Department of Physics and Astronomy
Heidelberg University

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Mathieu Kaltschmidt

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Gravity-Matter Systems
in
Asymptotically Safe Quantum Gravity

This Bachelor thesis has been carried out by

Mathieu Kaltschmidt

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Heidelberg University

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Prof. Dr. Jan M. Pawlowski

Gravity-Matter Systems in Asymptotically Safe Quantum Gravity

Mathieu Kaltschmidt

Abstract

In this work, we investigate asymptotically safe quantum gravity using non-perturbative Functional Renormalization Group methods. We compute the running of the Newton coupling and the cosmological constant and study the phase diagram of quantum gravity within the Einstein-Hilbert truncation in a background field approximation. In the pure-gravity sector, we find an UV-attractive interacting fixed point, providing further evidence for the Asymptotic Safety scenario for quantum gravity. Additionally, we include minimally coupled scalar, fermionic and gauge fields in our theory setting and analyze their impact on the fixed point structure. We find out, that scalar fields tend to destabilize the system, whereas fermionic fields and especially the gauge fields seem to be compatible with the theory setting. To conclude this work, we discuss some of the major problems of the background field approximation.

Zusammenfassung

In dieser Arbeit untersuchen wir asymptotisch sichere Quantengravitation mit Methoden der Funktionalen Renormierungsgruppe. Wir berechnen das Laufen der Newton-Konstante und der kosmologischen Konstante und analysieren das Phasendiagramm von Quantengravitation im Rahmen der Einstein-Hilbert Trunkierung in einer Hintergrundfeld-Näherung. Die Existenz eines ultraviolett attraktiven Fixpunktes im reinen Gravitationssektor wird nachgewiesen. Dies liefert bereits ein mögliches Indiz für die Realisierung einer asymptotisch sicheren Quantentheorie der Gravitation. Des Weiteren wird der Einfluss von minimal gekoppelten Skalar-, Fermion- und Eichfeldern auf die Fixpunktstruktur der Theorie untersucht. Wir finden heraus, dass Skalarfelder das System destabilisieren, wohingegen Fermionen und vor allem die Eichfelder mit der Theorie vereinbar zu sein scheinen. Am Ende der Arbeit werden einige der Probleme der Hintergrundfeld-Näherung erläutert.

Contents

Chapter 1.

Introduction

Einsteins theory of General Relativity successfully describes gravitational phenomena ranging from quotidian physics to the dynamics of whole galaxies very precisely in terms of the geometry of spacetime. This high precision has been proved once again recently in 2016, with the first observation of gravitational waves by the LIGO and VIRGO collaborations [1]. Nevertheless, General Relativity is a *classical* field theory and it is assumed, that the theory breaks down at some characteristic energy scale, i. e. the Planck scale

$$\Lambda_{\text{Planck}} \approx 10^{19}\ \text{GeV}.$$

For decades, physicists have been working on finding a microscopic, *quantum* theory of gravity, comparable to the successful description of the other three fundamental forces, namely the electromagnetic, the weak- and the strong nuclear forces, all unified in the Standard Model of Particle Physics. Such a quantum theory of gravity could provide interesting insides into the physics of e. g. the early universe or black holes. Among the most popular proposals for such a theory are e. g. String Theory or Loop Quantum Gravity.

One of the main problems in verifying predictions from such theories is, that currently it is all but impossible to probe (quantum) gravitational effects at energies near Λ_{Planck}. Particle colliders such as the Large Hadron Collider (LHC) nowadays reach maximum center-of-mass energies of about $\sqrt{s} \approx 14$ TeV $= 14 \cdot 10^3$ GeV. In addition to this experimental problem, it is well known, that the quantization of General Relativity leads to a (perturbatively) non-renormalizable theory due to the negative mass dimension of Newtons constant

$$[G] = -2$$

in $d = 4$ spacetime dimensions [10, 23]. During the last decades, the mathematical toolkit for theoretical physicists has evolved quite rapidly. Especially the development of the Functional Renormalization Group, in its modern formulation introduced by Kenneth Wilson in 1971 [27], offers a powerful, non-perturbative tool to solve path integrals in quantum field theory.

Proposed by Steven Weinberg in 1978 [24], the Asymptotic Safety scenario for Quantum Gravity provides a mechanism for constructing a fundamental quantum field theory of gravity in the language of the Functional Renormalization Group. It aims at generalizing the concept of Asymptotic Freedom, well-known e. g. in the context of Yang-Mills theories. The basic idea is, that the ultraviolet (= high energy) behavior of gravity may be governed by a Non-Gaussian Fixed Point (NGFP) of the underlying renormalization group flow. A first successful study of quantum gravity within the Asymptotic Safety

scenario was conducted by Martin Reuter in 1996 [19]. He derived the flow equations and proved the existence of such a NGFP in a pure-gravity setting in a truncated subsystem. The so-called Einstein-Hilbert truncation he employed back then, will also be the truncation of our choice for this work.
This thesis aims at investigating quantum gravity within the Asymptotic Safety scenario in the Einstein-Hilbert truncation after introducing the concepts needed for a general understanding of the subject. To get a first insight into the underlying structures, the theory is solved in a pure-gravity setting within a transverse-traceless spin-two graviton approximation. In order to probe a more realistic situation, the gravity-matter sector of the theory has to be taken into account. We include minimally coupled matter fields, i. e. scalar, fermionic and gauge fields and study their impact on the underlying fixed point structure of the Einstein-Hilbert truncation. Earlier studies put rather strict constraints on the matter content compatible with Asymptotic Safety, see e. g. [6]. In more recently published research papers, e. g. in [4, 14], strictly less severe constraints have been proposed. The latter results are based on calculations involving a vertex expansion of the effective average action. The flow of the couplings is then obtained from the flow of the n-point functions. Throughout this work, all computations will be performed in a background field approximation. Since this approximation has to be treated with care, we have to critically review our computations at the end of the thesis.

The structure of this work is the following: In chapter 2 the field theoretical language and the Functional Renormalization Group (FRG) are introduced. A derivation of Wetterich's exact renormalization group equation, a. k. a. the flow equation, completes our discussion of non-perturbative approaches to quantum field theory. Chapter 3 provides the background knowledge on gravity and curved spacetimes. In chapter 4, as a first step towards quantum gravity, the Asymptotic Safety approach is motivated and the flow equation is solved within the Einstein-Hilbert truncation in a transverse-traceless spin-two graviton approximation. The inclusion of matter is studied in chapter 5. In chapter 6 we critically review the background field approximation. To conclude this work, the results are summarized and discussed in chapter 7.
Throughout this thesis we use natural units such that $\hbar = c \equiv 1$. Einsteins sum convention is implicitly understood: Whenever an index appears twice in a single term, summation of that term over the whole index range is implied unless stated otherwise. As usual, greek indices refer to some d-dim. spacetime coordinates, ranging from 0 to $d-1$, i. e. $x^\mu = (x^0, x^1, \cdots, x^{d-1})$. Unless stated otherwise, we work in $d = 4$ spacetime dimensions.

Chapter 2.

Functional Methods in Quantum Field Theory

This chapter introduces a treatment of quantum field theory using functional methods. The main goal is to become familiar with the physical concepts and the notation used throughout this work and to derive the flow equation for the average effective action, introduced by Christof Wetterich in 1993 [26]. For the derivation of the flow equation we are following [9, 16].

2.1. Generating Functionals and Correlation Functions

Consider a theory setting of N real scalar fields $\varphi_a(x), a \in \{1, \dots, N\}$ in d-dimensional Euclidean space. The corresponding partition sum in presence of sources $J_a(x)$ reads

$$Z[J] = \frac{1}{\mathcal{N}} \int \mathcal{D}\varphi \, \mathrm{e}^{-\mathcal{S}[\varphi] + J\cdot\varphi} \,. \tag{2.1}$$

The action $\mathcal{S}$ is specified together with an ultraviolet cutoff scale Λ, later being the momentum scale where we initialize the flow equations and some normalization factor $\mathcal{N}$. In this notation, the scalar product sums over field components and integrates over all space,

$$J \cdot \varphi = \int_x J_a(x) \, \varphi_a(x) = \int_p \tilde{J}_a(p) \, \tilde{\varphi}_a(p), \tag{2.2}$$

with

$$\int_x = \int_{\mathbb{R}^d} \mathrm{d}^d x \qquad \text{and} \qquad \int_p = \int_{\mathbb{R}^d} \frac{\mathrm{d}^d p}{(2\pi)^d}. \tag{2.3}$$

The partition sum $Z[J]$ is called a *generating functional*. It directly allows us to compute field expectation values

$$\phi := \langle \varphi \rangle = \frac{1}{Z} \frac{\delta Z}{\delta J}\bigg|_{J=0} = \int \mathcal{D}\varphi \, \varphi \, \mathrm{e}^{-\mathcal{S}[\varphi] + J\cdot\varphi} \tag{2.4}$$

and higher order correlation functions

$$\langle \varphi(x_1) \cdots \varphi(x_n) \rangle := \langle \varphi^n \rangle = \frac{1}{Z} \frac{\delta^n Z}{\delta^n J}\bigg|_{J=0} = \int \mathcal{D}\varphi \, \overbrace{\varphi_1 \cdots \varphi_n}^{:= \varphi^n} \, \mathrm{e}^{-\mathcal{S}[\varphi] + J\cdot\varphi} \tag{2.5}$$

via functional differentiation. This means, we are basically able to compute all contributing Feynman diagrams for our theory setting, if we have knowledge of its corresponding (grand) canonical partition sum.

For a more efficient description of the theory in terms of only the *connected* correlation functions, we define the *Schwinger functional* $W[J]$ as the logarithm of $Z[J]$,

$$W[J] = \ln Z[J]. \tag{2.6}$$

It is the generating functional for the connected correlation functions. The normalization factor $\mathcal{N}$, introduced in (2.1) enters here as an additive constant, which drops out for all higher order correlation functions, except for the zero-point function. This term is connected to the thermodynamic quantities of the system and becomes important, when external parameters such as temperature, volume or the chemical potential are varied. For the case of quantum gravity, it is linked to the cosmological constant Λ. Nevertheless, in general we are only interested in correlation functions with $n \geq 1$ and therefore we drop this term.

Consider for example the connected two-point function $G_{ab}(x, y) = G_{\alpha\beta}$[1], known as the propagator, correlating the field φ_a at spacetime point x with the field φ_b at y,

$$\begin{aligned} G_{\alpha\beta} &= \frac{\delta^2 W[J]}{\delta J_\alpha \delta J_\beta} = \frac{\delta}{\delta J_\alpha}\left(\frac{1}{Z}\frac{\delta Z}{\delta J_\beta}\right) \\ &= \frac{1}{Z}\left(\frac{\delta^2 Z}{\delta J_\alpha \delta J_\beta}\right) - \frac{1}{Z^2}\left(\frac{\delta Z}{\delta J_\alpha}\right)\left(\frac{\delta Z}{\delta J_\beta}\right) \\ &= \langle \varphi_\alpha \varphi_\beta \rangle - \phi_\alpha \phi_\beta = \langle \varphi_\alpha \varphi_\beta \rangle_{\mathrm{c}}. \end{aligned} \tag{2.7}$$

The propagator is the key object in functional approaches to quantum field theory. It depends on the chosen background via J.

It is still possible to make our computations even more efficient, because $W[J]$ still contains some redundant information. Connected correlation functions can be separated into so-called one-particle irreducible (1PI) and one-particle reducible ones. The 1PI correlation functions are those, whose corresponding Feynman diagrams can *not* be separated into two disconnected ones by cutting a single internal line. As an example, contributing 1PI and reducible diagrams to the connected four-point function for Yukawa theory are depicted in figure (2.1).

The generating functional for the 1PI correlation functions, the *effective action* Γ, is obtained from the Schwinger functional via a Legendre transformation,

$$\Gamma[\phi] = \sup_J \left\{ \int_x J(x)\phi(x) - W[J] \right\} = \int_x J_{\sup}(x)\phi(x) - W\left[J_{\sup}\right], \tag{2.8}$$

1. To save on notation, we introduce collective indices $\alpha = (x, a)$ or (q, a) in momentum space.

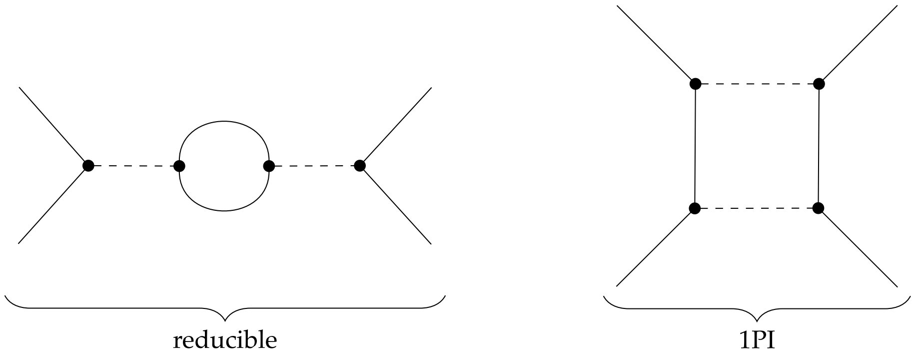

Figure 2.1. Contributing one-particle reducible and 1PI diagrams to the four-point-function in Yukawa theory, inspired by [8].

where J_{sup} has to be understood as a field-dependent current $J_{\mathrm{sup}}[\phi]$. In the following, we will drop the subscript, its meaning is implicitly understood. The quantum equation of motion derived from Γ reads

$$J(x) = \frac{\delta\Gamma[\phi]}{\delta\phi(x)}. \tag{2.9}$$

It allows us to understand the dynamics of field expectation values, taking the effects of all quantum fluctuations into account. From a physical point of view, the effective action Γ is the quantum analogue of the classical action $\mathcal{S}$. The performed Legendre transformation leads us to a mean field description of our theory with $\phi = \langle\varphi\rangle$ on a given background, as introduced before. The symmetries of the classical action are in general still present in the effective action.

In terms of the effective action, higher order correlation functions are again obtained by performing functional derivatives, but now w. r. t. the mean field ϕ,

$$\Gamma^{(n)}\left(x_1, \ldots, x_n\right) = \frac{\delta^n \Gamma}{\delta\phi\left(x_1\right)\cdots\delta\phi\left(x_n\right)}. \tag{2.10}$$

With the definition of the effective action (2.8), we find

$$\mathrm{e}^{-\Gamma[\phi]} = \int_\Lambda \mathcal{D}\varphi \exp\left(-\mathcal{S}[\phi+\varphi] + \int_x \frac{\delta\Gamma[\phi]}{\delta\phi(x)}\varphi(x)\right). \tag{2.11}$$

The solution of such functional integro-differential equations is highly non-trivial. To solve this problem, we want to make use of the Functional Renormalization Group. The general idea of this approach is to introduce a scale-dependent action Γ_k, interpolating between the bare, microscopic action $\mathcal{S}$ and the full quantum effective action Γ. A more formal motivation and a derivation of the equation governing this interpolation process is presented in the next section.

2.2. Functional Renormalization Group

The Functional Renormalization Group (FRG) is a mathematical tool, allowing us to investigate the dynamics of physical systems on different energy (momentum) scales. This idea is based on a continuous version of Leo P. Kadanoff's block spin model on the lattice [12] and was developed by Kenneth G. Wilson in 1971 [27]. It aims at solving the theory by integrating successively momentum shell by momentum shell, being the reason why the path integral approach to quantum field theory provides a suitable framework. The main advantage of the FRG approach is, that no regularization or renormalization procedure has to be applied. The latter one is already implemented systematically, which secures the self-consistency of the approach. As this section is only supposed to introduce the basics of the FRG, we refer the interested reader to more complete reviews, e. g. [9, 15], particularly for applications in different areas of physics.

