Around 380,000 years after the big bang, the recombination time, the universe cooled down enough for neutral hydrogen to form. Neutral hydrogen scatters photons much less efficient, so suddenly the mean free path of photons became much larger than the Hubble radius. As a result, the universe became approximately transparent for photons. These photons, redshifted by a factor $T_{\mathrm{now}}/T_{\mathrm{formation}} \sim 1100$, now form the 2.73 K microwave background. Due to the finite lifetime of our universe, not all lines of sight do end on stars. Yet, in an expanding universe of nonzero temperature one expects an all-sky background from the recombination surface. In contrast to the all-sky stellar radiation from an eternal universe, this background exists and thus provides strong evidence for an expanding universe.

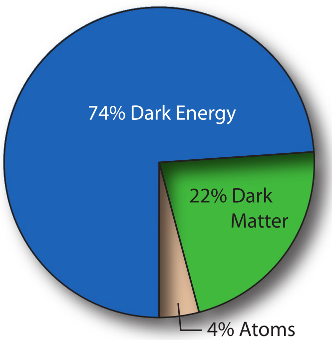


Figure 1.1: *The current measurements on the energy density components of our universe (image courtesy: LAMBDA, NASA.)*

Furthermore, from (eq. 1.3) it is clear that the rate of expansion of the universe should decrease ($\ddot{a} < 0$) for energy density components with $w > -1/3$, while it increases when $w < -1/3$. Recent evidence (Perlmutter et al., 1998, 1999, Riess et al., 1998, Spergel et al., 2003, Komatsu et al., 2010) shows the Hubble parameter H is currently increasing, which means that a cosmological constant or some other form of dark energy must actually give a large contribution to the total energy density.

Cold matter	$w = 0$
Radiation	$w = 1/3$
Curvature	$w = -1/3$
Cosmological constant	$w = -1$

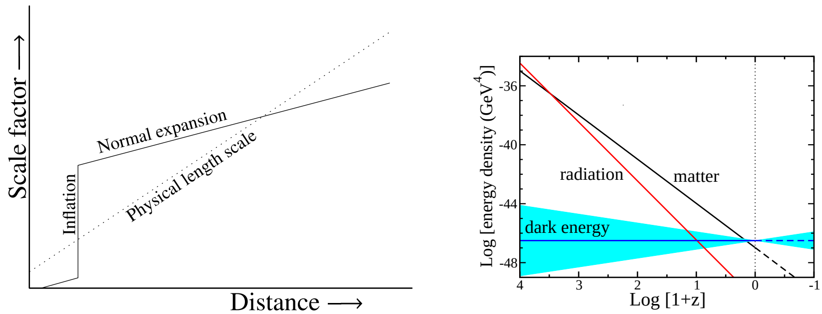
Table 1.1: *Equation of state for the different components of energy density.*

Besides dark energy, there is also a missing energy component that satisfies the matter equation of state, dark matter. From a wide range in scales it is known that normal matter cannot explain the observed gravitational fields. At small scales rotation curves in galaxies (Rubin and Ford, 1970) do not match the visible matter distribution. At intermediate scales, peculiar motion in clusters (Zwicky, 1933) and gravitational lensing (see eg. Bacon et al., 2000, Refregier, 2003 for constraints from weak lensing from the large scale structure) point to a mass component that is not visible. At large scales, the formation of the large scale structure (see Springel et al., 2005 for an overview of the simulations, Tegmark et al., 2006, Reid et al., 2010 for measurements) and the observations of the Cosmic Microwave background (Spergel et al., 2003, Komatsu et al., 2010) need a dark matter component to understand the data. We know that this dark matter component cannot be baryonic, as the baryonic matter component can be determined via primordial nucleonsynthesis (Kernan and Krauss, 1994, Copi et al., 1995, Coc et al., 2004). Furthermore, modified gravity with baryonic matter does not fit lensing observations of the Bullet cluster (Clowe et al., 2004) and the Cosmic Microwave Background data (Komatsu et al., 2010, for a nice theoretical discussion see Mukhanov, 2004). These observations point to an invisible component that must have weak interactions with all matter in the universe[2], and certainly does not interact with light. Combining the current knowledge about energy density components, we find that the universe is, within experimental bounds, spatially flat (Komatsu et al., 2010) and dominated by dark matter and dark energy (see figure 1.1).

In the Friedmann equations (eqs. 1.2 and 1.3) above the evolution of the universe is such that, for $w > -1/3$, the Hubble radius is growing faster than the distance between causally connected points, that is points that have been in causal contact. Our universe is currently dark energy dominated, but back in time it certainly was dominated by matter and earlier radiation (figure 1.2(b)). Current observations of

[2]This does not mean that dark matter is charged with respect to the the weak interaction, although a very popular candidate, the "Weakly Interacting Massive Particle", is.

the Cosmic Microwave Background show correlated perturbations up to the size of the visible universe. The Hubble horizon at the time of decoupling, and thus the maximum distance between causally connected events, is 380,000 light years, which corresponds to about one degree on the current sky (approximately the first peak in the angular power spectrum, figure 1.3). This means that the Microwave Background consists of approximately 40,000 patches that were causally disconnected at the time of decoupling. Furthermore, the current measured flatness, isotropy and homogeneity of the universe (Spergel et al., 2003, Komatsu et al., 2010) also violates causality or requires large amounts of finetuning, while the lack of phase transition defects (eg. magnetic monopoles) is unnatural. All these problems can be solved by introducing a period of inflation, as proposed by Guth in 1980, which will be the topic of the next paragraph.[3]



(a) Evolution of Hubble radius (solid line) and a physical distance scale (dotted line), such as the separation between two points.

(b) Evolution of the different energy density components as a function of redshift (image from Frieman et al., 2008).

Figure 1.2:

1.2 The inflationary paradigm

The solution proposed in Guth (1981) is to use the observation that for energy contributions with an equation of state parameter $w < -1/3$ the Hubble radius grows slower than the distance between causally connected points. If a period with $w < -1/3$,

[3]Isotropy is actually not solved by inflation, as different dimensions could inflate at different rates. String theory compactifications, discussed in section 1.3.1, actually need anisotropic inflation, as isotropic inflation would also inflate the hidden dimensions.

called an inflationary period, lasts only temporarily, it allows a region that is causally connected before the onset of inflation to cover a causally disconnected region after the inflationary period. When inflation lasts long enough, our current observable universe fits entirely within a region that was causally connected before inflation occurred (figure 1.2(a)). Thus, inflation can solve finetuning problems concerning the homogeneity of the visible universe. Furthermore, during the period of inflation the accelerated expansion reverses the growth of curvature and dilutes matter, thus solving the flatness problem and the issues regarding the absence of cosmological defects. A negative equation of state can be created by a field that has a small kinetic energy term and a large potential energy term. For $w = -1$ the energy density per unit volume is constant when the universe expands. A vacuum energy, generated by a nonzero vacuum expectation value of a scalar field, is so far the only field theoretical mechanism that can provide an equation of state with a negative w. The model proposed by Guth used quantum tunneling from an inflating state to the current vacuum to end inflation. Later work by Linde (1982) and Albrecht and Steinhardt (1982) found a new inflation paradigm that used a continuous potential along which the inflaton slowly rolls down.

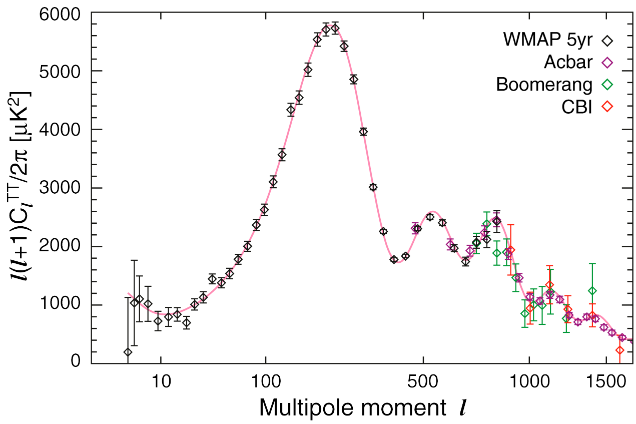


Figure 1.3: *Power spectrum of the angular cross correlation function of the temperature of the Cosmic Microwave Background. The red line is the best fit, based on WMAP-data alone, to the ΛCDM model (Dunkley et al., 2009). This fit agrees very well with the higher multipole data. Image from Nolta et al. (2009).*

In 1981 Mukhanov and Chibisov showed that perturbations of the metric and the inflaton field lead to perturbations with a power spectrum that depends on the potential of the inflaton field (see also Hawking, 1982, Starobinsky, 1982, Guth and Pi, 1982, Bardeen et al., 1983, Mukhanov, 1985). The power spectrum should be exactly scale invariant for $w = -1$, but in order to explain why inflation could end we need an inflaton field that evolves to lower energies. Therefore, a slightly red-tilted power spectrum is predicted. In 1991 the COBE satellite (Smoot et al., 1992) first observed these perturbations. In 1999, the ground based Toco experiment (Torbet et al., 1999) found evidence for the first acoustic peak, later confirmed by the BOOMERanG (Melchiorri et al., 2000) and MAXIMA (Hanany et al., 2000) balloon experiments. In 2001, the quality of the data improved significantly with data from the WMAP satellite (Spergel et al., 2003). Later, the WMAP 5-year data Hinshaw et al. (2009), together with data from balloon experiments (ACBAR (Reichardt et al., 2009), BOOMERanG (Jones et al., 2006) and CBI Readhead et al. (2004)) provided an accurate map of the perturbations to small angular scales (figure 1.3), and confirmed the red tilt of the power spectrum (Dunkley et al., 2009). The current best estimate is from the WMAP 7-year data (Komatsu et al., 2010), which is compatible with an approximately power law power spectrum $P(k) \propto k^{n_s-1}$, with $n_s = 0.963 \pm 0.012$ at 68% CL. Upcoming improved measurements are expected from the ground based Atacama Cosmology Telescope (see Fowler et al., 2010 for the first results) and the Planck satellite (Planck collaboration, 2006), whose results are expected late 2012 or early 2013.

In order to solve the problems that inflation was invented for, inflation needs to last for at least 55 H^{-1}, in which the universe thus expanded by e^{55} orders of magnitude or 55 e-folds. In order to have so many e-folds the potential energy is allowed to vary only very slowly. Being consistent with a slightly red tilted power spectrum, the most common mechanism to generate such a slow variation is called slow-roll inflation. Non-inflationary mechanisms that are compatible with the observations of the power spectrum include gauge/cosmology duality (McFadden and Skenderis, 2010a,b), the ekpyrotic scenario (Khoury et al., 2002, Steinhardt and Turok, 2001, 2002) or string gas cosmological models (Brandenberger and Vafa, 1989, Brandenberger et al., 2004). Also, an entirely causal mechanism has been proposed (Turok, 1996).

1.2.1 Slow-roll inflation

Single field slow-roll inflation successfully accounts for many of the observed properties of the cosmic microwave background (CMB), including the near scale invariance of the power spectrum of the primordial density fluctuations that seed the observed CMB anisotropies (Mukhanov and Chibisov, 1981). Slow-roll inflation in its simplest form (Linde, 1982, Albrecht and Steinhardt, 1982) still marginally fits the current observational bounds (Komatsu et al., 2010, Larson et al., 2010). Furthermore, all inflation models developed since these earliest models depend on the slow-roll approximation (Steinhardt and Turner, 1984, Salopek and Bond, 1990, Liddle and Lyth, 1992. See Lyth and Liddle, 2009 for a recent text-book discussion) that the Hubble rate is varying only slowly. In terms of the Hamilton-Jacobi slow-roll variables the requirement is that

$$\begin{aligned} \epsilon_H &= -\frac{\dot{H}}{H^2} \ll 1 \quad \text{and} \\ \eta_H &= -\frac{\ddot{\phi}_0}{H\dot{\phi}_0} \ll 1 \ , \end{aligned} \tag{1.5}$$

where $\dot{\phi}_0$ is the rate of change of the inflaton field. In a single-field inflationary scenario, one can also define the slow-roll parameters by requirements on the potential

$$\begin{aligned} \epsilon_V &= \frac{M_{\rm Pl}^2}{2}\left(\frac{V'}{V}\right)^2 \ll 1 \quad \text{and} \\ \eta_V &= M_{\rm Pl}^2 \frac{V''}{V} \ll 1 \ , \end{aligned} \tag{1.6}$$

which is the more commonly used definition. In the slow-roll limit the parameters (eq. 1.5) and (eq. 1.6) are related as

$$\epsilon_H = \epsilon_V \ , \qquad \eta_H = \eta_V - \epsilon_V \ . \tag{1.7}$$