As a first step towards a FRG equation we need to introduce an infrared cutoff scale k in our theory, below which the modes are not integrated out. A common way to introduce such a scale is by adding a scale-dependent cutoff term $\Delta\mathcal{S}_k$ in the definition of the partition sum (2.1) and therefore automatically also in the definition of the Schwinger functional (2.6):

$$W_k[J] = \ln Z_k[J] = \ln \int \mathcal{D}\varphi \, \mathrm{e}^{-\mathcal{S}[\varphi]+J\cdot\varphi-\Delta\mathcal{S}_k[\varphi]} \,. \tag{2.12}$$

The physical scale k we introduced here is known as *renormalization scale* and has units of inverse length, meaning large k correspond to small distances and vice versa. The cutoff term $\Delta\mathcal{S}_k$ is a quadratic functional depending on the field φ:

$$\Delta\mathcal{S}_k[\varphi] = \frac{1}{2}\varphi \cdot R_k \cdot \varphi = \frac{1}{2}\int_{x,y} \varphi_\alpha \, R_{k,\alpha\beta} \, \varphi_\beta. \tag{2.13}$$

The function R_k is called *regulator*. It plays an important role for this formulation of quantum field theory. The regulator is chosen such that only the propagation for momentum modes with $p^2 \lesssim k^2$ is suppressed. The most important physical limits are summarized in the following:

$$R_k(p^2) \to \begin{cases} k^2 & \text{for } p \to 0 \\ 0 & \text{for } p \to \infty \\ 0 & \text{for } k \to 0 \\ \infty & \text{for } k \to \Lambda \end{cases} \tag{2.14}$$

A convenient choice of the regulator is given by

$$R_k(p^2) = p^2 \cdot r_k(y), \tag{2.15}$$

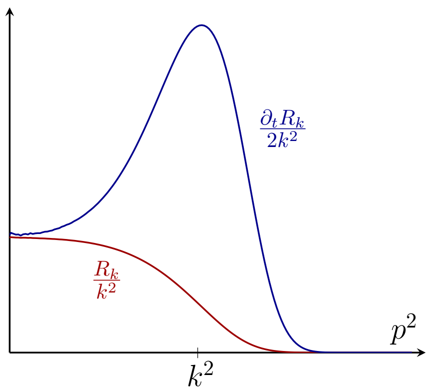

Figure 2.2. Shape of a typical exponential regulator function $R(p^2)$ and its derivative w. r. t. the RG time t. The regulator has a finite value for momenta smaller than k^2 and therefore acts as a suppressing mass term. The peak of $\partial_t R_k$ around $k^2 = p^2$ clearly shows the implementation of Wilsons idea of shell-wise momentum integration.

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with $y := \frac{p^2}{k^2}$, and a dimensionless regulator shape function r_k, only depending on the dimensionless momentum ratio y. There is a plethora of different types of shape functions, but for the computations performed in this work we restrict ourselves to a class of rather simple, so-called Litim-type regulators with shape functions

$$r_k(y) = \left(\frac{1}{y} - 1\right)\theta(1-y), \tag{2.16}$$

where θ is the Heaviside step function. This class of *sharp* regulators is a good choice for finding analytic FRG equations in simple approximations. For numerical approaches, exponential regulators such as depicted in figure (2.2) are well suited.

At this point it is quite convenient to introduce the *RG time t* as

$$t = \ln\left(\frac{k}{\Lambda}\right) \qquad \longrightarrow \qquad \partial_t = \frac{\partial}{\partial \ln(k/\Lambda)} = \frac{k}{\Lambda}\frac{\partial}{\partial(k/\Lambda)} = k\partial_k, \tag{2.17}$$

where Λ is a fixed reference scale. Usually one chooses the ultraviolet cutoff scale, where the flow is initialized.

In this setting, (2.12) provides a good starting point for solving the theory by successively lowering the cutoff scale k infinitesimally and integrating out all momentum modes $\varphi_{p\approx k}$. This procedure can be formalized by taking a scale derivative of our scale-dependent functional (2.12):

$$\begin{aligned}\partial_t W_k[J] &= -\frac{1}{2}\int \mathcal{D}\varphi\, \varphi(-p)\partial_t R_k(p)\varphi(p)\, \mathrm{e}^{-\mathcal{S}[\varphi]+J\cdot\varphi-\Delta\mathcal{S}_k[\varphi]} \\ &= -\frac{1}{2}\int_p \partial_t R_k(p) G_k(p) + \partial_t \Delta S_k[\phi],\end{aligned} \tag{2.18}$$

where we used the definition of the connected propagator:

$$G_k = \frac{\delta^2 W_k[\phi]}{\delta\phi(x)\delta\phi(y)}. \tag{2.19}$$

The *flowing* or *effective average action* Γ_k is then again defined via a modified Legendre transformation, including the insertion of ΔS_k:

$$\Gamma_k[\phi] = \sup_J \left(\int_x J(x)\phi(x) - W_k[J]\right) - \Delta S_k[\phi]. \tag{2.20}$$

This yields the modified, scale-dependent quantum equation of motion:

$$J(x) = \frac{\delta\Gamma_k[\phi]}{\delta\phi(x)} + (R_k\phi)(x). \tag{2.21}$$

Compared to the scale-independent version (2.9), we find an additional, regulator dependent term, but with the properties of the regulator presented in (2.14) in mind, we see that in the limit $k \to 0$ the initial equation of motion is restored. We find

$$\frac{\delta J(x)}{\delta\phi(y)} = \frac{\delta^2\Gamma_k[\phi]}{\delta\phi(x)\delta\phi(y)} + R_k(x,y). \tag{2.22}$$

With the help of these relations we are able to show that

$$\begin{aligned}\delta\left(x-x'\right) = \frac{\delta J(x)}{\delta J(x')} &= \int_y \frac{\delta J(x)}{\delta\phi(y)}\frac{\delta\phi(y)}{\delta J(x')} \\ &= \int_y \left(\Gamma_k^{(2)}[\phi] + R_k\right)(x,y)\, G_k\left(y-x'\right).\end{aligned} \tag{2.23}$$

Here, we used (2.22) and the definition of G_k (2.19). This yields the following important identity:

$$G_k = \left(\Gamma_k^{(2)} + R_k\right)^{-1}. \tag{2.24}$$

Altogether, we arrive at the *flow equation*, a. k. a. the *Wetterich equation* for the average effective action:

$$\begin{aligned}\partial_t\Gamma_k[\phi] &= -\partial_t W_k + \int (\partial_t J)\,\phi - \partial_t\Delta S_k[\phi] = -\partial_t W_k[J] - \partial_t\Delta S_k[\phi] \\ &\overset{(2.18)}{=} \frac{1}{2}\int_p G_k(p)\,\partial_t R_k(p) \\ &\overset{(2.24)}{=} \frac{1}{2}\,\mathrm{STr}\left[\left(\Gamma_k^{(2)}[\phi] + R_k\right)^{-1}\partial_t R_k\right].\end{aligned} \tag{2.25}$$

The supertrace STr sums over all internal indices and integrates over momentum space. For Grassmann fields, it also involves the inclusion of a minus sign. We will drop the S

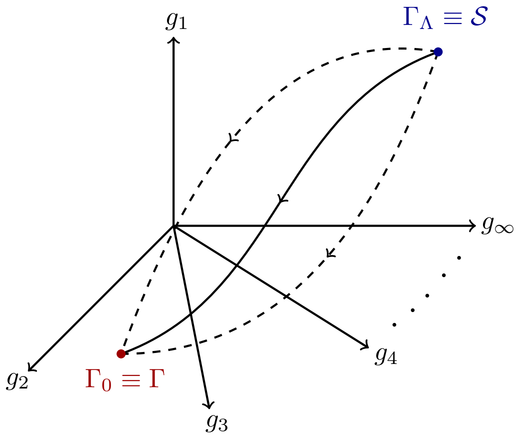

Figure 2.3. Flow of Γ_k through infinite-dimensional theory space for different regulators, inspired by [9]. Although the trajectories in theory space, governed by the flow equation (2.25) may be different, they flow towards the same quantum effective action $\Gamma_{k\to 0} \equiv \Gamma$.

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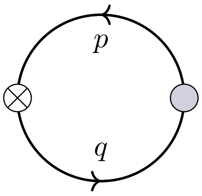
for the rest of this work, its meaning should be understood implicitly. The flow equation can be represented diagrammatically as a 1-loop equation:

$$\partial_t \Gamma_k[\phi] = \frac{1}{2} \sum_{i,j=1}^{N} \int_{p,q} \partial_t R_{k,ij}\,(p,q)$$



$$\left[\Gamma_k^{(2)}[\phi] + R_k\right]^{-1}_{ji} (q,p). \qquad (2.26)$$

The full propagator $\left[\Gamma_k^{(2)} + R_k\right]^{-1}$ is represented as usual as a single, double, dashed etc. line, dependent on the field content. The crossed circle $\otimes$ denotes the insertion of the respective regulator or more precisely its derivative w. r. t. the RG time t. Here $\partial_t R_{k,ij}(p,q) = \partial_t R_k(p^2)(2\pi)^d\, \delta_{ij}\, \delta(p-q)$ and therefore the trace on the r. h. s. effectively sums over just one index i and integrates over one loop momentum p.

It is important to mention, that the Wetterich equation is an *exact* equation, no approximations have been made. The only modification, the implementation of ΔS_k vanishes in the limit $k \to 0$. Solutions of the flow equation correspond to trajectories in *theory space*, the space spanned by all (= infinitely many) dimensionless couplings g_α. The choice of the regulator has direct impact on the exact form of the trajectory. This is often referred to as *scheme dependence*. Nevertheless, for all regulators satisfying the properties (2.14) it is guaranteed, that the flow will lead to the same quantum effective action Γ. For a visualization of this idea, have a look at figure (2.3). In principle this means, that $\lim_{k\to 0} \Gamma_k \equiv \Gamma$, but in most practical cases it is unavoidable to employ truncation schemes to be able to solve the flow equation. A plethora of different truncation schemes has been developed recently, details concerning the most important schemes can be found e. g. in the reviews about the FRG we referred to at the beginning of this section. We want to conclude this chapter with a more formal discussion of the concept of theory space, before proceeding to an introduction of the basic concepts of (classical) gravity.

2.3. Renormalization Group Flow and Theory Space

We want to use this section to formalize the concept of theory space we introduced in the last section and to discuss important characteristics of the renormalization group flow such as the β-functions and their zeros, the fixed points of the flow. For this part, we mainly follow [20].

The theory space is defined as the space spanned by all dimensionless couplings of the theory. To be more precise, it consists of all (action) functionals $A : \Phi \mapsto A[\Phi]$, that are compatible with the imposed symmetries of the theory such as e. g. diffeomorphism invariance in the case of (quantum) gravity.

The flow equation (2.25) defines a vector field $\vec{\beta}$ in theory space whose integral curves are the trajectories Γ_k parametrized by the scale k. Assuming the existence of a complete set of basis functionals $\{P_\alpha[\,\cdot\,]\}$, we can expand Γ_k as follows:

$$\Gamma_k[\Phi, \bar{\Phi}] = \sum_{\alpha=1}^{\infty} \bar{g}_\alpha(k) P_\alpha[\Phi, \bar{\Phi}]. \tag{2.27}$$

Here, the expansion coefficients $\bar{g}_\alpha(k)$ are given by the generalized couplings. Inserting this ansatz into the flow equation (2.25), yields a set of infinitely many coupled differential equations for the couplings:

$$k\partial_k \bar{g}_\alpha(k) = \bar{\beta}_\alpha\left(\bar{g}_1, \bar{g}_2, \cdots; k\right), \qquad \alpha = 1, 2, \cdots \tag{2.28}$$

The *beta functions* $\bar{\beta}_\alpha\left(\bar{g}_1, \bar{g}_2, \cdots; k\right)$ are the components of the vector field $\vec{\beta}$ and arise from an expansion of the trace on the r. h. s. of the flow equation in terms of the functional basis[2].

Up to this point, we are still dealing with dimensionful couplings $\bar{g}$, but as mentioned earlier usually the flow equation is expressed in terms of *dimensionless couplings*

$$g_\alpha \equiv k^{-d_\alpha} \bar{g}_\alpha, \tag{2.29}$$

where d_α is the canonical mass dimension of the respective coupling. The *essential* couplings[3] provide a set of coordinates for the theory space. This allows us to interpret the idea of renormalization theory in a new, geometrical way: We need to construct "infinitely long" trajectories Γ_k, that lie *entirely* in theory space. In this case, the couplings are prevented from diverging and we are able to define a consistent quantum field theory.

2. The expansion reads: $\frac{1}{2}\,\mathrm{Tr}\,[\cdots] = \sum_{\alpha=1}^{\infty} \bar{\beta}_\alpha\left(\bar{g}_1, \bar{g}_2, \cdots; k\right) P_\alpha[\Phi, \bar{\Phi}]$.
3. Essential in this sense means, that they can not be absorbed into the fields via a rescaling.

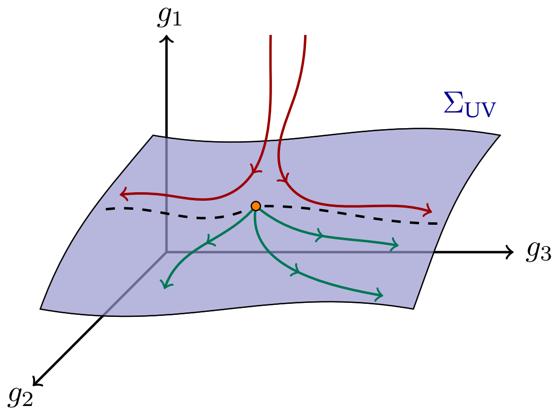


Figure 2.4. Visualization of a fixed point g^* (orange dot) with its corresponding UV hypersurface Σ_{UV} and trajectories starting at g^* (green) in theory space. The flow points towards the IR. Trajectories starting off the surface (red) are pulled towards the FP along the irrelevant direction (here: g_1) until the IR repulsive directions g_2 and g_3 dominate and drive the flow away from g^*. This figure is inspired by [7].

A *fixed point* g^* of the flow is a zero point of the vector field $\vec{\beta}$, i. e. $\beta_\alpha(g^*) \equiv 0\ \forall\alpha$. The existence of such fixed points is crucial for our discussion of Asymptotic Safety as an approach to quantum gravity, based on the concepts we introduced here.

In general, one distinguishes different classes of fixed points. The *Gaussian* or *non-interacting* fixed points (GFP) are classified by $g^*_\alpha = 0\ \forall\alpha$. This class of fixed points is relevant for perturbation theory, where the limit $k \to \infty$ is taken at such a GFP[4]. If at least one of the couplings $g^*_\alpha \neq 0$, the fixed point is classified as *Non-Gaussian* or *interacting* (NGFP). The idea of Asymptotic Safety relies on the existence of such a NGFP, rendering the theory "safe" from divergences in the ultraviolet (UV) regime. An important characteristic of a fixed point is its stability or more precisely if it is *attractive* or *repulsive* for near RG trajectories. Additionally one distinguishes between infrared ($k \to 0$) and ultraviolet ($k \to \infty$) attractive (repulsive) fixed points. To analyze this behavior, the flow near a fixed point is linearized, i. e.

$$\partial_t g_\alpha(k) = \sum_{j=1}^{\infty} B_{\alpha j}\left(g_j - g_j^*\right), \tag{2.30}$$

where we defined the *stability matrix* $\mathbf{B} = B_{\alpha j} = \partial_j \beta_\alpha(g^*_\alpha)$. The solution of the differential equation (2.30) reads:

$$g_\alpha(k) = g^*_\alpha + \sum_{j=1}^{\infty} C_j V^j_\alpha \left(\frac{k}{k_0}\right)^{\theta_j}. \tag{2.31}$$

4. E. g. in Yang-Mills theory, the concept of *Asymptotic freedom*, where the couplings tend to zero in the limit $k \to \infty$, is based on the existence of an UV attractive Gaussian fixed point, rendering the theory perturbatively renormalizable [11].