The slow-roll parameters are thus related to the velocity and acceleration of the inflaton field along a field space trajectory. These are determined from an equation of motion

$$\ddot{\phi}_0 + 3H(t)\dot{\phi}_0 + V_\phi = 0 \ . \tag{1.8}$$

In the slow-roll limit, the potential is friction dominated, therefore the kinetic energy $\ddot{\phi}_0$ is very small compared to the other two terms. Together with the assumption that $|\dot{H}|/H \ll 1$, which follows from the assumption that $\dot{\phi}$ is small and thus the motion of the inflaton is potential energy dominated, the above equation reduces to

$$3H(t)\dot{\phi}_0 = -V_\phi \ , \tag{1.9}$$

which straightforwardly translates into the slow-roll parameters defined in (eq. 1.6) make sense. When $\ddot{\phi}_0$ cannot be neglected, the Hamilton-Jacobi parameters (eq. 1.5) are still well defined. This is the so-called fast-roll limit, which cannot be used to generate many e-folds of inflation. It can be used, however, to end inflation, by assuming that after some period of exponential expansion a change in the potential leads to a rapidly decreasing Hubble parameter. Then, via some reheating mechanism the kinetic energy of the inflaton must be transferred to other degrees of freedom to populate the universe with the particles we see today.

The above scenario assumes there is only one single light degree of freedom during inflation, that is not coupled to heavy degrees of freedom. Although a minimal assumption, it is certainly not a natural assumption. For example, it might be that there was more than one light field present during inflation, leading to models of multi-field inflation (eg. Groot Nibbelink and van Tent, 2000, 2002, Hwang and Noh, 2002, Wands et al., 2002, van Tent, 2004, Rigopoulos et al., 2006b, Seery and Lidsey, 2005, Rigopoulos et al., 2006a, 2007, Byrnes and Wands, 2006, Lalak et al., 2007b, Wands, 2008, Malik and Wands, 2009, Langlois et al., 2008c, Langlois and Renaux-Petel, 2008, Langlois et al., 2008a,b, Peterson and Tegmark, 2010). Another possibility is that the kinetic terms are noncanonical, with higher order derivative operators present (eg. Armendariz-Picon et al., 1999, Garriga and Mukhanov, 1999, Alishahiha et al., 2004, Bezrukov and Shaposhnikov, 2008, Barvinsky et al., 2008). Also, couplings to heavy degrees of freedom will lead to non-trivial results (chapters 4 and 5 and Tolley and Wyman, 2010, Chen and Wang, 2010b, Cremonini et al., 2010b. However, currently a large subset of the simplest models of single-field inflation remain perfectly compatible with current CMB precision measurements (Komatsu et al., 2010, Larson et al., 2010), predicting a nearly scale invariant power law inflation (Starobinsky, 1992, Adams et al., 2001, Tocchini-Valentini et al., 2005, Gong, 2005, Covi et al., 2006, Hunt and Sarkar, 2007, Ichiki et al., 2010, Peiris and Verde, 2010, Hamann et al., 2010)). Upcoming data, such as that from the Planck satellite promises to provide new handles on the overall shape of the spectrum and, particularly in combination with other data sets, could help us determine the precise nature of any possible features in it. If present, such features will lead to qualitative new tests on the single-field slow-roll paradigm (Kosowsky and Turner, 1995, Copeland et al., 1998) and could constitute strong evidence in favour of the existence of additional degrees of freedom present during the evolution of density perturbations as the universe inflated.

1.2.2 The power spectrum

Given an inflaton trajectory parametrised by a parameter ϕ_0, one can define perturbations around the trajectory as

$$\phi = \phi_0 + \delta\phi \; . \tag{1.10}$$

This definition is not gauge invariant, as also the spacetime metric will experience perturbations. Therefore, it is convenient for the calculation of the perturbations to introduce a gauge-invariant parameter (Sasaki, 1986, Mukhanov, 1988)

$$Q \equiv \delta\phi + \frac{\dot{\phi}}{H}\psi \; , \tag{1.11}$$

with ψ the scalar perturbation associated with the metric. Then, for Q the equation of motion is

$$\frac{d^2Q}{dt^2} + 3H\frac{dQ}{dt} - \frac{\nabla^2}{a^2}Q + m_{\text{eff}}^2 Q = 0 \; , \tag{1.12}$$

with $m_{\text{eff}}^2 = H^2(2 + 2\epsilon_H - 3\eta_H)$ for slow-roll inflation. Furthermore, a parameter $v = aQ$ can be introduced, which allows a convenient expression of (eq. 1.12) in conformal time, $dt = ad\tau$,

$$\frac{d^2v}{d\tau^2} - \nabla^2 v + a^2\left(H^2(2-\epsilon) + m_{\text{eff}}\right)v = 0 \; . \tag{1.13}$$

Considering the Fourier transformed equation,

$$\frac{d^2v}{d\tau^2} - k^2 v + a^2\left(H^2(2-\epsilon) + m_{\text{eff}}\right)v = 0 \; . \tag{1.14}$$

one sees there are two obvious limits. First, one has the short wavelength limit where $k \gg aH$. In this case (eq. 1.14) reduces to a normal oscillator equation. The other limit is when $k^2 \ll aH$, the late time limit, where the solution to (eq. 1.14) is $v \to a \times \text{const.}$. The behaviour changes when $k \sim aH$, when the wavelength k is of the same order as the Hubble radius, and thus crosses the horizon.

These perturbations are sourced by the everpresent quantum fluctuations at small scales, which are in a de Sitter universe described by a Bunch-Davies vacuum (Bunch and Davies, 1978). In order to match the classical perturbations of (eq. 1.13) these perturbations have to be quantised. One proceeds by Fourier transforming the perturbation equation, and expand the mode $v(\mathbf{k}, \tau)$ as

$$(2\pi)^3 v(\mathbf{k}, \tau) = v(k, \tau)\hat{a}(\mathbf{k}) + v^*(k, \tau)\hat{a}^\dagger(\mathbf{k}) \; , \tag{1.15}$$

with initial condition

$$v(k, \tau) = \frac{1}{\sqrt{2k}}e^{-ik\tau} \; , \tag{1.16}$$

since we need the growing mode. The mode that satisfies the initial condition is

$$v(k, \tau) = \frac{e^{-ik\tau}}{\sqrt{2k}}\frac{k\tau - i}{k\tau} \; , \tag{1.17}$$

which goes, well after horizon crossing, to

$$v(k,\tau) = \frac{i}{\sqrt{2k}}\frac{1}{k\tau} \ . \tag{1.18}$$

This derivation is for a single-field, the multi-field derivation is presented in section 4.3. The application to slow-roll two-field inflation is presented in section 5.4. In the remainder of this introductory section the single-field scenario will be discussed.