Here, the V^j are the eigenvectors of the stability matrix with eigenvalues θ_j a. k. a. *critical exponents*. In general, the θ_j, are complex numbers. We use the real part of the critical exponents to classify the coupling as *relevant* (= attractive) or *irrelevant* (= repulsive):

$$g_\alpha^* \text{ is } \begin{cases} \text{relevant} & \text{for } \mathfrak{Re}\left(\theta_j\right) > 0 \\ \text{irrelevant} & \text{for } \mathfrak{Re}\left(\theta_j\right) < 0 \end{cases} . \tag{2.32}$$

Fixed points with critical exponents $\theta_j = 0$ are called *marginal*. Based on this classification, it follows quite naturally to define an UV (or IR) *critical hypersurface* $\Sigma_{\rm UV}$ in theory space for a NGFP, consisting of all points that are pulled into the NGFP for increasing k. The dimension of $\Sigma_{\rm UV}$ is equal to the number of UV relevant couplings. This means, that trajectories lying on such a hypersurface tend to flow towards the fixed point in the UV limit. To visualize this idea, a schematic sketch of such a hypersurface in a 3-dim. theory space is depicted in figure (2.4).

Chapter 3.

Curved Spacetimes and Gravity

Our current understanding of gravity is manifested in Einsteins theory of General Relativity. In contrast to the treatment of the other fundamental forces, which are all described by gauge theories and summarized in the Standard Model of Particle Physics, gravity is based on the concept of curved spacetime. This chapter summarizes some of the general concepts and notions of General Relativity, needed for a basic understanding of the subject. For most of the concepts we present here, we are following Sean Carroll's lecture notes [3]. At the end of this chapter, we show why gravity can not be quantized in a perturbative manner, in opposition to the other three fundamental forces. For this part we follow [16].

3.1. An Introduction to Spacetime Geometry

When talking about the concept of curved spacetimes, one first needs a mathematical framework to quantify curvature and to understand how mathematical concepts such as differentiation and integration are generalized to curved spaces. The central objects in our discussion of curved spaces are *differentiable manifolds*, i e. topological spaces, that are locally diffeomorphic to $\mathbb{R}^n$. Locally in this sense means, that we can find coordinate maps $\phi_i : M \underset{\text{open}}{\supset} U_i \to \mathbb{R}^n$, such that the image $\phi_i(U_i)$ is open in $\mathbb{R}^n$, for every point on M, whereas globally the manifold may have a very complicated topology. A set of such coordinate maps $\{(U_\alpha, \phi_\alpha)\}$ that covers the entire manifold and where the charts are smoothly sewed together is called an *atlas*. For overlapping charts $U_\alpha \cap U_\beta \neq \emptyset$, the maps $(\phi_\alpha \circ \phi_\beta^{-1})$, a. k. a. coordinate transformations, must be smooth and differentiable. They are directly connected to the coordinates x^μ we'll work with later on.

Further, we need to introduce additional structures, such as vectors and tensors on manifolds, since they are the objects we are interested in when it comes to the discussion of physical models. To be able to talk about vectors, one needs to associate a *tangent space* T_p to every point p of the manifold. The tangent space is the set of all vectors at p and has the structure of a vector space with the same dimension as M. The disjoint union of all tangent spaces on M is called the *tangent bundle*. To specify the concept of the tangent space we claim, that it can be identified with the space of directional derivative operators along curves $\gamma : \mathbb{R} \to M$ through p. In this case, we find a basis of T_p as the set $\{\hat{\partial}_\mu\}$ of directional derivatives at p. It can be shown, that the directional derivatives can be decomposed into a sum of real numbers times partial derivatives, i. e. $\frac{d}{d\lambda} = \frac{dx^\mu}{d\lambda}\partial_\mu$, where λ is the parameter of the curve γ. This allows us to represent a vector $V = V^\mu \partial_\mu$ independently of the chosen coordinates. The basis vectors in some different coordinate system $x^{\mu'}$ are then simply related to the initial basis via $\partial_{\mu'} = \frac{\partial x^{\mu}}{\partial x^{\mu'}} \partial_{\mu}$ which yields the transformation law for vector components under general coordinate transformations,

$$V^{\mu'} = \frac{\partial x^{\mu'}}{\partial x^{\mu}} V^{\mu}. \tag{3.1}$$

Components obeying this transformation law are called *contravariant*. At this point it follows quite naturally to define the *cotangent space* T_p^* as the set of linear maps $\omega : T_p \to \mathbb{R}$. Elements of the cotangent space are called one-forms or dual vectors and similarly to the discussion of the tangent space, we find a suitable basis for T_p^* as the gradients $\{\mathrm{d}\hat{x}^{\mu}\}$, allowing us to represent arbitrary one-forms as $\omega = \omega_{\mu} \mathrm{d}x^{\mu}$. As before, we are interested in the transformation behavior of our basis one-forms, i. e. $\mathrm{d}x^{\mu'} = \frac{\partial x^{\mu'}}{\partial x^{\mu}} \mathrm{d}x^{\mu}$, and the dual vector components

$$\omega_{\mu'} = \frac{\partial x_{\mu}}{\partial x^{\mu'}} \omega_{\mu}. \tag{3.2}$$

This transformation behavior differs from the one found for vectors. We call components transforming as in equation (3.2) *covariant*. Now we are able to generalize these concepts by introducing tensors T of type (k, l) as

$$T = T^{\mu_1 \cdots \mu_k}{}_{\nu_1 \cdots \nu_l}\, \partial_{\mu_1} \otimes \cdots \otimes \partial_{\mu_k} \otimes \mathrm{d}x^{\nu_1} \otimes \cdots \otimes \mathrm{d}x^{\nu_l}. \tag{3.3}$$

Here $\otimes$ denotes the usual tensor product. The general transformation law for tensors follows naturally as expected from equations (3.1) and (3.2),

$$T^{\mu'_1 \cdots \mu'_k}{}_{\nu'_1 \cdots \nu'_l} = \frac{\partial x^{\mu'_1}}{\partial x^{\mu_1}} \cdots \frac{\partial x^{\mu'_k}}{\partial x^{\mu_k}} \frac{\partial x^{\nu_1}}{\partial x^{\nu'_1}} \cdots \frac{\partial x^{\nu_l}}{\partial x^{\nu'_l}} T^{\mu_1 \cdots \mu_k}{}_{\nu_1 \cdots \nu_l}. \tag{3.4}$$

Having understood the basic structures and their respective behavior under coordinate transformations, we are now able to present some of the most important tensors in General Relativity.

Maybe the most important object to quantify curved space is the *metric tensor* $g_{\mu\nu}$[5] and its inverse $g^{\mu\nu}$, related via $g^{\mu\nu} g_{\nu\sigma} = \delta^{\mu}{}_{\sigma}$. The metric and its inverse can be used to raise and lower indices, e. g. $x^{\mu} = g^{\mu\nu} x_{\nu}$. Additionally we can compute path lengths and proper time via the definition of the line element

$$\mathrm{d}s^2 = g_{\mu\nu} \mathrm{d}x^{\mu} \mathrm{d}x^{\nu}. \tag{3.5}$$

For arbitrary vector fields V and W the scalar product induced by the metric tensor reads

$$g(V, W) = g_{\mu\nu} V^{\mu} W^{\nu} = V^{\mu} W_{\mu} = g^{\mu\nu} V_{\mu} W_{\nu} = V_{\mu} W^{\mu}. \tag{3.6}$$

5. It is convenient to write the components $T^{\mu_1 \cdots \mu_k}{}_{\nu_1 \cdots \nu_l}$ when speaking about tensors T.

We will see, that the metric tensor already contains all the information on the geometrical structure of the respective manifold, whose curvature we want to quantify. Nevertheless, we first have to think about differentiation of general tensors again.
In flat space, the partial derivative is a map from (k, l) to $(k, l+1)$ tensor fields satisfying linearity and the Leibniz product rule. We want to generalize this concept to curved space by introducing the *covariant derivative* ∇[6]. In contrast to the usual partial derivative, the covariant derivative is independent of the chosen set of coordinates. Consider for example the covariant derivative of a vector field V, which can be written as a partial derivative plus some correction term due to its property to obey the Leibniz rule:

$$\nabla_\mu V^\nu = \partial_\mu V^\nu + \Gamma^\nu{}_{\mu\lambda} V^\lambda. \tag{3.7}$$

Here, the correction term is specified by the so-called *Christoffel symbols*, a. k. a. *connection coefficients*. They are determined by derivatives of the metric tensor:

$$\Gamma^\alpha{}_{\mu\nu} = \frac{1}{2} g^{\mu\lambda} \left(\partial_\mu g_{\nu\lambda} + \partial_\nu g_{\mu\lambda} - \partial_\lambda g_{\mu\nu} \right)^7. \tag{3.8}$$

It can be shown, that the connection coefficients themselves do *not* transform like tensor components, but are constructed in a way such that the combination (3.7) does. Note, that the covariant derivative reduces to the partial when applied to scalars. With this definition of the connection, we are now finally able to introduce the remaining tensor structures needed for the understanding of the calculations presented later on in this work.
The central object in our discussion of curvature is the *Riemann tensor* $R^\alpha{}_{\beta\gamma\delta}$. It is a $(1,3)$-tensor given by

$$R^\alpha{}_{\beta\gamma\delta} = \partial_\gamma \Gamma^\alpha{}_{\beta\delta} - \partial_\delta \Gamma^\alpha{}_{\beta\gamma} + \Gamma^\epsilon{}_{\beta\delta} \Gamma^\alpha{}_{\epsilon\gamma} - \Gamma^\epsilon{}_{\beta\gamma} \Gamma^\alpha{}_{\epsilon\delta}. \tag{3.9}$$

It contains all the information about the curvature of the respective manifold. Another useful definition of the Riemann tensor is related to the commutator of two covariant derivatives, acting on a vector field:

$$\left[\nabla_\mu, \nabla_\nu \right] A^\sigma = R^\sigma{}_{\rho\mu\nu} A^\rho. \tag{3.10}$$

We are also interested in contractions of the Riemann tensor, especially the *Ricci tensor*

$$R_{\mu\nu} = R^\alpha{}_{\mu\alpha\nu} = g_{\alpha\beta} R^\beta{}_{\mu\alpha\nu} \tag{3.11}$$

6. In the context of quantum field theory, the gauge covariant derivative is often written as D. Nevertheless, throughout this thesis we will use ∇ to indicate any kind of covariant derivative.
7. This holds only true, if the connection is *torsion free*, i. e. $T^\lambda{}_{\mu\nu} = \Gamma^\lambda{}_{\mu\nu} - \Gamma^\lambda{}_{\nu\mu} = 2\Gamma^\lambda{}_{[\mu\nu]} = 0$, and fullfils *metric compatibility*, i. e. $\nabla_\rho g_{\mu\nu} = 0$. For the most important connection in the context of General Relativity, the *Levi-Civita connection*, these properties are fullfilled. The fundamental theorem of Riemannian geometry states, that for every Riemannian manifold there exists a unique Levi-Civita connection. It is determined by the Koszul formula.

and the *Ricci scalar*

$$\mathcal{R} = g_{\mu\nu} R^{\mu\nu} = R^{\mu}{}_{\mu}. \tag{3.12}$$

At this point, we also want to introduce the *Einstein tensor*, defined as

$$G_{\mu\nu} = R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} \mathcal{R}. \tag{3.13}$$

Having introduced the setup for the calculations performed in this work, we are now ready to introduce the *Einstein-Hilbert action*, providing the starting point for an investigation of quantum gravity within the Functional Renormalization Group approach.

3.2. From Geometry to Einsteins Equations

The Einstein-Hilbert action, given by

$$\mathcal{S}_{\text{EH}} = \frac{1}{16\pi G} \int_x \sqrt{g}\, (\mathcal{R} - 2\Lambda), \tag{3.14}$$

where G is Newtons coupling and Λ is the cosmological constant, describes a minimally coupled theory of gravity, leading to a $^1\!/_r$ gravitational potential in the non-relativistic limit. Note, that compared to the usual spacetime measure a factor of $\sqrt{g} := \sqrt{-\det g_{\mu\nu}}$ is included to preserve diffeomorphism invariance.[8]
Varying the Einstein-Hilbert action w. r. t. the inverse metric $g^{\mu\nu}$ yields Einsteins equations in absence of matter:

$$G_{\mu\nu} + \Lambda g_{\mu\nu} = 0. \tag{3.15}$$

The non-vacuum Einstein equations are obtained the same way, after the inclusion of matter in this setting by adding a matter part to the Einstein-Hilbert action:

$$\mathcal{S} = \frac{1}{8\pi G} \mathcal{S}_{\text{EH}} + \mathcal{S}_{\text{matter}}. \tag{3.16}$$

With the definition of the Energy-Momentum tensor $T_{\mu\nu}$, given by

$$T_{\mu\nu} = \frac{-2}{\sqrt{g}} \frac{\delta \mathcal{S}_{\text{matter}}}{\delta g^{\mu\nu}}, \tag{3.17}$$

8. Diffeomorphism invariance, i. e. the freedom of choosing an appropriate coordinate system, is the central symmetry in the context of General Relativity, based on the assumption, that coordinates do not exist a priori in nature, but are rather a mathematical tool used to describe it, that should not change the fundamental laws of physics.

we arrive at

$$\frac{1}{8\pi G}\left[G_{\mu\nu}+\Lambda g_{\mu\nu}\right]=T_{\mu\nu}. \tag{3.18}$$

In this form, Einsteins equations perfectly embody the direct correlation between curvature (l. h. s.) and the dynamics of the matter content of the theory (r. h. s.).
At the end of this chapter we want to emphasize the problem of perturbative non-renormalizability in the context of finding a quantum field theoretical description of gravity.

3.3. Perturbative Non-Renormalizability of Gravity

Naively, one could try to quantize gravity via the path integral formalism with a generating functional, given by $\int_{g_{\mu\nu}} \mathrm{e}^{-\mathcal{S}_{\mathrm{EH}}}$, as usual. The main problem in this approach is the lack of positivity of $\mathcal{S}_{\mathrm{EH}}$ causing problems with unitarity of the theory. In quantum gravity one usually introduces a linear split of the *full* metric $g_{\mu\nu}$, to perform expansions about a given background $\bar{g}_{\mu\nu}$, comparable to classical perturbation theory, which is based on coupling or amplitude expansions around the free Gaussian theory. The linear split reads

$$g_{\mu\nu}=\bar{g}_{\mu\nu}+\sqrt{G}h_{\mu\nu}, \tag{3.19}$$

with the metric fluctuation $h_{\mu\nu}$ defined as $h_{\mu\nu} = 1/\sqrt{G}\left(g_{\mu\nu}-\bar{g}_{\mu\nu}\right)$. This allows us to write the path integral in terms of the fluctuation field as

$$Z\left[J^{\mu\nu};\bar{g}_{\mu\nu}\right] \propto \int_{h_{\mu\nu}} \mathrm{e}^{-S_{\mathrm{EH}}\left[\bar{g}_{\mu\nu}+\sqrt{G}h_{\mu\nu}\right]+\int_x \sqrt{\bar{g}}\, J^{\mu\nu}h_{\mu\nu}}\,. \tag{3.20}$$

Note, that the source term depends on the determinant of the background metric, otherwise the usual $J^{\mu\nu}$ derivatives would not generate the n-point functions of the fluctuation field $h_{\mu\nu}$. We will come back to this problem, which is often referred to as *background independence*, at the end of this thesis in chapter 6.
After a suitable tensor decomposition of the fluctuation field and a gauge fixing procedure à la Faddeev-Popov[9], we are left with the gauge fixed Einstein-Hilbert action

$$\mathcal{S}_{\mathrm{grav}}[\bar{g},\Phi]=\mathcal{S}_{\mathrm{EH}}[g]+\mathcal{S}_{\mathrm{gf}}[\bar{g},h]+\mathcal{S}_{\mathrm{gh}}[\bar{g},\Phi]. \tag{3.21}$$

9. The functional quantization of gauge theories requires a gauge fixing procedure due to redundancies in the path integral measure. The idea of Faddeev and Popov is to represent the gauge fixing condition, which is implemented in the functional integral, as an additional functional integral over a set of Grassmann fields c and $\bar{c}$, known as *Faddeev-Popov ghosts*. Even though they are anticommuting Grassmann fields, they transform as scalars under Lorentz transformations. They also violate spin statistics. Nevertheless, they can be treated as additional particles in the computation of Feynman diagrams. For a detailed discussion, see e. g. ch. 16 in [18] or sec. 5.2 in [16].

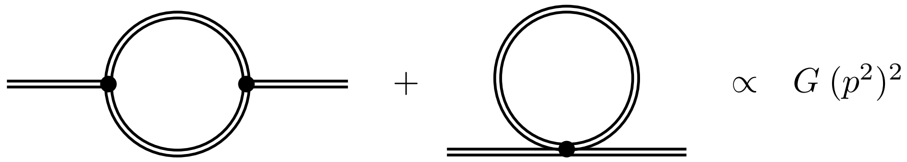


Figure 3.1. Vacuum polarization diagrams up to 1-loop order. The double lines represent the graviton propagator.

Here the pure gravity multi-field $\Phi = (h_{\mu\nu}, c_\mu, \bar{c}_\mu)$ was introduced. Altogether, this yields the gauge-fixed path integral representation of quantum gravity:

$$Z[J;\bar{g}] = \int_\Phi \mathrm{e}^{-\mathcal{S}_{\text{grav}}[\bar{g}_{\mu\nu},\Phi]+\int_x \sqrt{\bar{g}}\; J\cdot\Phi} \,. \tag{3.22}$$

An analysis of the canonical momentum dimensions of the essential couplings of this theory, G and Λ, results in:

$$[G] = \left[\mathrm{d}^d x\sqrt{g}\,\mathcal{R}\right] = 2-d, \qquad\qquad [\Lambda] = 2. \tag{3.23}$$

This implies, that the Newton coupling has a negative mass dimension in $d = 4$ spacetime dimensions. To investigate the consequences of this, one can consider the grade of divergence $\Lambda^{\delta(\gamma)}$ for a general graph γ with E external lines, I internal propagators and L loops. Here, Λ is an UV cutoff for the momentum integrals and $\delta(\gamma)$ is the index of the graph,

$$\delta(\gamma) = dL - 2\left(I - \sum_{n=3}^{\infty}\nu_n\right), \tag{3.24}$$

where the ν_n represent n-graviton vertices. After expressing the number of loops in terms of the internal lines and the n-graviton vertices and restricting ourselves to graphs satisfying $E + 2I = \sum_{n=3}^{\infty}\nu_n$, we find

$$\delta(\gamma) = d - \frac{d-2}{2}E + \sum_{n=3}^{\infty} v_n\delta\left(v_n\right), \tag{3.25}$$

where $\delta\left(v_n\right) = \frac{1}{2}(n-2)(d-2)$. After fixing the number of external lines, e. g. to $E = 2$, representing the case of vacuum polarization as depicted in figure (3.1), one is now able to investigate the grade of divergence for diagrams of different loop orders. t'Hooft and Veltman proved that the theory is renormalizable up to 1-loop order [23], but already at 2-loop order, Goroff and Sagnotti showed, that non-vanishing counterterms are generated [10]. In general, this is interpreted as the failure of perturbative quantization of gravity due to the negative mass dimension of the Newton coupling. This leads us to our discussion of Asymptotic Safety as a non-perturbative approach based on the FRG we presented in the last chapter.

Chapter 4.

Functional Renormalization and Quantum Gravity

After the formal introduction of the physical and mathematical concepts in the last chapters, we are now able to motivate and formulate the key idea of the Asymptotic Safety approach to quantum gravity, which aims at finding a quantum field theoretical description of gravity within the language of the Functional Renormalization Group. This chapter briefly discusses the requirements for the existence of such a theory. We proceed by solving the flow equation for quantum gravity within the Einstein-Hilbert truncation in a transverse-traceless spin-two graviton approximation and investigate its fixed point structure.

4.1. Asymptotic Safety

Asymptotic Safety generalizes the concept of Asymptotic Freedom known as a property of certain gauge theories. The latter one is based on the idea, that the interaction of particles becomes asymptotically weak in the high-energy limit. In the language of the FRG this means, that the UV behavior of such theories is governed by a non-interacting, Gaussian fixed point, allowing perturbative calculations. Maybe the most popular example for an asymptotically free theory is Quantum Chromodynamics (QCD), the theory of the strong interaction. In 1978, Steven Weinberg proposed the Asymptotic Safety scenario for gravity, based on the existence of an interacting, Non-Gaussian fixed point, rendering the theory non-perturbatively renormalizable in the UV [24]. The main advantage compared to other approaches is, that there is no need to introduce new symmetries, extra dimensions or other additional complications. The key requirements for an asymptotically safe quantum field theory of gravity can be summarized as follows:

1. The existence of an UV-attractive NGFP of the Renormalization Group flow of Γ_k has to be guaranteed.
2. For the predictivity of the theory, it is crucial to be able to fix the trajectory Γ_k by a finite amount of measurements, i. e. the corresponding UV hypersurface Σ_{UV} of the NGFP should be of finite dimension:

$$\dim \Sigma_{\mathrm{UV}} < \infty.$$

We want to probe these requirements in the following by computing the running of the Newton constant G_k and the cosmological constant Λ_k in a pure-gravity setting.

For a more detailed discussion of current Asymptotic Safety research, we refer the interested reader to [7]. Recently, very detailed textbooks covering both, the physical and the mathematical concepts of Asymptotic Safety, have been released. For a complete treatment of the subject it is worth to have a look at [17] or [21].

4.2. Einstein-Hilbert Truncation

Solving the flow equation analytically is nothing but impossible. Therefore it is unavoidable to truncate the initially infinite dimensional theory space to a finite subspace, to be able to find approximated solutions. It is important, that all terms, that are invariant under the imposed symmetry, i. e. invariant under diffeomorphism transformations, need to be taken into account. The easiest truncation fulfilling this requirement is the *Einstein-Hilbert truncation*. In 1996, Martin Reuter was the first to successfully investigate quantum gravity within the Einstein-Hilbert truncation [19]. He was able to prove the existence of a NGFP for a pure-gravity setting. This truncation takes only the scalar curvature $\mathcal{R}$ and the cosmological constant Λ into account and therefore the truncated subspace is 2-dimensional[10].

The full Einstein-Hilbert truncation reads:

$$\Gamma_k = 2\kappa^2 Z_{h,k} \int_x \sqrt{g}\,[-\mathcal{R} + 2\Lambda_k] + \mathcal{S}_{\text{gf}} + \mathcal{S}_{\text{gh}}, \tag{4.1}$$

where we abbreviated $\kappa^2 = (32\pi G)^{-1}$. Since we don't need to consider the gauge fixing action $\mathcal{S}_{\text{gf}}$ and the Non-Abelian ghost action $\mathcal{S}_{\text{gh}}$ for the calculation performed in this chapter, we won't specify their explicit forms at this point. We will come back to this discussion at the beginning of chapter 5.

The index k indicates the *running* i. e. the scale dependence of the cosmological constant and the wave function renormalization $Z_{h,k}$. It is quite convenient to define the running Newton coupling as

$$G_k = G \cdot \left(Z_{h,k}\right)^{-1}. \tag{4.2}$$

Up to this point, we did not care about the dimensionality of the two couplings. As already mentioned, one usually works with dimensionless couplings when solving the flow equation.

We therefore define the dimensionless renormalized cosmological constant

$$\lambda_k = \Lambda_k \cdot k^{-2} \tag{4.3}$$

and Newton constant

$$g_k = G_k \cdot k^{d-2} \overset{(d=4)}{=} G_k \cdot k^2. \tag{4.4}$$

10. Recently, more sophisticated truncations including higher-order curvature terms ($\mathcal{R}^2$, $f(\mathcal{R})$, $R^{\mu\nu}R_{\mu\nu}$ ...) have been investigated, see e. g. [2].

After these redefinitions, we can already compute the l. h. s. of the flow equation, i. e. the scale derivative w. r. t. the RG time t. The derivative only acts on λ_k and g_k.

$$\partial_t \Gamma_k = 2\kappa^2 Z_{k,h} \int_x \sqrt{g} \left\{ \eta_h \mathcal{R} + 2 \left(k^2 (\partial_t \lambda_k) + \Lambda_k (2 - \eta_h) \right) \right\}. \tag{4.5}$$

Here we introduced the *anomalous dimension* η_h, defined as

$$\eta_h = -\partial_t \ln Z_{h,k} = -\frac{\partial_t Z_{h,k}}{Z_{h,k}}. \tag{4.6}$$

We find terms of order $\sim \sqrt{g}$ and order $\sim \sqrt{g}\mathcal{R}$. The r. h. s. of the flow equation is assumed to admit an expansion in terms of invariants, which will also be of order $\sim \sqrt{g}$ and $\sim \sqrt{g}\mathcal{R}$ since we are working in the Einstein-Hilbert truncation[11]. This will allow us to determine the β-function for λ_k and the explicit form of η_h by a comparison of the terms of these orders on both sides of the flow equation. The β-function for g_k follows directly from the anomalous dimension:

$$\beta_g = \partial_t g_k = (2 + \eta_h)\, g_k. \tag{4.7}$$

At this point, we want to use the idea of the background field approximation. We assume a linear split of the metric into a background field $\bar{g}$ and a fluctuation field h, as described in equation (3.19). The approximation at this point is, that in the following, we set all fluctuations to zero, i. e. we evaluate all derivatives etc. at $g = \bar{g}$. We will critically review this approximation in chapter 6. In general, there is no conceptual need to fix the background metric, it nevertheless is really useful to choose specific classes of backgrounds for certain computations, since this can simplify calculations a lot. We exploit the freedom of choosing a spherical background. On the four-sphere $\mathbb{S}^4$, which is a maximally symmetric space[12], the Riemann tensor and the Ricci tensor can be written as multiples of the curvature scalar:

$$\begin{aligned} \bar{R}_{\mu\nu} &= \frac{1}{4}\, \bar{g}_{\mu\nu} \bar{\mathcal{R}} \\ \bar{R}_{\mu\nu\rho\sigma} &= \frac{1}{12}\, (\bar{g}_{\mu\rho}\bar{g}_{\nu\sigma} - \bar{g}_{\mu\sigma}\bar{g}_{\nu\rho}) \bar{\mathcal{R}}. \end{aligned} \tag{4.8}$$

The bar in this notation refers to the background $\bar{g}$. The inversion of the two-point function on general curved backgrounds is non-trivial. We need to find a suitable tensor basis for the decomposition of the metric fluctuation $h_{\mu\nu}$. We choose a *York decomposition* and find:

$$h_{\mu v} = h^{\mathrm{TT}}_{\mu v} + \bar{\nabla}_\mu \xi_v + \bar{\nabla}_v \xi_\mu + \left(\bar{\nabla}_\mu \bar{\nabla}_v - \frac{1}{d} \bar{g}_{\mu v} \bar{\Delta} \right) \sigma + \frac{1}{d} \bar{g}_{\mu v} h. \tag{4.9}$$

11. There will also appear terms of higher order in curvature, but we will drop them.
12. Maximally symmetric spaces are characterized by the fact, that they have the same number of symmetries as ordinary Euclidean space.

Here, $h^{\mathrm{TT}}_{\mu\nu}$ is a transverse-traceless, spin-two degree of freedom, ξ_μ is transverse and carries a spin-one d. o. f. and σ and h have spin zero. For a more detailed motivation of the York decomposition, we refer to appendix A.

As a further approximation, we only take the contribution from the spin-two graviton mode $h^{\mathrm{TT}}_{\mu\nu}$ into account[13]. The other modes are included later on in chapter 5, when we couple matter to the system.

Our starting point for the computation of the r. h. s. of the flow equation is the transverse-traceless graviton two-point function:

$$\Gamma^{(2)}_{k,h^{\mathrm{TT}}h^{\mathrm{TT}}} = \frac{Z_k}{32\pi}\left(\bar{\Delta} - 2\Lambda_k + \frac{2}{3}\bar{\mathcal{R}}\right). \tag{4.10}$$

Here, we used the definition of the Laplace operator, given by $\Delta = -\nabla^2$. We use a regulator of the form

$$R_k = \Gamma^{(2)}_{h^{\mathrm{TT}}h^{\mathrm{TT}}}\Big|_{\Lambda_k=\bar{\mathcal{R}}=0} \cdot r_k\left(\frac{\bar{\Delta}}{k^2}\right), \tag{4.11}$$

with a Litim-type shape function (2.16). This allows us to determine the full propagator

$$G_{k,h^{\mathrm{TT}}h^{\mathrm{TT}}} := \left(\Gamma^{(2)}_{k,h^{\mathrm{TT}}h^{\mathrm{TT}}} + R_k\right)^{-1} = \frac{32\pi}{Z_{h,k}}\left(\bar{\Delta}\left(1+r_k\right) - 2\Lambda_k + \frac{2}{3}\bar{\mathcal{R}}\right)^{-1} \tag{4.12}$$

and the scale derivative of the regulator

$$\partial_t R_k = \frac{Z_{h,k}}{32\pi}\bar{\Delta}\left(\partial_t r_k - \eta_h r_k\right). \tag{4.13}$$

The computation of the r. h. s. of the flow equation is in general very hard because it involves the computation of a functional trace of a function depending on the Laplacian on a curved background. We can use *heat-kernel techniques* to solve such equations. Heat-kernel computations are based on a curvature expansion in powers of the curvature scalar $\mathcal{R}$. For more details, have a look at the second part of appendix A.

The previously obtained results lead us to

$$\frac{1}{2}\,\mathrm{Tr}\left[G_{k,h^{\mathrm{TT}}h^{\mathrm{TT}}}\,\partial_t R_k\right] = \frac{1}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k - \eta_h r_k\right)}{\bar{\Delta}(1+r_k) - 2\Lambda_k + \frac{2}{3}\bar{\mathcal{R}}}\right]. \tag{4.14}$$

We proceed by expanding this expression around vanishing curvature:

$$\tfrac{1}{2}\,\mathrm{Tr}\left[G_{k,h^{\mathrm{TT}}h^{\mathrm{TT}}}\,\partial_t R_k\right] = \tfrac{1}{2}\,\mathrm{Tr}\left[\tfrac{\bar{\Delta}(\partial_t r_k - \eta_h r_k)}{\bar{\Delta}(1+r_k)-2\Lambda_k}\right] - \tfrac{1}{3}\bar{\mathcal{R}}\,\mathrm{Tr}\left[\tfrac{\bar{\Delta}(\partial_t r_k - \eta_h r_k)}{\left(\bar{\Delta}(1+r_k)-2\Lambda_k\right)^2}\right] + \mathcal{O}(\mathcal{R}^2). \tag{4.15}$$

13. This choice is motivated by the fact, that the spin-two mode carries the most degrees of freedom. From the 10 initial d. o. f. in this decomposition, 4 can be removed via gauge fixing, the remaining 6 are divided into 5 from the $h^{\mathrm{TT}}_{\mu\nu}$ mode and 1 for the trace mode h.