The initial curvature power spectrum is then given by the two-point correlation function of the curvature perturbations

$$P_R(\mathbf{k}) = \langle v(\mathbf{k},\tau_{\text{end}})v(-\mathbf{k},\tau_{\text{end}})\rangle \ . \tag{1.19}$$

where statistical isotropy can be assumed to perform the sum over all angles $P_R(k) = \sum_{\text{angles}} P_R(\mathbf{k})$. For single-field slow-roll inflation one finds for the power spectrum

$$P_R(k,\tau_{\text{end}}) = \left.\frac{H^2}{24\pi^2 M_{\text{Pl}}^4\epsilon}\right|_{k=aH} \tag{1.20}$$

where the power spectrum is evaluated at horizon exit, $k = aH$ and stays constant until $\tau = \tau_{\text{end}}$. For comparison with the cosmic microwave background, we need a spherical expansion (Lyth and Liddle, 2009)

$$v(\mathbf{k}) = \int_0^\infty dk \sum_{lm} v_{lm}(k,\tau_{\text{end}})Z_{klm}(\mathbf{x}) \tag{1.21}$$

with $v_{lm}(k,\tau_{\text{end}})$ the expansion coefficients of $v(\mathbf{k},\tau_{\text{end}})$ and

$$Z_{klm}(\mathbf{x}) \equiv \sqrt{\frac{2}{\pi}}k\, j_l(k|\mathbf{x}|)Y_{lm}(\theta,\phi) \ . \tag{1.22}$$

In this equation, j_l is the spherical Bessel function, (θ,ϕ) are the spherical directions of $\mathbf{x}$ and Y_{lm} is the spherical harmonic. Then, integrating v_{lm} over the sphere and, using statistical isotropy, summing over l and m, one obtains

$$\langle v^*_{lm}(k,\tau_{\text{end}})v_{l'm'}(k',\tau_{\text{end}})\rangle = (2\pi)^3 P_R(k)\delta(k-k')\delta_{ll'}\delta_{mm'} \ . \tag{1.23}$$

The scale dependence of $P_R(k)$ is defined as

$$n - 1 \equiv \frac{d\log P_R(k)}{d\log k} \ . \tag{1.24}$$

Assuming a constant spectral tilt, $n-1$, one finds that $P_R(k) \propto k^{n-1}$. From (eq. 1.20) we obtain, using $d\log(aH) \simeq Hdt$ and

$$-\frac{d(\log H)}{DN} \simeq -\epsilon\ , \qquad -\frac{d(\log \epsilon)}{DN} \simeq 4\epsilon - 2\eta\ , \tag{1.25}$$

where $dN = -Hdt$, that the spectral tilt is related to the slow-roll parameters as

$$n_s = 1 - 6\epsilon + 2\eta\ . \tag{1.26}$$

For multi-field inflation, also entropy perturbations, commonly known as isocurvature perturbations, are possible. These perturbations do not change the energy density but redistribute energy differently among particle species. They occur when perturbations in a direction normal to the inflaton trajectory are generated. Entropy perturbations are defined similarly as curvature perturbations, as a variance of a statistical field. Given an isocurvature perturbation $S(\mathbf{x},\tau) = (\delta n_i(\mathbf{x},\tau))/n_i$, evaluated at a slice of uniform energy density, we define the isocurvature power spectrum

$$P_S = \langle S_{\mathbf{x}} S_{-\mathbf{x}} \rangle \tag{1.27}$$

and similarly a cross correlation power spectrum

$$P_{RS} = \langle v_{\mathbf{x}} S_{-\mathbf{x}} \rangle\ . \tag{1.28}$$

Besides curvature and isocurvature perturbations, also tensor perturbations are possible. Sourced by perturbations of the metric, they have a different energy dependence

$$P_h(k) = \left.\frac{2H^2}{3\pi^2 M_{\rm Pl}^4}\right|_{k=aH}\ , \tag{1.29}$$

so that the tensor to scalar ratio r becomes (Liddle and Lyth, 1992)

$$r \equiv \frac{P_h(k)}{P_R(k)} = 16\epsilon\ . \tag{1.30}$$

This ratio can be used to determine the necessary field variation

$$r < 0.003\left(\frac{50}{N}\right)^2\left(\frac{\Delta\phi}{M_{\rm Pl}}\right)^2\ , \tag{1.31}$$

which is super-Planckian for large field inflationary models (Lyth, 1997).

1.2.3 Nongaussianities

One of the prime objectives of the Planck mission is to put better constraints on the existence of nongaussianities (Planck collaboration, 2006). Nongaussianities are given by the three point function (see Bartolo et al., 2004, for a review)

$$\langle v(\mathbf{k}_1, \tau_{\text{end}}) v(\mathbf{k}_2, \tau_{\text{end}}) v(\mathbf{k}_3, \tau_{\text{end}}) \rangle = \left(\frac{3}{5}\right)^3 (2\pi)^3 \delta(\mathbf{k}_1 + \mathbf{k}_2 + \mathbf{k}_3) F(k_1, k_2, k_3) \ . \quad (1.32)$$

In this equation, the δ-function ensures momentum conservation, meaning that the vectors k_1, k_2 and k_3 should form a triangle in phase space. The shape of the triangle is determined by the shape function $F(k_1, k_2, k_3) = f_{NL} S(k_1, k_2, k_3)$, which consists of an identity shape function $S(k_1, k_2, k_3)$ and a factor f_{NL} that determines the relative weight of this shape.

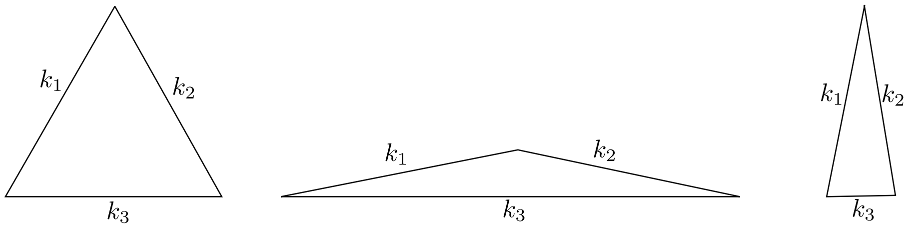


Figure 1.4: *A pictorial version of three commonly used shape functions for the nongaussianities*

Commonly, three momentum triangle shapes are studied (see figure 1.4). First, the local shape corresponds to $k_1 \approx k_2$, while $k_3 \rightarrow 0$. Second, the equilateral shape corresponds to $k_1 \approx k_2 \approx k_3$. Finally, the orthogonal shape is defined by the momenta $k_1 \approx k_2$, while $k_3 \rightarrow \infty$. These relations are used to define identity shape functions, that are multiplied with a factor f_{NL} as a measure of the amount of nongaussianities in the cosmic microwave background. Current best estimates, at 68% confidence level, are $f_{NL}^{(local)} = 32 \pm 21$, $f_{NL}^{(eq.)} = 26 \pm 140$ and $f_{NL}^{(ort.)} = 202 \pm 104$ (Komatsu et al., 2010).