We evaluate these two terms separately using the heat-kernel formulas presented in appendix A. The result for the first term reads

$$\begin{aligned}\frac{1}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k-\eta_h r_k\right)}{\bar{\Delta}(1+r_k)-2\Lambda_k}\right]&=\frac{1}{2}\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\left[5\Phi_2^1(-2\Lambda_k)-\frac{5}{6}\bar{\mathcal{R}}\Phi_1^1(-2\Lambda_k)\right]\\&=\frac{1}{2}\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\left[\frac{1}{1-2\lambda_k}\left(5\left(1-\frac{\eta_h}{6}\right)-\frac{5}{3}\bar{\mathcal{R}}\left(1-\frac{\eta_h}{4}\right)\right)\right],\end{aligned} \tag{4.16}$$

with the definition of threshold functions

$$\begin{aligned}\Phi_n^p(\omega)&=\frac{1}{\Gamma(n)}\int_0^\infty \mathrm{d}z\; z^{n-1}\frac{z(-2zr_k(z)-\eta_\Psi r_k(z))}{(z(1+r_k(z))+\omega)^p}\\&=\frac{1}{\Gamma(n)}\frac{1}{(1+\omega)^p}\left(\frac{2}{n}-\frac{\eta_\Psi}{n(n+1)}\right).\end{aligned} \tag{4.17}$$

In the last step, we evaluated the threshold functions for the Litim-type shape function. We used η_Ψ, since we want to keep this formula as general as possible and we will use it multiple times for different fields throughout this thesis. Analogously, the second term in our expansion reads

$$\begin{aligned}-\frac{1}{3}\bar{\mathcal{R}}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k-\eta_h r_k\right)}{\left(\bar{\Delta}(1+r_k)-2\Lambda_k\right)^2}\right]&=-\frac{1}{2}\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\;\frac{10}{3}\bar{\mathcal{R}}\;\Phi_2^2(-2\Lambda_k)\\&=-\frac{1}{2}\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\left[\frac{1}{(1-2\lambda_k)^2}\left(\frac{10}{3}\bar{\mathcal{R}}\left(1-\frac{\eta_h}{6}\right)\right)\right].\end{aligned} \tag{4.18}$$

With these results, we are able to determine the β-function for the cosmological constant, simply by comparing the order $\sim\sqrt{\bar{g}}$ terms occurring on the l. h. s. of the flow equation and the results from the heat-kernel expansion of the functional trace:

$$\beta_\lambda=\partial_t\lambda_k=-4\lambda_k+\frac{\lambda_k}{g_k}\partial_t g_k+\frac{5}{4\pi}g_k\frac{1-\frac{\eta_h}{6}}{1-2\lambda_k}. \tag{4.19}$$

Additionally, we find an expression for the anomalous dimension η_h by comparing the terms of order $\sim\sqrt{\bar{g}}\bar{\mathcal{R}}$:

$$\eta_h=-\frac{5g_k}{3\pi}\left(\frac{1-\frac{\eta_h}{4}}{1-2\lambda_k}+2\frac{1-\frac{\eta_h}{6}}{(1-2\lambda_k)^2}\right). \tag{4.20}$$

To find the fixed points for this truncated solution, we use `Mathematica` to solve the equation

$$\vec{\beta}=\begin{pmatrix}\beta_g\\ \beta_\lambda\end{pmatrix}\overset{!}{=}\begin{pmatrix}0\\0\end{pmatrix}. \tag{4.21}$$

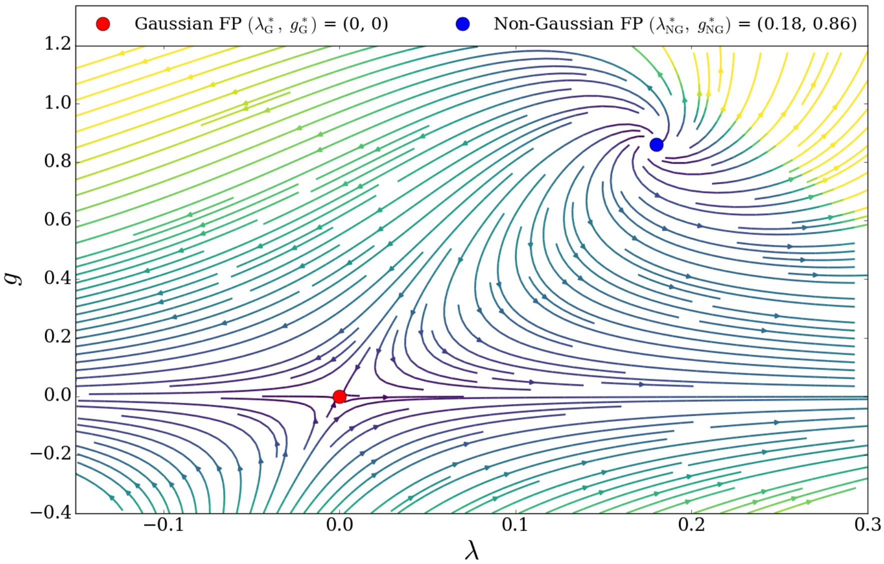


Figure 4.1. Flow diagram for the Einstein-Hilbert truncation in h^{TT} approximation as computed in this work. The flow points towards the infrared. The postulated UV-attractive Non-Gaussian fixed point (blue dot) is clearly visible.

We arrive at the following values for the Newton coupling and the cosmological constant at the NGFP:

$$(g_k^*, \lambda_k^*) = (0.86, 0.18). \tag{4.22}$$

The critical exponents of the fixed point are given by the complex conjugated pair

$$\theta_{1,2} = 2.9 \pm 2.6i. \tag{4.23}$$

The corresponding flow diagram is depicted in figure (4.1). The "swirl" around the NGFP is explained by the fact, that $\mathfrak{Im}\ \theta_{1,2} \neq 0$. The visualization and the plotting of the diagram was done in `Python`.

Based on the results obtained in the scope of this calculation, we can confirm the existence of a UV-attractive Non-Gaussian fixed point for this truncated pure-gravity system, providing a good starting point for further analysis. As a first extension of this truncation, we want to go beyond the h^{TT} approximation by taking the contributions for the other modes into account. Then, we will investigate the impact of minimally coupled matter fields on the system. All these calculations will be explained in full detail in the next chapter.

Chapter 5.

Asymptotic Safety of Gravity-Matter Systems

The calculation in h^{TT} approximation we performed in the last chapter, already allowed us to investigate the characteristic fixed point structure of the Einstein-Hilbert truncation. Nevertheless, in this part of the thesis, where the impact of minimally coupled matter fields is analyzed, we want to work with the full result, including also the vector and scalar modes arising from the York decomposition (4.9) of the fluctuation field.

5.1. Inclusion of the other Graviton Modes

The gauge conditions and the results for the full graviton propagator are taken from [16]. For the full solution, we have to take care of an additional gauge fixing term

$$\mathcal{S}_{\mathrm{gf}} = \frac{1}{2\alpha}\int_x \sqrt{\bar{g}}\, \bar{g}^{\mu\nu} F_\mu F_\nu, \tag{5.1}$$

with a gauge fixing parameter α and a De-Donder-type gauge condition given by

$$F_\mu = \bar{\nabla}^\nu h_{\mu\nu} - \frac{1+\beta}{4}\bar{\nabla}_\mu h^\nu{}_\nu. \tag{5.2}$$

We also have to include a ghost term $\mathcal{S}_{\mathrm{gh}}$ arising from the Faddeev-Popov procedure:

$$\mathcal{S}_{\mathrm{gh}} = Z_c \int_x \sqrt{\bar{g}}\, \bar{g}^{\mu\mu'}\bar{g}^{\nu\nu'}\bar{c}_{\mu'}\mathcal{M}_{\mu\nu}c_{\nu'}. \tag{5.3}$$

The Faddeev-Popov operator $\mathcal{M}_{\mu\nu}(\bar{g}, h)$ reads

$$\mathcal{M}_{\mu\nu} = \bar{\nabla}^\rho(g_{\mu\nu}\nabla_\rho + g_{\rho\nu}\nabla_\mu) - \bar{\nabla}_\mu\nabla_\nu. \tag{5.4}$$

If we choose to work in Landau gauge, where $\alpha \equiv 0$ and also set $\beta \equiv 0$, we are able to reduce the dynamical degrees of freedom to two, namely the transverse-traceless mode h^{TT} and the trace mode h^{Tr}. This choice allows us to simplify the full graviton propagator a lot. We are left with:

$$G_{k,hh} = \left(\Gamma_k^{(2)} + R_k\right)^{-1}_{hh} = \frac{32\pi}{Z_h}\begin{pmatrix} \frac{1}{\bar{\Delta}[1+r_k]-2\Lambda_k+\frac{2}{3}\mathcal{R}} & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & \frac{-\frac{8}{3}}{\bar{\Delta}[1+r_k]-\frac{4}{3}\Lambda_k} & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}. \tag{5.5}$$

For this calculation, the regularized 2-point-function has been determined using a regulator of the form

$$R_k = \Gamma_k^{(2)}\Big|_{\bar{R}=0,\Lambda=0} \cdot r_k\left(\frac{\bar{\Delta}}{k^2}\right), \tag{5.6}$$

where the shape function r_k is again of Litim-type (2.16). Since we already computed the contributions arising from the h^{TT} mode in the previous chapter, we only need to compute those from the trace mode h^{Tr} and the Faddeev-Popov ghosts associated with the graviton sector in the following.

5.1.1. Trace Mode

For the trace mode h^{Tr}, we take the propagator given in (5.5):

$$G_{k,h^{\mathrm{Tr}}h^{\mathrm{Tr}}} = \left(\frac{32\pi}{Z_h}\right)\left(-\frac{8}{3}\right)\frac{1}{\bar{\Delta}\left[1+r_k\right]-\frac{4}{3}\Lambda_k}. \tag{5.7}$$

The respective regulator term reads

$$R_{k,h^{\mathrm{Tr}}} = -\frac{3}{8}\left(\frac{Z_h}{32\pi}\right)\cdot\bar{\Delta}\cdot r_k\left(\frac{\bar{\Delta}}{k^2}\right). \tag{5.8}$$

We proceed by computing its scale derivative:

$$\partial_t R_{k,h^{\mathrm{Tr}}} = -\frac{3}{8}\left(\frac{Z_h}{32\pi}\right)\cdot\bar{\Delta}\left(\partial_t r_k - \eta_h r_k\right). \tag{5.9}$$

In total, the contribution from the r. h. s. of the flow equation is given by

$$\frac{1}{2}\,\mathrm{Tr}\left[G_{k,h^{\mathrm{Tr}}h^{\mathrm{Tr}}}\;\partial_t R_{k,h^{\mathrm{Tr}}}\right] = \frac{1}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k - \eta_h r_k\right)}{\bar{\Delta}\left[1+r_k\right]-\frac{4}{3}\Lambda_k}\right]. \tag{5.10}$$

To evaluate the functional trace, we make again use of the heat-kernel techniques we already used for the computation in h^{TT} approximation:

$$\begin{aligned}\frac{1}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k - \eta_h r_k\right)}{\bar{\Delta}\left[1+r_k\right]-\frac{4}{3}\Lambda_k}\right] &= \frac{1}{2}\frac{1}{(4\pi^2)}\left[\int_x\sqrt{g}\;\Phi_2^1\left(-\frac{4}{3}\Lambda_k\right)+\frac{1}{6}\int_x\sqrt{g}\;\bar{\mathcal{R}}\;\Phi_1^1\left(-\frac{4}{3}\Lambda_k\right)\right] \\ &= \frac{1}{2}\frac{1}{(4\pi)^2}\int_x\sqrt{g}\left[\frac{1}{1-\frac{4}{3}\lambda_k}\left(\left(1-\frac{\eta_h}{6}\right)+\frac{\bar{\mathcal{R}}}{3}\left(1-\frac{\eta_h}{4}\right)\right)\right].\end{aligned} \tag{5.11}$$

In the last step, we evaluated the threshold functions using equation (4.17). This already completes the computation of the additional contribution from the trace mode.

5.1.2. Faddeev-Popov Ghosts for the Graviton Sector

Since we work in a background field approximation, where we assume $g_{\mu\nu} = \bar{g}_{\mu\nu}$, we can simplify the calculations for this part a lot. Nevertheless, we have to be very careful since we will be confronted with a modified Laplace operator, i. e. a spin-1 Laplacian $\bar{\Delta}'^{(1)}_{\mu\nu}$, occurring as kinetic operator in the ghost two-point function. A more detailed discussion on how these modified Laplacians effect the values of the heat-kernel coefficients is presented at the end of appendix A.
First we have a look at the Faddeev-Popov operator (5.4). In background field approximation, $\mathcal{M}_{\mu\nu}$ is given by

$$\begin{aligned} \mathcal{M}_{\mu\nu}\Big|_{g=\bar{g}} &= -\bar{g}_{\mu\nu}\bar{\nabla}^2 + \bar{\nabla}_\nu\bar{\nabla}_\mu - \bar{\nabla}_\mu\bar{\nabla}_\nu \\ &= -\bar{g}_{\mu\nu}\bar{\nabla}^2 - \left[\bar{\nabla}_\mu, \bar{\nabla}_\nu\right]. \end{aligned} \tag{5.12}$$

The minus sign in front of the first term arises from integrating by parts to obtain the $\bar{\nabla}^2$ operator. As usual we assumed vanishing boundary terms. Inserting this result into the ghost action (5.3) yields

$$\begin{aligned} \mathcal{S}_{\rm gh} &= Z_c \int_x \sqrt{\bar{g}}\, \bar{c}^\mu \left[-\bar{g}_{\mu\nu}\bar{\nabla}^2 - \left[\bar{\nabla}_\mu, \bar{\nabla}_\nu\right]\right] c^\nu \\ &\overset{\rm (B.8)}{=} Z_c \int_x \sqrt{\bar{g}}\, \bar{c}^\mu \underbrace{\left[-\bar{g}_{\mu\nu}\bar{\nabla}^2 - \bar{R}_{\mu\nu}\right]}_{=:\, \bar{\Delta}'^{(1)}_{\mu\nu}} c^\nu. \end{aligned} \tag{5.13}$$

Note, that the bar referring to the background field should not be confused with the bar in the notation for the conjugated ghost. We can directly read off the ghost two-point function:

$$\left[\Gamma^{(2)}_{\bar{c}c}\right]_{\mu\nu} = \frac{\delta^2 \mathcal{S}_{\rm gh}}{\delta c\, \delta\bar{c}} = Z_c \cdot \bar{\Delta}'^{(1)}_{\mu\nu}. \tag{5.14}$$

With the following choice of the regulator

$$\left[R_{k,c}\right]_{\mu\nu} = Z_c \cdot \bar{\Delta}'^{(1)}_{\mu\nu} \cdot r_k\left(\frac{\bar{\Delta}'^{(1)}}{k^2}\right) \tag{5.15}$$

and its scale derivative

$$\left[\partial_t R_{k,c}\right]_{\mu\nu} = Z_c \cdot \bar{\Delta}'^{(1)}_{\mu\nu} \left(\partial_t r_k - \eta_c r_k\right), \tag{5.16}$$

we arrive at the regularized two-point function for the ghosts:

$$\left[\Gamma^{(2)}_{k,\bar{c}c}\right]_{\mu\nu} = \left[\Gamma^{(2)}_{\bar{c}c} + R_{k,c}\right]_{\mu\nu} = Z_c \cdot \bar{\Delta}'^{(1)}_{\mu\nu}\left(1 + r_k\left(\frac{\bar{\Delta}'^{(1)}}{k^2}\right)\right). \tag{5.17}$$

With equation (A.28), we find the new coefficients for the heat-kernel expansion for the modified Laplacian:

$$\begin{aligned}\operatorname{Tr}\mathbf{b}_0\left(\bar{\Delta}'^{(1)}\right) &= 4\\ \operatorname{Tr}\mathbf{b}_2\left(\bar{\Delta}'^{(1)}\right) &= \frac{5}{3}\bar{\mathcal{R}}.\end{aligned}\tag{5.18}$$

All together, we compute the following contribution:

$$\begin{aligned}-\operatorname{Tr}\left[G_{k,\bar{c}c}\,\partial_t R_{k,c}\right] &= -\operatorname{Tr}\left[\frac{\bar{\Delta}'^{(1)}_{\mu\nu}\left(\partial_t r_k - \eta_c r_k\right)}{\bar{\Delta}'^{(1)}_{\mu\nu}\left(1+r_k\right)}\right]\\ &= -\frac{1}{(4\pi^2)}\left[4\int_x \sqrt{\bar{g}}\,\Phi^1_2(0) + \frac{5}{3}\int_x \sqrt{\bar{g}}\,\bar{\mathcal{R}}\,\Phi^1_1(0)\right]\\ &= -\frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[4\left(1-\frac{\eta_c}{6}\right) + \frac{10}{3}\bar{\mathcal{R}}\left(1-\frac{\eta_c}{4}\right)\right].\end{aligned}\tag{5.19}$$

The minus sign in front of the trace is due to the Grassmannian nature of the ghost fields. With the completion of this calculation, we are ready to consider minimally-coupled matter fields.

5.2. Matter Contributions in Background Field Approximation

The inclusion of matter in this theory setting is in principle straightforward. We extend our truncation (4.1) by including an additional matter term:

$$\Gamma_k = \Gamma_{\mathrm{EH}} + \mathcal{S}_{\mathrm{gf}} + \mathcal{S}_{\mathrm{gh}} + \Gamma_{\mathrm{matter}},\tag{5.20}$$

where Γ_{matter} consists of scalar, fermion and gauge field contributions, denoted with $\mathcal{S}_S, \mathcal{S}_D$ and $\mathcal{S}_V$ respectively:

$$\Gamma_{\mathrm{matter}} = \mathcal{S}_S + \mathcal{S}_D + \mathcal{S}_V.\tag{5.21}$$

The different actions will be specified later on, every matter type will be treated separately. For conventions regarding the choice of the respective regulators and the general structure of this calculation, we are following [6].

In this truncation we have two essential couplings, G and Λ, and five inessential[14] wave function renormalizations Z_Ψ with $\Psi = (h, c, S, D, V)$. As before, the wave function renormalizations Z_Ψ do not enter the β-functions for G and Λ directly, but are still present in a non-trivial way via the anomalous dimension η_Ψ, defined as

$$\eta_\Psi = -\partial_t \ln Z_\Psi.\tag{5.22}$$

14. Inessential in this sense means, that they can be eliminated by field rescalings.

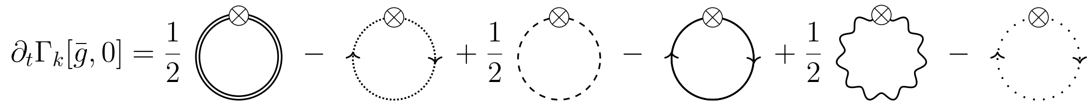


Figure 5.1. Flow equation (5.25) for the average effective action Γ_k including different matter contributions in diagrammatic representation. The double, densely dotted, dashed, solid, wiggly and loosely dotted lines correspond to the graviton, graviton ghost, scalar, fermion, gauge field and gauge field ghost propagators, respectively. The crossed circles denote the insertion of the respective regulator.

For the scalar and gauge field regulators we choose

$$R_{k,S/V}(z) = Z_{S/V} \cdot \mathbb{1}_{S/V} \cdot \tilde{\Delta} \cdot r_k \left(\frac{\tilde{\Delta}}{k^2} \right) , \tag{5.23}$$

where $\tilde{\Delta} = -\nabla^2 \mathbb{1}_\Psi + \mathbf{E}_\Psi$ is a modified Laplacian, occurring as kinetic operator in the different matter field actions. The Litim-type shape function r_k is in this case the same as the one defined in equation (2.16), now as a function of the modified Laplacian $\tilde{\Delta}$. The regulator choice for the Dirac fermions is slightly different, details are discussed in the respective subsection. Nevertheless, we already present the values of $\mathbf{E}_\Psi$ for all three kinetic operators:

$$\mathbf{E}_\Psi = \begin{cases} 0 & \text{for } \Psi = S \\ \frac{\mathcal{R}}{4} & \text{for } \Psi = D \\ R^\mu{}_\nu & \text{for } \Psi = V. \end{cases} \tag{5.24}$$

After having introduced the setup for the following calculation, we are now able to determine the different contributions from the matter fields step by step, by evaluating the functional traces ocurring on the r. h. s. of the flow equation separately. For the matter configuration in our setting the complete flow equation reads

$$\begin{aligned} \partial_t \Gamma_k = \frac{1}{2} \operatorname{Tr}\Big[G_k \, \partial_t R_k\Big]_{hh} - \operatorname{Tr}\Big[G_k \, \partial_t R_k\Big]_{\bar{c}c} + \frac{1}{2} \operatorname{Tr}\Big[G_k \, \partial_t R_k\Big]_{\phi\phi} \\ - \operatorname{Tr}\Big[G_k \, \partial_t R_k\Big]_{\bar{\psi}\psi} + \frac{1}{2} \operatorname{Tr}\Big[G_k \, \partial_t R_k\Big]_{AA} - \operatorname{Tr}\Big[G_k \, \partial_t R_k\Big]_{\bar{C}C} \end{aligned} \tag{5.25}$$

As before, the minus signs arise due to the fact, that the ghosts and the fermions are Grassmann fields. A digrammatic representation of the flow equation (5.25) is depicted in figure (5.1).

For our computation in the background field approximation, we expand the different actions on some background $\bar{g}_{\mu\nu}$ and drop all contributions of $\mathcal{O}(h)$.

5.2.1. Scalar Fields

The action for N_S scalar fields, minimally coupled to gravity reads

$$
\begin{aligned}
\mathcal{S}_S &= \frac{Z_S}{2} \int_x \sqrt{g}\, g^{\mu\nu} \sum_{i=1}^{N_\mathrm{S}} \partial_\mu \phi^i \partial_\nu \phi^i \\
&= \frac{Z_S}{2} \int_x \sqrt{\bar{g}}\, \bar{g}^{\mu\nu} \sum_{i=1}^{N_\mathrm{S}} \partial_\mu \phi^i \partial_\nu \phi^i + \mathcal{O}(h) \\
&= \frac{Z_S}{2} \int_x \sqrt{\bar{g}} \sum_{i=1}^{N_\mathrm{S}} \phi^i \left(-\bar{\nabla}^2\right) \phi^i + \mathcal{O}(h).
\end{aligned}
\tag{5.26}
$$

In the last step, we use integration by parts and assume vanishing boundary terms. Since $\mathbf{E} = 0$ for scalars, we use the initial definition of the Laplacian $\bar{\Delta} = -\bar{\nabla}^2$ for further calculations. These simple manipulations directly allow us to read off the corresponding two-point function

$$
\Gamma^{(2)}_{\phi\phi} = \frac{\delta^2 \mathcal{S}_S}{\delta\phi^i\, \delta\phi^j} = Z_S \cdot \bar{\Delta} \cdot \mathbb{1}_S + \mathcal{O}(h), \tag{5.27}
$$

where $\mathbb{1}_S$ has to be understood as the identity in field space. Using the regulator defined in (5.23), we find the regularized two-point-function:

$$
\Gamma^{(2)}_{k,\phi\phi} = \left[\Gamma^{(2)}_{\phi\phi} + R_{k,S}\right] = Z_S \cdot \bar{\Delta} \cdot \mathbb{1}_S \left(1 + r_k\left(\frac{\bar{\Delta}}{k^2}\right)\right). \tag{5.28}
$$

This expression is already diagonal in field space, meaning we are directly able to invert it to obtain the propagator. Together with the scale derivative of the regulator

$$
\partial_t R_{k,S} = Z_S \cdot \mathbb{1}_S \cdot \bar{\Delta}\left(\partial_t r_k - \eta_S r_k\right), \tag{5.29}
$$

we can start to evaluate the r. h. s. of the flow equation:

$$
\begin{aligned}
\frac{1}{2}\,\mathrm{Tr}\left[G_{k,\phi\phi}\, \partial_t R_{k,S}\right] &= \frac{1}{2}\,\mathrm{Tr}\left[\frac{Z_S \cdot \bar{\Delta}\left(\partial_t r_k - \eta_s r_k\right)}{Z_S \cdot \bar{\Delta}\left(1 + r_k\right)} \mathbb{1}_S\right] \\
&= \frac{N_S}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k - \eta_s r_k\right)}{\bar{\Delta}\left(1 + r_k\right)}\right].
\end{aligned}
\tag{5.30}
$$

Here, we already performed the trace operation on the internal indices, leading to an overall factor of N_S. The functional trace is again evaluated using heat-kernel techniques.

$$
\begin{aligned}
\frac{N_S}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}\left(\partial_t r_k - \eta_s r_k\right)}{\bar{\Delta}\left(1+r_k\right)}\right] &= \frac{N_S}{2}\frac{1}{(4\pi^2)}\left[\int_x \sqrt{\bar{g}}\,\Phi^1_2(0) + \frac{1}{6}\int_x \sqrt{\bar{g}}\,\bar{\mathcal{R}}\,\Phi^1_1(0)\right] \\
&= \frac{N_S}{2}\frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[\left(1-\frac{\eta_S}{6}\right) + \frac{\bar{\mathcal{R}}}{3}\left(1-\frac{\eta_S}{4}\right)\right].
\end{aligned}
\tag{5.31}
$$

5.2.2. Fermionic Fields

For the fermionic contribution, we proceed slightly different. First, we present the action for N_D minimally coupled Dirac fermions:

$$
\begin{aligned}
\mathcal{S}_D &= iZ_D \int_x \sqrt{g}\,\sum_{i=1}^{N_D} \bar{\psi}^i \nabla\!\!\!/\, \psi^i \\
&= iZ_D \int_x \sqrt{\bar{g}}\,\sum_{i=1}^{N_D} \bar{\psi}^i \bar{\nabla\!\!\!/}\, \psi^i + \mathcal{O}(h).
\end{aligned}
\tag{5.32}
$$

The Dirac operator $\nabla\!\!\!/$ satisfies $\left(i\nabla\!\!\!/\right)^2 = -\nabla^2 + \frac{\mathcal{R}}{4} =: \Delta_{(1/2)}$. Once again, the notation for the conjugated field $\bar{\psi} = \psi^\dagger \mathfrak{h}$[15] should not be confused with the bar referring to the background field. As usual, the slashed notation implies contraction with gamma matrices[16], i.e. $\nabla\!\!\!/ = \gamma^\mu \nabla_\mu$.

In principle, this allows us to read off the fermion two-point function

$$
\Gamma^{(2)}_{\bar{\psi}\psi} = \frac{\delta^2 \mathcal{S}_D}{\delta\psi^i\,\delta\bar{\psi}^j} = Z_D \cdot \left(i\bar{\nabla\!\!\!/}\right) \cdot \mathbb{1}_D + \mathcal{O}(h).
\tag{5.33}
$$

For the regularized two-point-function, we choose the following form:

$$
\Gamma^{(2)}_{k,\bar{\psi}\psi} = \left[\Gamma^{(2)}_{\bar{\psi}\psi} + R_{k,D}\right] = Z_D \cdot \mathbb{1}_D \cdot \left(i\bar{\nabla\!\!\!/}\right)\left(1 + r_{k,D}\left(\frac{\bar{\Delta}_{(1/2)}}{k^2}\right)\right).
\tag{5.34}
$$

This allows to determine the expression for $R_{k,D}$ as the difference between the regularized and the initial two-point function:

$$
\begin{aligned}
R_{k,D} &= \Gamma^{(2)}_{k,\bar{\psi}\psi} - \Gamma^{(2)}_{\bar{\psi}\psi} \\
&= Z_D \cdot \left(i\bar{\nabla\!\!\!/}\right) \cdot r_{k,D} \cdot \mathbb{1}_D.
\end{aligned}
\tag{5.35}
$$

15. The spin metric $\mathfrak{h}$ satisfies $|\det \mathfrak{h}| = 1$ and $\mathfrak{h}^{-1} = -\mathfrak{h}$.
16. In the discussion of Dirac fermions, the gamma matrices $\left\{\gamma^0, \gamma^1, \gamma^2, \gamma^3\right\}$ are a set of complex valued matrices, that constitute an irreducible representation of the Clifford algebra, defined by the anticommutation relation $\{\gamma_\mu, \gamma_\nu\} = 2g_{\mu\nu}\mathbb{1}_{d_\gamma \times d_\gamma}$, with $d_\gamma = 2^{\lfloor d/2 \rfloor}$. A more formal treatment of fermions in curved spacetimes is presented in [13].

When computing fermion propagators in other theory settings, it follows quite naturally to consider the Dirac dispersion as the "square root" of the scalar Klein-Gordon dispersion. For a more detailed discussion of this idea in the context of Fermi-Bose mixtures, we refer to chapter 2 of [16].
This assumption allows us to express the shape function for the fermions as a function of the scalar shape function:

$$\left(1 + r_{k,D}\right)^2 = 1 + r_{k,S} \qquad \longrightarrow \qquad r_{k,D} = \sqrt{1 + r_{k,S}} - 1 \tag{5.36}$$

Inserting this relation into the expression for the two-point function (5.34) and the regulator (5.35), we arrive at:

$$\begin{aligned} \Gamma^{(2)}_{k,\bar{\psi}\psi} &= \sqrt{1 + r_{k,S}} \cdot Z_D \cdot \left(i\bar{\nabla\!\!\!/}\right) \cdot \mathbb{1}_D \\ R_{k,D} &= \left(\sqrt{1 + r_{k,S}} - 1\right) \cdot Z_D \cdot \left(i\bar{\nabla\!\!\!/}\right) \cdot \mathbb{1}_D. \end{aligned} \tag{5.37}$$

To invert the Dirac operator $\left(i\nabla\!\!\!/\right)$, we use

$$\left(i\nabla\!\!\!/\right)^{-1} = \left(i\nabla\!\!\!/\right)^{-1} \cdot \left(\left(i\nabla\!\!\!/\right)^{-1} \cdot \left(i\nabla\!\!\!/\right)\right) = \left(\Delta_{(1/2)}\right)^{-1} \cdot \left(i\nabla\!\!\!/\right) \tag{5.38}$$

and therefore the full fermion propagator reads

$$G_{k,\bar{\psi}\psi} = \frac{\left(i\bar{\nabla\!\!\!/}\right)}{Z_D \cdot \bar{\Delta}_{(1/2)} \sqrt{1 + r_{k,S}}} \mathbb{1}_D. \tag{5.39}$$

For the scale derivative of the regulator we find

$$\begin{aligned} \partial_t R_{k,D} &= Z_D \cdot \mathbb{1}_D \left(i\bar{\nabla\!\!\!/}\right) \left(\partial_t r_{k,D} - \eta_D r_{k,D}\right) \\ &= Z_D \cdot \mathbb{1}_D \left(i\bar{\nabla\!\!\!/}\right) \left(\frac{\partial_t r_{k,S}}{2\sqrt{1 + r_{k,S}}} - \eta_D \left(\sqrt{1 + r_{k,S}} - 1\right)\right) \end{aligned} \tag{5.40}$$

Since we have to deal with the modified spin-$\frac{1}{2}$ Laplacian $\Delta_{(1/2)}$, we have to be careful with the heat-kernel coefficients once again. Using equation (A.28), we find:

$$\begin{aligned} \mathrm{Tr}\,\mathbf{b}_0 \left(\Delta_{(1/2)}\right) &= 4 \\ \mathrm{Tr}\,\mathbf{b}_2 \left(\Delta_{(1/2)}\right) &= \frac{5}{12} \bar{\mathcal{R}}. \end{aligned} \tag{5.41}$$