For a free field theory $f_{\text{NL}} = 0$. For single-field inflation, the only source for nongaussianities is due to the nonlinear character of gravity, which generates local nongaussianities $f_{\text{NL}}^{(\text{local})}(5/12)(1 - n_s) \approx 10^{-2}$ (Maldacena, 2003, Acquaviva et al., 2003)

1.3 Heavy physics and inflation

Inflationary physics takes place at the highest energy scales — the natural scale for inflation is $O(10^{14})$GeV — where string theory becomes relevant. It is for this reason that many attempts have been made to find string theory models for inflation. A short overview is given in the next subsection. In this subsection, also the vacuum structure of string theory is discussed. Almost all 4 dimensional reductions of string theory bring many degrees of freedom, moduli, that need some mechanism to be stabilised. No matter the mechanism, these modes will be lighter than the Planck mass and thus be relatively light as seen from the point of view of inflation. This warrants a study of an effective field theory description of string inflation, which in turn requires a discussion of the role, or rather the removal, of these light fields. A short overview of this method is presented in the second subsection, a further study is the topic of this thesis.

1.3.1 String theory and inflation

Flux compactifications

String theory is conventionally and conveniently defined in ten dimensions. In order to get an effective four dimensional theory, six dimensions have to be compactified on a six dimensional manifold. Using the supergravity descriptions of string theory, it is known that compactifying on a torus leaves the maximal amount of supersymmetry, while compactifying on manifolds with less internal symmetry leads to less supersymmetric effective theories (Aspinwall, 2000, Freedman and van Proeyen, 2009). Compactifying IIB supergravity on a Calabi-Yau manifold leads to an $\mathcal{N} = 2$ supergravity. Adding orientifolds, branes and/or fluxes can reduce this to $\mathcal{N} = 1$ supergravity. For this reason, Calabi-Yau manifolds are widely used as compact manifold for compactification.

Calabi-Yau manifolds have many internal symmetries (Candelas and de la Ossa, 1991) that will show up as light degrees of freedom, moduli, in the effective description, unless some stabilising potential is generated. Work by Gukov et al. (2000), Dasgupta et al. (1999), Greene et al. (2000) and Giddings et al. (2002) has shown that three form fluxes can provide the necessary potential for the complex structure moduli and the dilaton, leaving only the volume moduli unfixed. A first method to fully stabilise all moduli, the so called KKLT framework, was provided by Kachru et al. (2003a), where nonperturbative effects provide a potential for the volume moduli. In the KKLT paper, also a method to provide an additional contribution to the vacuum expectation value is provided. This is necessary, as without it the model predicts a vacuum in anti-de Sitter space, while our universe has always been characterised by a positive vacuum expectation value.

Another method of stabilising all moduli was found by Balasubramanian et al. (2005), the so-called large volume compactifications. The method invoked is that a combination of stringy corrections (α' corrections) and nonperturbative effects generate a potential that is zero at the origin and at infinity, while it is negative in between. It can be shown that at an exponentially large volume there exists a minimum with a negative vacuum expectation value and broken supersymmetry. These models generally predict a low scale of supersymmetry breaking and can be phenomenologically successful (eg. Conlon et al., 2005, 2007).

Inflation in string theory

A good overview of the current state-of-the-art of string theory inflation is given in McAllister and Silverstein (2008), Baumann and McAllister (2009) and references therein. String theory models fall in two broad classes, small-field and large-field models, depending on the range of variation of the inflaton. Small-field models are models where the inflaton degree of freedom is represented by a parameter that can only move small distances on the internal compactification manifold, and by invoking the Lyth bound (eq. 1.31) thus operates at low energy scales. The typical model consists of a D-brane moving along a warped throat (Klebanov and Strassler, 2000). A first model, using the potential of a $D3-\overline{D3}$ brane pair in a KKLT compactification, was provided by Kachru et al. (2003b). Further work has shown that the true story is more complicated (eg. Baumann et al., 2008, 2007, Krause and Pajer, 2008) although the strongly coupled dynamics of branes can also be used for inflation (Silverstein and Tong, 2004, Alishahiha et al., 2004). An overview of other available degrees of freedom can be found in Binetruy and Gaillard (1986)

In general, the KKLT framework has lead to a large class of inflation models. Further addition of an extra nonperturbative term leads to the racetrack scenario (Blanco-Pillado et al., 2004, Lalak et al., 2007c, Blanco-Pillado et al., 2006b), Kähler moduli inflation (Conlon and Quevedo, 2006, Blanco-Pillado et al., 2010) or other volume moduli inflation models (Misra and Shukla, 2008, Conlon et al., 2008, Badziak and Olechowski, 2009, 2008). Models that use the volume modulus as inflaton also draw inspiration from the large volume compactification framework.

As discussed above, inflation using a string degree of freedom allows the inflaton to move only sub-Planckian distances in field space and thus allows only for a low scale of inflation (Kallosh and Linde, 2007b). A detection of tensor modes (eq. 1.30), pointing to a large scale of inflation (eq. 1.31) thus suggests a non-stringy inflation mechanism. A loophole in this argument is exploited by monodromy inflation (Silverstein and Westphal, 2008, McAllister et al., 2010), by D-branes and axions respectively. Later models include and Cicoli et al. (2009), Kaloper and Sorbo (2009) and Dong et al. (2010). Furthermore, renewed interest in shift symmetries (Kallosh and Linde, 2010, eq.) also suggests the possibility of large field models.

The complicated vacuum landscape has sparked many multi-field modes of stringy inflation (eg. Groot Nibbelink and van Tent, 2000, 2002, Hwang and Noh, 2002, Wands et al., 2002, van Tent, 2004, Rigopoulos et al., 2006b, Seery and Lidsey, 2005, Rigopoulos et al., 2006a, 2007, Byrnes and Wands, 2006, Lalak et al., 2007b, Wands, 2008, Malik and Wands, 2009, Langlois et al., 2008c, Langlois and Renaux-Petel, 2008, Langlois et al., 2008a,b, Peterson and Tegmark, 2010). Given the presence of multiple fields and interactions, one expects observable features in the power spectrum (chapter 5, Cremonini et al., 2010a,b) and nongaussianities, see chapter 5. There is a lot of work on this subject, see eg. Maldacena (2003), Bernardeau and Uzan (2002), Creminelli (2003), Bartolo et al. (2004), Rigopoulos et al. (2006a), Seery and Lidsey (2005), Langlois et al. (2008a), Langlois et al. (2008b), Langlois et al. (2008c), Tolley and Wyman (2010), Chen and Wang (2010b) and Barnaby (2010).

An important feature is that almost all models are set up such that the inflationary physics is happening in a sector well separated from all other physics and that the sectors are only coupled due to gravity. In principle, one could then integrate out these hidden degrees of freedom. However, integrating out physics is very challenging, therefore this separation is used as a justification to truncate the additional degrees of freedom. In this thesis, however, it is shown that such truncations are usually not well justified, and discusses the non-decoupling of these gravitationally coupled degrees of freedom. Truncation only makes sense when it is done consistently, such that the equation of motion obtained from the truncated theory is the same as the equation of motion obtained from the full theory. In this thesis, it is shown that consistently truncating degrees of freedom is far more subtle than usually assumed. A recent application of this knowledge in terms of the η-problem, the problem that the natural scale for η if $O(1)$ instead of $O(10^{-3})$ is given by Hardeman et al. (2010).

1.3.2 Supergravity and effective field theory

In this thesis, we will study the applicability of supergravity as an effective theory. For this to be possible, it is necessary to consistently truncate heavy degrees of freedom, as discussed in chapter 2. It is generally assumed that having gravitational strength couplings between two sectors is enough to truncate one of these sectors. Before discussing and properly defining this statement later in this section, let us first introduce the notation used.

Notation and conventions

Supergravity is defined by an action

$$S = M_{\rm Pl}^2 \int \mathrm{d}^4x \sqrt{g} \left[\frac{1}{2} R - g^{\mu\nu} G_{I\bar{J}} \nabla_\mu \xi^I \nabla_\nu \bar{\xi}^{\bar{J}} - V M_{\rm Pl}^2 \right] , \tag{1.33}$$

in which $G^{I\bar{J}}$ is the inverse field space metric $G_{I\bar{J}} = \partial_I \partial_{\bar{J}} G$ and $g_{\mu\nu}$ is the spacetime metric with associated Riemann scalar R. Greek indices run over spacetime coordinates $\{\mu, \nu\}$, capital indices run over all fields $\{I, \bar{J}\}$. For calculational convenience we have defined the scalar fields ξ and functions V, K and W to be dimensionless. The (F-term) potential V of the scalar sector is defined as

$$V = e^G \left(G_I G^I - 3 \right) . \tag{1.34}$$

Through the metric defined as above and $G_I = \partial_I G$ the action (eq. 1.33) is completely specified by the real Kähler function $G(\xi, \bar{\xi})$, which is, when $W \neq 0$, related to global supersymmetry quantities through

$$G(\xi, \bar{\xi}) = K(\xi, \bar{\xi}) + \log\left(W(\xi)\right) + \log\left(\bar{W}(\bar{\xi})\right) \tag{1.35}$$

in terms of the real Kähler potential $K(\xi, \bar{\xi})$ and the holomorphic (dimensionless) superpotential $W(\xi)$. The definition for G is convenient as it is invariant under Kähler transformations, i.e. it is invariant under the simultaneous transformation of

$$\begin{aligned} &K(\xi, \bar{\xi}) \to K(\xi, \bar{\xi}) + f(\xi) + \bar{f}(\bar{\xi}) \quad \text{and} \\ &W(\xi) \to e^{-f(\xi)} W(\xi) , \end{aligned} \tag{1.36}$$

for an arbitrary holomorphic function $f(\xi)$. Furthermore, the first and second derivatives of the potential (eq. 1.34) are easily calculated and read

$$\nabla_I V = e^G \left(G_I + G^J \nabla_I G_J \right) + V G_I , \tag{1.37}$$

$$\begin{aligned} \nabla_I \nabla_{\bar{J}} V = e^G \left(G_{I\bar{J}} + \nabla_J G_K \nabla_{\bar{J}} G^K - \mathcal{R}_{I\bar{J}K\bar{L}} G^K G^{\bar{L}} \right) + G_I V_{\bar{J}} + G_{\bar{J}} V_I \\ + (G_{I\bar{J}} - G_I G_{\bar{J}}) V , \end{aligned} \tag{1.38}$$

$$\nabla_I \nabla_J V = e^G \left(2 \nabla_I G_J + G^K \nabla_I \nabla_J G_K \right) + G_I V_J + G_J V_I + \left(\nabla_I G_J - G_I G_J \right) V . \tag{1.39}$$

In the action (eq. 1.33) we have denoted the gauge covariant derivatives by $\nabla_\mu \xi^I = \partial_\mu \xi^I - A^a_\mu k_a{}^I(\xi)$, and $k_a{}^I(\xi)$ are the Killing vectors that define the gauge transformations of the scalars,

$$\delta_{gauge} \xi^I = k_a{}^I(\xi) \alpha^a , \qquad a = 1, \ldots, n_v \ , \tag{1.40}$$

where α^a are the gauge parameters. The kinetic terms of the gauge fields are determined by the (holomorphic) gauge kinetic functions $f_{ab}(\xi)$

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} (\text{Re} f_{ab}) F^a{}_{\mu\nu} F^{b\mu\nu} + \frac{1}{4\sqrt{-g}} (\text{Im} f_{ab}) F^a{}_{\mu\nu} \epsilon^{\mu\nu\rho\sigma} F^b{}_{\rho\sigma} \ . \tag{1.41}$$