Altogether, we can start to evaluate the r. h. s. of the flow equation:

$$
\begin{aligned}
-\operatorname{Tr}\left[G_{k,\bar{\psi}\psi}\,\partial_t R_{k,D}\right] &= -\operatorname{Tr}\left[\frac{Z_D\cdot\bar{\Delta}_{(1/2)}\cdot\left(\frac{\partial_t r_{k,S}}{2\sqrt{1+r_{k,S}}}-\eta_D\left(\sqrt{1+r_{k,S}}-1\right)\right)}{Z_D\cdot\bar{\Delta}_{(1/2)}\sqrt{1+r_{k,S}}}\mathbb{1}_D\right] \\
&= -N_D\operatorname{Tr}\left[\frac{k^2\cdot\theta\left(k^2-\bar{\Delta}_{(1/2)}\right)}{\bar{\Delta}_{(1/2)}\left(1+r_{k,S}\right)}+\eta_D\left(\frac{1}{\sqrt{1+r_{k,S}}}-1\right)\right],
\end{aligned}
\tag{5.42}
$$

The Heaviside function θ enters via the scale derivative of $r_{k,S}$. Since we are working with a different shape function, we won't use the definition of the threshold functions this time. Nevertheless, it is still possible, to solve this functional trace in terms of the heat kernel expansion. For this explicit trace, we solved the integrals occurring in the associated Q-functionals analytically. The most important steps of this calculation can be found in appendix B. The final contribution after solving all four integrals is given by:

$$
-\operatorname{Tr}\left[G_{k,\bar{\psi}\psi}\partial_t R_{k,D}\right] \overset{(B.6)}{=} -N_D\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\left[2\left(1-\frac{\eta_D}{3}\right)+\frac{5}{12}\bar{\mathcal{R}}\left(1-\frac{\eta_D}{2}\right)\right]. \tag{5.43}
$$

5.2.3. Gauge Fields

The structure of the gauge field contribution is more complex than for the other fields. This is due to the fact, that we have to employ a gauge fixing procedure w. r. t. the background field $\bar{g}_{\mu\nu}$. This ensures gauge invariance w. r. t. background gauge transformations. The action for N_V gauge fields, minimally coupled to gravity reads

$$
\begin{aligned}
\mathcal{S}_V = &\frac{Z_{\mathrm{V}}}{4}\int_x\sqrt{g}\sum_{i=1}^{N_{\mathrm{V}}}g^{\mu\nu}g^{\kappa\lambda}F^i_{\mu\kappa}F^i_{\nu\lambda}+\frac{Z_{\mathrm{V}}}{2\xi}\int_x\sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}}\left(\bar{g}^{\mu\nu}\bar{\nabla}_\mu A^i_\nu\right)^2 \\
&+\int_x\sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}}\bar{C}_i(-\bar{\nabla}^2)C_i,
\end{aligned}
\tag{5.44}
$$

where the second term is the gauge fixing term with gauge parameter ξ and the third term is the associated ghost term. Since the two-point function is obtained from a functional derivative w. r. t. the fields A^i, we have to evaluate the ghost-term separately. We start by manipulating the first term:

$$
\begin{aligned}
\frac{Z_V}{4}\int_x\sqrt{g}\sum_{i=1}^{N_{\mathrm{V}}}g^{\mu\nu}g^{\kappa\lambda}F^i_{\mu\kappa}F^i_{\nu\lambda} &= \frac{Z_V}{4}\int_x\sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}}\bar{g}^{\mu\nu}\bar{g}^{\kappa\lambda}\bar{F}^i_{\mu\kappa}\bar{F}^i_{\nu\lambda}+\mathcal{O}(h) \\
&\overset{(B.7)}{=} \frac{Z_V}{2}\int_x\sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}}A^i_\lambda\left[\bar{\nabla}^\mu\bar{\nabla}^\lambda-\bar{g}^{\mu\lambda}\bar{\nabla}^2\right]A^i_\mu+\mathcal{O}(h).
\end{aligned}
\tag{5.45}
$$

The steps we skipped can be found in appendix B. For the gauge fixing term we find:

$$\begin{aligned}\frac{Z_{\mathrm{V}}}{2\xi}\int_x \sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}}\left(\bar{g}^{\mu\nu}\bar{\nabla}_\mu A^i_\nu\right)^2 &= \frac{Z_{\mathrm{V}}}{2\xi}\int_x \sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}}\bar{g}^{\mu\nu}\bar{\nabla}_\mu A^i_\nu g^{\kappa\lambda}\bar{\nabla}_\kappa A^i_\lambda \\ &= \frac{Z_{\mathrm{V}}}{2\xi}\int_x \sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}} A^i_\lambda\left[-\bar{\nabla}^\lambda\bar{\nabla}^\mu\right]A^i_\mu .\end{aligned} \tag{5.46}$$

In the last step, we integrated by parts and assumed vanishing boundary terms. This allows us to write

$$\mathcal{S}_V = \frac{Z_V}{2}\int_x \sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}} A^i_\lambda\left[-\bar{g}^{\mu\lambda}\bar{\nabla}^2+\bar{\nabla}^\mu\bar{\nabla}^\lambda-\frac{1}{\xi}\bar{\nabla}^\lambda\bar{\nabla}^\mu\right]A^i_\mu \;+\; \text{ghost term} \tag{5.47}$$

In Feynman gauge, where we set $\xi \equiv 1$, this simplifies to

$$\begin{aligned}\mathcal{S}_V &= \frac{Z_V}{2}\int_x \sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}} A^i_\lambda\left[-\bar{g}^{\mu\lambda}\bar{\nabla}^2+\left[\bar{\nabla}^\mu,\bar{\nabla}^\lambda\right]\right]A^i_\mu \;+\; \text{ghost term} \\ &\overset{(\mathrm{B.8})}{=} \frac{Z_V}{2}\int_x \sqrt{\bar{g}}\sum_{i=1}^{N_{\mathrm{V}}} A^i_\lambda\left[-\bar{g}^{\mu\lambda}\bar{\nabla}^2+\bar{R}^{\mu\lambda}\right]A^i_\mu \;+\; \text{ghost term}\end{aligned} \tag{5.48}$$

In this form, we are again directly able to read off the two-point-function:

$$\left[\Gamma^{(2)}_{AA}\right]^{\mu\nu} = \left[\frac{\delta^2\mathcal{S}_V}{\delta A^i\,\delta A^j}\right]^{\mu\nu} = Z_V\,\underbrace{\left[-\bar{g}^{\mu\lambda}\bar{\nabla}^2+\bar{R}^{\mu\lambda}\right]}_{=:\,\bar{\Delta}^{\mu\nu}_{(1)}}\mathbb{1}_V+\mathcal{O}(h), \tag{5.49}$$

where $\bar{\Delta}^{\mu\nu}_{(1)}$ is another modified spin-one Laplacian, but this time with a different sign in front of the Ricci tensor. With the respective regulator we find

$$\left[\Gamma^{(2)}_{k,AA}\right]^{\mu\nu} = \left[\Gamma^{(2)}_{AA}+R_{k,V}\right]^{\mu\nu} = Z_V\cdot\bar{\Delta}^{\mu\nu}_{(1)}\cdot\mathbb{1}_V\left(1+r_k\left(\frac{\bar{\Delta}_{(1)}}{k^2}\right)\right) \tag{5.50}$$

and

$$\left[\partial_t R_{k,V}\right]^{\mu\nu} = Z_V\cdot\bar{\Delta}^{\mu\nu}_{(1)}\cdot\mathbb{1}_V\left(\partial_t r_k-\eta_V r_k\right). \tag{5.51}$$

In this case, the heat-kernel coefficients for $\bar{\Delta}^{\mu\nu}_{(1)}$ read

$$\begin{aligned}\operatorname{Tr}\mathbf{b}_0\left(\bar{\Delta}_{(1)}\right) &= 4 \\ \operatorname{Tr}\mathbf{b}_2\left(\bar{\Delta}_{(1)}\right) &= -\frac{\bar{\mathcal{R}}}{3}\end{aligned} \tag{5.52}$$

and therefore, the result for the heat-kernel expansion for the gauge fields is given by:

$$\begin{aligned}
\frac{1}{2}\,\mathrm{Tr}\left[\frac{Z_V\cdot\bar{\Delta}^{\mu\nu}_{(1)}\left(\partial_t r_k-\eta_V r_k\right)}{Z_V\cdot\bar{\Delta}^{\mu\nu}_{(1)}(1+r_k)}\mathbb{1}_V\right] &= \frac{N_V}{2}\,\mathrm{Tr}\left[\frac{\bar{\Delta}^{\mu\nu}_{(1)}\left(\partial_t r_k-\eta_V r_k\right)}{\bar{\Delta}^{\mu\nu}_{(1)}(1+r_k)}\right]\\
&= \frac{N_V}{2}\frac{1}{(4\pi)^2}\left[\int_x\sqrt{\bar{g}}\;\Phi^1_2(0)-\frac{1}{3}\int_x\sqrt{\bar{g}}\;\bar{\mathcal{R}}\;\Phi^1_1(0)\right]\\
&= \frac{N_V}{2}\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\left[\left(1-\frac{\eta_V}{6}\right)-\frac{2}{3}\bar{\mathcal{R}}\left(1-\frac{\eta_V}{4}\right)\right]
\end{aligned} \tag{5.53}$$

To finish the calculation of the gauge field contribution, we need to take the associated ghost term into account. Fortunately, it has already the desired form, where we can directly read off the two-point function:

$$\Gamma^{(2)}_{\bar{C}C}=\frac{\delta^2\mathcal{S}_V}{\delta C^i\,\delta\bar{C}^j}=\mathbb{1}_V\cdot\bar{\Delta}. \tag{5.54}$$

Note, that we have the usual Laplacian $\bar{\Delta}=-\bar{\nabla}^2$ as kinetic operator and that no additional wave function renormalization was introduced. The ghost regulator $R_{k,\mathrm{gh}}$ is the same as for the scalar fields and therefore the regularized two-point function reads

$$\Gamma^{(2)}_{k,\bar{C}C}=\left[\Gamma^{(2)}_{\bar{C}C}+R_{k,\mathrm{gh}}\right]=\bar{\Delta}\cdot\mathbb{1}_V\left(1+r_k\left(\frac{\bar{\Delta}}{k^2}\right)\right). \tag{5.55}$$

In absence of a wave function renormalization, the scale derivative only acts on the shape function r_k and therefore the final contribution is given by

$$\begin{aligned}
-\,\mathrm{Tr}\left[\frac{\bar{\Delta}\partial_t r_k}{\bar{\Delta}\,(1+r_k)}\mathbb{1}_V\right] &= -N_V\,\mathrm{Tr}\left[\frac{\bar{\Delta}\partial_t r_k}{\bar{\Delta}\,(1+r_k)}\right]\\
&= -N_V\frac{1}{(4\pi)^2}\left[\int_x\sqrt{\bar{g}}\;\Phi^1_2(0)+\frac{1}{6}\int_x\sqrt{\bar{g}}\bar{\mathcal{R}}\;\Phi^1_1(0)\right]\\
&= -N_V\frac{1}{(4\pi)^2}\int_x\sqrt{\bar{g}}\left[1+\frac{1}{3}\bar{\mathcal{R}}\right].
\end{aligned} \tag{5.56}$$

In the next section, we combine the obtained results and present the final expressions for the β-functions.

5.3. Beta-Functions and Perturbative Approximation

We follow the same strategy as last chapter. Simply by comparing the different orders in curvature occurring on both sides of the flow equation, we are able to read off the β-function for the cosmological constant. We collect all terms of order $\sim \sqrt{\bar{g}}$ and find:

$$\partial_t \lambda_k = (\eta_h - 2)\,\lambda_k + \frac{g_k}{4\pi}\left[\left(\frac{5\left(1-\frac{\eta_h}{6}\right)}{1-2\lambda_k}\right) + \left(\frac{5\left(1-\frac{\eta_h}{6}\right)}{1-\frac{4}{3}\lambda_k}\right) - 8\left(1-\frac{\eta_c}{6}\right)\right. \\ \left. + N_S\left(1-\frac{\eta_S}{6}\right) - 4N_D\left(1-\frac{\eta_D}{3}\right) - N_V\left(1-\frac{\eta_V}{6}\right)\right]. \tag{5.57}$$

The expression for the graviton anomalous dimension follows analogously from a comparison of all terms of order $\sim \sqrt{\bar{g}}\bar{\mathcal{R}}$:

$$\eta_h = \frac{g_k}{\pi}\left[\frac{2\left(1-\frac{\eta_h}{4}\right)}{1-2\lambda_k} - \frac{10}{3}\frac{1-\frac{\eta_h}{6}}{(1-2\lambda_k)^2} + \frac{1}{6}\frac{1-\frac{\eta_h}{4}}{1-\frac{4}{3}\lambda_k} - \frac{10}{3}\left(1-\frac{\eta_c}{4}\right)\right. \\ \left. + \frac{N_S}{6}\left(1-\frac{\eta_S}{4}\right) - \frac{5N_D}{12}\left(1-\frac{\eta_D}{2}\right) - 2N_V\left(\frac{1}{3}-\frac{\eta_V}{12}\right)\right]. \tag{5.58}$$

This allows us also to determine also the β-function for the Newton coupling using equation (4.7) with the new expression for the graviton anomalous dimension:

$$\partial_t g_k = 2g_k + \frac{g_k^2}{\pi}\left[\frac{2\left(1-\frac{\eta_h}{4}\right)}{1-2\lambda_k} - \frac{10}{3}\frac{1-\frac{\eta_h}{6}}{(1-2\lambda_k)^2} + \frac{1}{6}\frac{1-\frac{\eta_h}{4}}{1-\frac{4}{3}\lambda_k} - \frac{10}{3}\left(1-\frac{\eta_c}{4}\right)\right. \\ \left. + \frac{N_S}{6}\left(1-\frac{\eta_S}{4}\right) - \frac{5N_D}{12}\left(1-\frac{\eta_D}{2}\right) - 2N_V\left(\frac{1}{3}-\frac{\eta_V}{12}\right)\right]. \tag{5.59}$$

To get a rough idea on how the different matter fields contribute to the RG flow, we want to perform a perturbative approximation, i. e. we neglect all contributions from the different anomalous dimensions and expand our β-functions up to second order in the couplings. This yields:

$$\begin{aligned} \partial_t g_k &\simeq 2g_k + \frac{g_k^2}{6\pi}\left(N_S - \frac{5}{2}N_D - 4N_V - 27\right) \\ \partial_t \lambda_k &\simeq -2\lambda_k + \frac{g_k}{4\pi}\Big(N_S - 4N_D - N_V + 2\Big) + \frac{g_k\lambda_k}{6\pi}\left(N_S - \frac{5}{2}N_D - 4N_V - 27\right) \end{aligned} \tag{5.60}$$

The integer numbers -27 and $+2$ represent the contributions from the gravitons and the Faddeev-Popov ghosts. In this approximation, the β-functions attain a non-trivial zero, corresponding to the Non-Gaussian fixed point we are interested in.
We find the following fixed point values:

$$\begin{aligned} g_k^* &= \frac{-12\pi}{N_S - \frac{5}{2}N_D - 4N_V - 27} \\ \lambda_k^* &= -\frac{3}{4}\frac{N_S - 4N_D - N_V + 2}{N_S - \frac{5}{2}N_D - 4N_V - 27}. \end{aligned} \tag{5.61}$$

In the following we present some plots to visualize the effects on the fixed point values when a single matter species is increased while the other numbers are fixed to zero.
The plots, depicted in figure (5.2) on the next page, indicate, that the scalar fields tend to destabilize the system, after a certain amount is surpassed. The values for both couplings increase drastically and will probably diverge. The fermionic and the gauge fields seem to be compatible with the theory, since they tend to decrease the values of both couplings. The value for λ^* is nearly constant over almost the whole depicted range for the gauge fields.