The scalar potential includes a contribution from F-terms and D-terms

$$V = V_F + V_D \ , \tag{1.42}$$

where V_F and V_D can be written in as a function of the auxiliary fields of the chiral and gauge superfields, F^I and D^a respectively,

$$V_F = G_{I\bar{J}} F^I F^{\bar{J}} - 3e^G = \mathrm{e}^G (G^{I\bar{J}} G_I G_{\bar{J}} - 3) \ , \tag{1.43}$$

$$V_D = \frac{1}{2} \text{Re}(f_{ab}) D^a D^b \ . \tag{1.44}$$

The auxiliary fields have equations of motion that can be solved algebraically in terms of the chiral fields and read

$$F^I = e^{G/2} G^{I\bar{J}} G_{\bar{J}} \ , \tag{1.45}$$

$$D^a = i (\text{Re} f)^{-1ab} k^I_b G_I = -i (\text{Re} f)^{-1ab} k^{\bar{I}}_b G_{\bar{I}} \ . \tag{1.46}$$

The two expressions given for the D-terms are equivalent due to the gauge invariance of the Kähler function $G(\xi, \bar{\xi})$ (Groot Nibbelink and van Holten, 2000, Binetruy et al., 2004)

$$\delta_{\text{gauge}} G = (k^I_a G_I + k^{\bar{I}}_a G_{\bar{I}}) \alpha^a = 0, \quad \text{for all} \quad a = 1, \ldots, n_v \ . \tag{1.47}$$

In this thesis we will assume that there are no constant Fayet-Iliopoulos terms (Fayet and Iliopoulos, 1974) present. Fayet-Iliopoulos terms require a more careful treatment that is outside the scope of this thesis.

The $\mathcal{N} = 1$ supersymmetry transformations of the fermions in the chiral and vector multiplets χ^I and λ^a are

$$\begin{aligned} \delta \chi^I{}_L &= \tfrac{1}{2} \gamma^\mu \nabla_\mu \xi^I \epsilon_R - \tfrac{1}{2} \, \mathrm{e}^{\frac{1}{2} K} K^{I\bar{J}} \mathcal{D}_{\bar{J}} \bar{W} \epsilon_L \ , \\ \delta \lambda^a &= \tfrac{1}{4} \gamma^{\mu\nu} F^a{}_{\mu\nu} \epsilon + \tfrac{1}{2} i D^a \gamma_5 \epsilon \ . \end{aligned} \tag{1.48}$$

Here ϵ is the parameter of the supersymmetry transformations, and γ^μ represent the gamma matrices as usual. The subscripts R and L of the fermions stand for right and left chirality respectively,

$$\chi^I{}_R = \tfrac{1}{2}(1-\gamma^5)\chi^I{}_R \qquad \chi^I{}_L = \tfrac{1}{2}(1+\gamma^5)\chi^I{}_L \tag{1.49}$$

From (eq. 1.48) we can see that in a homogeneous background ($\nabla^\mu \xi^I = F^a{}_{\mu\nu} = 0$), a set of necessary conditions for unbroken supersymmetry is

$$\mathcal{D}_I W = 0 \qquad \text{for all} \qquad I = 1, \ldots, n_c \ . \tag{1.50}$$

Equivalently this condition can be written in terms of the Kähler function as

$$\partial_I G(\xi, \bar{\xi}) = 0 \qquad \text{for all} \qquad I = 1, \ldots, n_c \ , \tag{1.51}$$

which, using (eq. 1.37) immediately shows that supersymmetric solutions are automatically extrema of the scalar potential (eq. 1.34). Furthermore, note that although it is always possible to break supersymmetry spontaneously by non-vanishing F-terms and zero D-terms (eq. 1.48), the relations (eq. 1.45) and (eq. 1.46) imply that non-vanishing D-terms necessarily require non-vanishing F-terms, and therefore supersymmetry can never be broken by D-terms alone (Choi et al., 2005).

The result (eq. 1.51) implies, together with the expression for the scalar potential (eq. 1.43) and (eq. 1.44), that supersymmetric critical points ξ_0^I with non vanishing superpotential $W(\xi_0) \neq 0$ always have a negative vacuum energy, i.e. they are Anti-de Sitter critical points

$$V(\xi_0) = -3e^{G(\xi_0)} < 0 \ . \tag{1.52}$$

Interestingly, supersymmetric critical points are always perturbatively stable, regardless of being local minima, maxima or saddle points. The reason is that in an Anti-de Sitter background a fluctuation with a tachyonic mass might still be stable as long as it satisfies the Breitenlohner-Freedman bound (Breitenlohner and Freedman, 1982)

$$m^2 \geq \frac{3}{4}\, V(\xi_0), \tag{1.53}$$

which is always fulfilled by supersymmetric critical points.

Gravitational couplings in supergravity and rigid supersymmetry

Having introduced the notation, we can now focus on defining gravitational couplings as discussed in this thesis. As is clear from (eq. 1.46) D-terms can never appear without F-terms. Therefore, in the following discussion we focus in the simpler case of only F-terms. The only effect D-terms can have is coupling decoupled sectors via gauge couplings, leading to an entirely different scenario that is not discussed in this thesis.

To describe a two-sector system we will consider a minimally coupled scenario (Cremmer et al., 1983a, Binetruy and Gaillard, 1985)

$$G(L,\bar{L},H,\bar{H}) = G^{(1)}(L,\bar{L}) + G^{(2)}(H,\bar{H}) \,, \tag{1.54}$$

with L, H denoting the fields in the two sectors respectively. In the following, we will take the indices $\{i, \bar{j}\}$ to run over the L fields, while $\{\alpha, \bar{\beta}\}$ denote the fields in the H sector. The L fields are assumed to be in the visible sector and thus allowed to be dynamical, while the H fields reside in another sector which is assumed not to take part in the dynamics and is hence called hidden sector. In chapter 2, a mass hierarchy between a light visible sector L and a heavy hidden sector H is present, which is why the sectors are labelled with L and H. However, currently L and H can be of arbitrary mass and are thus not necessarily light or heavy. This split of the Kähler function $G(L,\bar{L},H,\bar{H})$, (eq. 1.54), is invariant under Kähler transformations in each sector separately (chapter 2 and Choi et al., 2004, de Alwis, 2005a,b, Achúcarro and Sousa, 2008) and thus defines a sensible way of splitting up the action in multiple sectors. In terms of K and W, this definition has a conventional separation of the Kähler function

$$K(L,\bar{L},H,\bar{H}) + \log|W(L,H)|^2 = K^{(1)}(L,\bar{L}) + K^{(2)}(H,\bar{H}) + \log|W^{(1)}(L)W^{(2)}(H)|^2 \,, \tag{1.55}$$

but the superpotentials in each sector combine multiplicatively rather than add.

Let us illustrate the importance of this multiplicative superpotential in the situation in which the hidden sector resides in a supersymmetric vacuum, i.e. $\partial_a V(H_0) = 0$ and $\partial_a G^{(2)}(H_0) = 0$. We write the superpotential of the hidden sector as $W^{(2)}(H) = W_0^{(2)} + W_{\mathrm{global}}^{(2)}(H - H_0)$. The second term in this expression is what determines the potential for fluctuations around the minimum of the hidden sector, while the first constant term is just an overall contribution and hence not interesting for the internal hidden sector dynamics at energies much less than the Planck scale. However, for the gravitational dynamics and the remaining H^α sector this "vacuum energy contribution" $W_0^{(2)}$ is of crucial importance as it sets the scale of the potential (Davis and Postma, 2008, Hardeman et al., 2010)

$$V = e^{K^{(2)}}|W_0^{(2)}|^2 e^{G^{(1)}}\left(G_i^{(1)}G^{(1)^i} - 3\right) \,, \tag{1.56}$$

which is evaluated at $H = H_0$ such that all terms depending on $W^{(2)}_{\text{global}}$ vanish. The normal practise of setting $W^{(2)}_0$ to zero as an overall contribution to the hidden sector is neglecting the fact that gravity also feels the constant part of the potential energy, as opposed to field theory. The inflationary sector feels the presence of the hidden sector through this coupling and as such it may be more intuitive to regard $W^{(2)}_0$ to contain information about the inflationary sector rather than the hidden sector. Making a similar split in $W^{(1)}$, the constant part $W^{(1)}_0$ is the overall contribution to the hidden sector due to the inflaton sector.

The multiplicative superpotential also means that the zero-gravity limit to a global supersymmetry is more subtle than just taking $M_{\text{Pl}} \to \infty$ as is usually done. One must first determine a ground state which sets $W^{(1)}_0$ and $W^{(2)}_0$, and then send both $W^{(1)}_0 \to \infty$ and $W^{(2)}_0 \to \infty$ in such a way that the combinations $W^{(1)}_0 W^{(2)}_{\text{global}}$ and $W^{(2)}_0 W^{(1)}_{\text{global}}$ remain constant. The total superpotential

$$W = W^{(1)}_0 W^{(2)}_0 + W^{(1)}_0 W^{(2)}_{\text{global}} + W^{(2)}_0 W^{(1)}_{\text{global}} + W^{(1)}_{\text{global}} W^{(2)}_{\text{global}} \tag{1.57}$$

then consists of an overall infinite contribution, a finite sum of two terms and a negligible product. Only in this decoupling limit, does one recover the two independent global supersymmetry sectors with the naive additive behaviour in both the superpotential and the Kähler potential:

$$\begin{aligned} K(L, \bar{L}, H, \bar{H}) &= K^{(1)}(L, \bar{L}) + K^{(2)}(H, \bar{H}) \ , \\ W(L, H) &= W^{(1)}(L) + W^{(2)}(H) \ . \end{aligned} \tag{1.58}$$

However, one cannot use this split (eq. 1.58) and couple gravity back in (Davis and Postma, 2008). As explained, in supergravity the definition (eq. 1.58) is not invariant under Kähler transformations in each sector separately and is valid only in a specific Kähler frame or, say, gauge dependent (Achúcarro and Sousa, 2008). Another way to understand the result is to realise that the definition (eq. 1.58) does not lead to a Kähler metric and mass matrix that can be made block diagonal in the same basis (chapter 2), and thus there is no sense of "independent" sectors.

Insisting on the separate Kähler invariance of (eq. 1.54), the two-sector action (eq. 1.33) reads

$$S = M_{\text{Pl}}^2 \int \mathrm{d}^4x \sqrt{g} \left[\frac{1}{2} R - g^{\mu\nu} (G^{(1)}_{i\bar{j}} \partial_\mu L^i \partial_\nu \bar{L}^{\bar{j}} + G^{(2)}_{\alpha\bar{\beta}} \partial_\mu H^\alpha \partial_\nu \bar{H}^{\bar{\beta}}) - V M_{\text{Pl}}^2 \right], \tag{1.59}$$

with

$$V(L, \bar{L}, H, \bar{H}) = e^{G^{(1)} + G^{(2)}} \left(G^{(1)}_i G^{(1)i} + G^{(2)}_\alpha G^{(2)\alpha} - 3 \right) \ . \tag{1.60}$$

We will allow ourselves to drop the sector label from G in the remainder, since $G_L^{(1)} = G_L$ and similarly for H.

Non-decoupling in effective theories

After defining gravitational strength couplings we can continue the discussion on integrating out degrees of freedom in general field theories. As discussed in chapter 2, the requirement for truncating a degree of freedom is

$$\left.\frac{\delta\widehat{S}}{\delta L}\right|_{H_0} = \frac{\delta\widehat{S}|_{H_0}}{\delta L} = \frac{\delta S}{\delta L}\,, \tag{1.61}$$

where L and H point to a visible and hidden sector. Furthermore, $\widehat{S}$ refers to the full action for both the L and H sector, while S is the effective action for the L sector only.

Due to gravity, in any effective theory the gravitational force will couple everything to everything. This leads to Planck-suppressed corrections that are usually not relevant due to the huge scale difference between everyday physics and the Planck mass. Yet, in case of inflation there is a difference. On the one hand, the scale of the problem is much closer to the Planck scale. On the other hand, fields will move considerable distances in field space, also probing Planckian corrections. As shown in chapters 4 and 5 gravitational size couplings in non-linear sigma models can lead to corrections on the light degree of freedom. These corrections will manifest themselves as a reduced speed of sound for the light perturbations, leading to features in the power spectrum and nongaussianities.

In the context of supergravity, chapter 3 focuses on a model with a Kähler function of the form (eq. 1.54), where one of the sectors is in its supersymmetric minimum. This meets the requirement for consistent decoupling as derived in chapter 2. In fact, due to the separable Kähler function the field space manifold is actually a product manifold, making it possible to consistently decouple the heavy sector globally, thus consistently truncating the heavy physics (Groot Nibbelink and van Holten, 2000). However, instead of focusing on the supersymmetry broken sector, we focus on the supersymmetric sector, that *will* receive corrections from the supersymmetry broken sector. The reason is that the supersymmetry broken sector can provide an uplifting term, allowing de Sitter models in supergravity. Yet, we show that the coupling of the supersymmetry broken sector destabilises uplifted local minima after some amount of uplifting. In contrast, local maxima, that are actually stable in anti-de Sitter space due to the Breitenlohner-Freedman bound (Breitenlohner and Freedman, 1982), become stable when uplifted to Minkowski or de Sitter space.