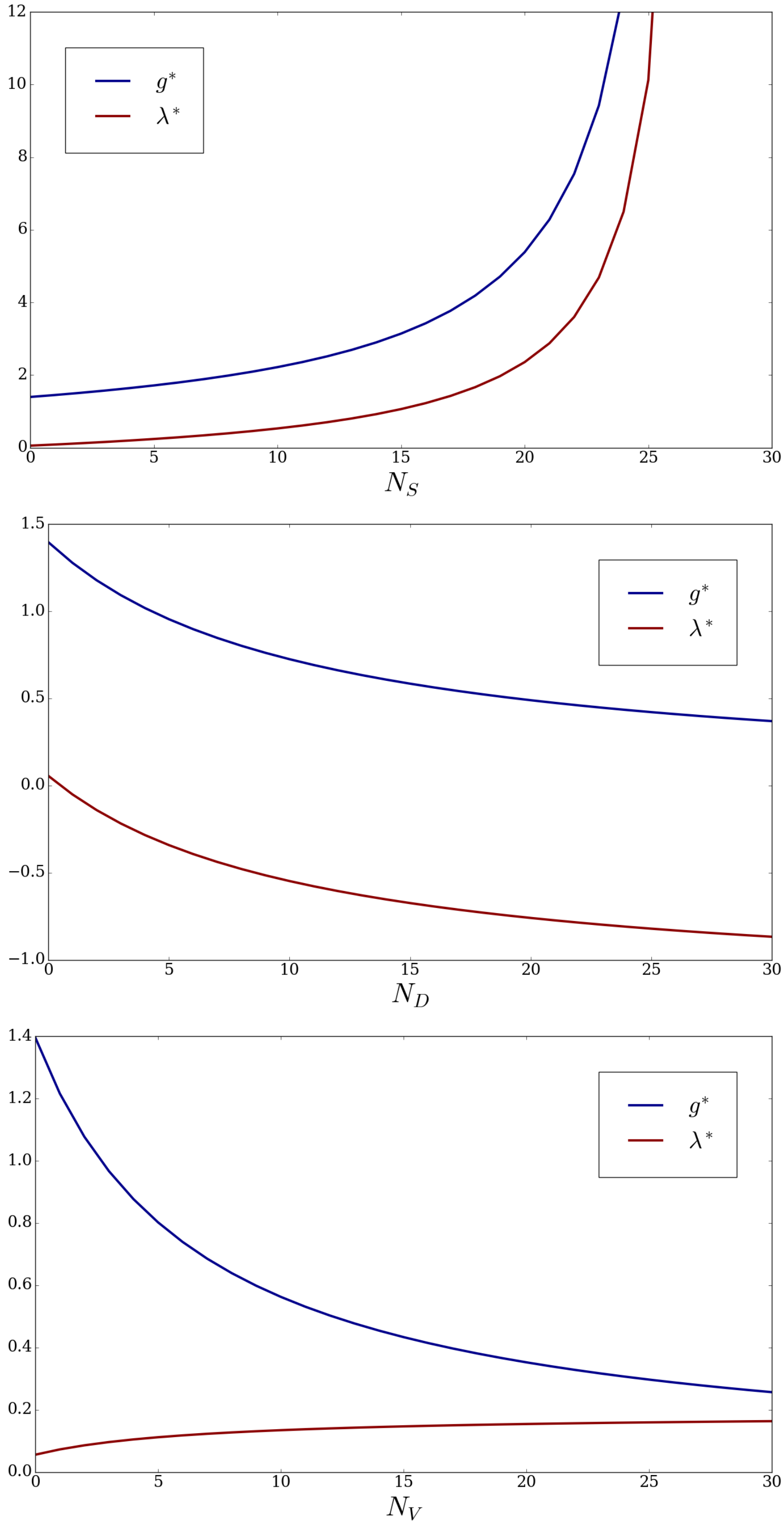


Figure 5.2. Fixed point values for increasing number of matter fields. Initially the matter content is set to zero, meaning we are only considering the respective matter type in each plot.

Chapter 6.

Background Independence in Quantum Gravity

Before we come to an end, we briefly want to discuss the issue of background independence in the context of quantum gravity and emphasize that our calculations in the background field approximation have to be treated very carefully. The formulas presented in the scope of this chapter are taken from [16].

Throughout this work we used a linear split of the metric into some background $\bar{g}_{\mu\nu}$ and a fluctuation field $h_{\mu\nu}$. This changes nothing about the fact, that in general, the flowing action $\Gamma_k = \Gamma_k[\bar{g}+h]$ is still a functional of the full metric g, but from a computational point of view it may be easier accessible, since we can expand it in powers of the fluctuations on the given background, i. e.

$$\Gamma_k\left[\bar{g}+h\right] = \Gamma_k\left[\bar{g}\right] + \Gamma_k^{(0,1)}\left[\bar{g}\right]\cdot h + \frac{1}{2}\Gamma_k^{(0,2)}\left[\bar{g}\right]\cdot h^2 + \mathcal{O}(h^3), \tag{6.1}$$

where we introduced the shorthand notation $\Gamma_k^{(n,m)} = \frac{\delta^{n+m}\Gamma_k}{\delta^n \bar{g}_{\mu\nu}\delta^m h_{\mu\nu}}$ to distinguish derivatives w. r. t. the background and the fluctuation field. But since this split is only a mathematical trick to be able to perform calculations in a more efficient way, there should be some relation between the correlations of the fluctuation field and those of the background field. This idea is encoded in the *Nielsen identities*, given by

$$\mathrm{NI} = \frac{\delta\Gamma}{\delta\bar{g}_{\mu\nu}} - \frac{\delta\Gamma}{\delta h_{\mu\nu}} - \left\langle \left[\frac{\delta}{\delta\bar{g}_{\mu\nu}} - \frac{\delta}{\delta\hat{h}_{\mu\nu}}\right]\Big(\mathcal{S}_{\mathrm{gf}} + \mathcal{S}_{\mathrm{gh}}\Big)\right\rangle = 0. \tag{6.2}$$

Here $h_{\mu\nu} = \langle\hat{h}_{\mu\nu}\rangle$. The difference between the background derivatives and the fluctuation derivatives is connected to derivatives of the gauge fixing sector. At finite k the regulator, which plays a crucial role in our approach to the subject, introduces another background dependence. This leads us to the *modified Nielsen identities*:

$$\mathrm{mNI} = \mathrm{NI} - \frac{1}{2}\,\mathrm{Tr}\left[\frac{\delta}{\delta\sqrt{g}}\frac{\delta\sqrt{\bar{g}}R_k[\bar{g}]}{\delta\bar{g}_{\mu\nu}}G_k\right] = 0. \tag{6.3}$$

In background field approximation we assume $\frac{\delta\Gamma_k}{\delta h} \approx \frac{\delta\Gamma_k}{\delta\bar{g}}$, this violates the Nielsen identities and therefore background independence is lost. This may seem contradictory at first sight, since in background field approximation we only deal with a single metric. Nevertheless, up until now, the background field approximation is somehow the standard approach to calculations in this area. As already mentioned in the introduction, more recently some progress has been made in surmounting the background field approximation. Results based on vertex expansions in powers of the fluctuation field, such as in (6.1), were used to determine the flow of the couplings by computation of the flow of the

n-point functions. Details can be found in [4, 14]. The results obtained from this approach differ to some extend strictly from the results found in computations in background field approximations. For future projects, one should keep these problems concerning the background field approximation in mind.

Chapter 7.

Conclusion

In this thesis, we investigated the Asymptotic Safety scenario for quantum gravity and studied its phase diagram within the Einstein-Hilbert truncation in a background field approximation. We used non-perturbative Functional Renormalization Group techniques to compute the running of Newton's constant g_k and of the cosmological constant λ_k.

In chapters 1 to 3 we introduced the background knowledge, needed for a general understanding of the conducted calculations. After briefly discussing the general idea of the Asymptotic Safety conjecture, in chapter 4 we solved the flow equation for a pure-gravity system in a transverse-traceless spin-two graviton approximation to get a first insight into the underlying structures of the theory setting. The mathematical tools we used to solve the flow equation, such as the York decomposition of the fluctuation field and the heat-kernel techniques to compute the functional traces, were introduced very detailed. We computed the β-functions for g_k and λ_k and determined the fixed points of the flow. Besides the non-interacting Gaussian fixed point at $(g_k, \lambda_k) = (0, 0)$ we found an UV-attractive fixed point at $(g_k, \lambda_k) = (0.86, 0.18)$, providing further evidence for the Asymptotic Safety scenario as a promising candidate for a non-perturbative renormalizable quantum field theory of gravity. In chapter 5 we extended our truncation: First of all, the previously neglected trace mode and the Faddeev-Popov ghosts associated with the graviton sector were included to complete the calculation from chapter 4. Then we investigated the impact of minimally coupled scalars, fermions and gauge fields on the Non-Gaussian fixed point. We explained in full detail how to solve the functional traces for all three matter types separately and presented the most important and insightful steps. After finishing the computation of all contributions, we were able to determine the β-functions for g_k and λ_k as a function of the number of matter fields. Neglecting all the contributions from the different anomalous dimensions and expanding the β-functions in some neighborhood of the Gaussian fixed point up to second order in the couplings, we qualitatively analyzed the behavior of the values for both couplings. We found out, that for an increasing amount of scalar fields, the values of the couplings tend to increase drastically. The fermionic fields and especially the gauge fields seem to have a stabilizing effect on the system. These results are almost in agreement with the results from earlier investigations, where similar conventions have been chosen, see e. g [6]. At the end, in chapter 6, we highlighted some problems associated with the background field approximation, e. g. the loss of background independence as a consequence of violating the Nielsen identities.

This work may not provide fundamentally new results, but it still demonstrates some of the central concepts and calculations related to the subject. Due to the fact, that we worked in a rather simple truncation and chose a Litim-type cutoff, we we able to solve all the problems in this thesis analytically. Interesting modifications of our setup could include the employment of more sophisticated regulators or the inclusion of higher-order curvature terms. Compared to more recent results, the outcomes of our calculations should be treated with care. Nevertheless, this thesis provides a suitable framework for further investigations of asymptotically safe quantum gravity. Based on our discussion in chapter 6, one may consider abandoning the background field approximation in future projects.

Appendix A.

Mathematical Background

In this part of the appendix we want to discuss some of the mathematical tools we used during the calculations presented in the scope of this thesis in a more formal manner. The part on the York decomposition is mainly inspired by [17], whereas the conventions for the heat-kernel computations are taken from [16] and extended for the matter part, using the conventions from [5].

A.1. York Decomposition

In the discussion of gauge theories, it is often very useful to decompose the gauge field A_μ into transversal and longitudinal parts:

$$A_\mu = A_\mu^{\mathrm{T}} + \nabla_\mu \phi. \tag{A.1}$$

The transversal part is characterized by the fact, that $\nabla^\mu A_\mu^{\mathrm{T}} = 0$. Using this decomposition, we are able to separate the pure gauge spin-0 degrees of freedom from the physical ones, contained in the spin-1 part A_μ^{T}.
Assuming vanishing boundary terms, integration by parts allows us to change the integration variables in the functional integral, i. e.

$$\int_x \sqrt{g}\, A_\mu A^\mu = \int_x \sqrt{g}\, A_\mu^{\mathrm{T}} A^{\mathrm{T},\mu} + \int_x \sqrt{g}\, \phi \left(-\nabla^2\right) \phi. \tag{A.2}$$

Note, that we have to take care of the Jacobian J of this variable transformation:

$$(\mathrm{d}A_\mu) \longrightarrow J \left(\mathrm{d}A_\mu^{\mathrm{T}}\right) (\mathrm{d}\phi)\,. \tag{A.3}$$

To be able to determine the Jacobian for our transformation, the integration measure needs to be normalized. A quite convenient choice is to evaluate the Gaussian integral over the different fields ψ and set the result to one:

$$\int (\mathrm{d}\psi) \exp\left\{-\int \mathrm{d}x\, \sqrt{g}\, \psi^2\right\} = 1, \tag{A.4}$$

where we are assuming an Euclidean signature and a curved background metric. With this condition we find:

$$1 = J \int \left(\mathrm{d}A_\mu^{\mathrm{T}}\right) \mathrm{e}^{-\int \mathrm{d}x \sqrt{g}\, A_\mu^{\mathrm{T}} A^{\mathrm{T},\mu}} \int (\mathrm{d}\phi)\, \mathrm{e}^{-\int \mathrm{d}x \sqrt{g}\, \phi\left(-\nabla^2\right)\phi} = J \left(\mathrm{det}'_\phi \left(-\nabla^2\right)\right)^{-1/2}. \tag{A.5}$$

This allows us to determine the Jacobian J as follows:

$$J = \left(\det{}'_{\phi}\left(-\nabla^2\right)\right)^{1/2}. \tag{A.6}$$

The prime denotes the fact, that the zero mode has to be removed, when computing the determinant to obtain a consistent result. Physically this is in accordance with the fact, that a constant ϕ does not contribute to A_μ.
For our computation in chapters 4 and 5, we used the background field method, where we assume a linear split of the *full* metric $g_{\mu\nu}$ into a background metric $\bar{g}_{\mu\nu}$ and a fluctuation field $h_{\mu\nu}$. There is an analogous way of decomposing the fluctuation field in the background field formalism. First, we split $h_{\mu\nu}$ into

$$h_{\mu\nu} = h^{\mathrm{T}}_{\mu\nu} + \frac{1}{d}\,\bar{g}_{\mu\nu} h, \tag{A.7}$$

where $h^{\mathrm{T}}_{\mu\nu}$ is traceless, i. e. $\bar{g}^{\mu\nu} h^{\mathrm{T}}_{\mu\nu} = 0$ and $h = \bar{g}^{\mu\nu} h_{\mu\nu}$. The traceless part can be further decomposed in flat space using the irreducible representations of the Lorentz group with spins 0, 1 and 2 respectively, but in our case a more sophisticated approach, the so-called *York decomposition* is chosen:

$$h_{\mu\nu} = h^{\mathrm{TT}}_{\mu\nu} + \bar{\nabla}_\mu \xi_\nu + \bar{\nabla}_\nu \xi_\mu + \left(\bar{\nabla}_\mu \bar{\nabla}_\nu - \frac{1}{d}\,\bar{g}_{\mu\nu}\bar{\nabla}^2\right)\sigma + \frac{1}{d}\,\bar{g}_{\mu\nu} h. \tag{A.8}$$

Here, $h^{\mathrm{TT}}_{\mu\nu}$ is a transverse-traceless, spin-2 degree of freedom, ξ_μ is transverse and carries a spin-1 d. o. f. and σ and h possess spin-0. As before, we want to find the Jacobian J for this variable transformation:

$$(\mathrm{d}h_{\mu\nu}) \longrightarrow J\left(\mathrm{d}h^{\mathrm{TT}}_{\mu\nu}\right)(\mathrm{d}\xi_\mu)(\mathrm{d}\sigma)(\mathrm{d}h). \tag{A.9}$$

This is again possible after specifying a suitable normalization of the functional measure as

$$\int (\mathrm{d}h_{\mu\nu}) \exp\left\{-\mathcal{G}(h,h)\right\} = 1, \tag{A.10}$$

where $\mathcal{G}$ is an inner product in the space of symmetric two-tensors, defined as

$$\begin{aligned}\mathcal{G}(h,h) &= \int_x \sqrt{\bar{g}}\,\left(h_{\mu\nu}h^{\mu\nu} + \frac{a}{2}h^2\right)\\ &= \int_x \sqrt{\bar{g}}\,\left[h^{\mathrm{TT}}_{\mu\nu}h^{\mathrm{TT},\mu\nu} + 2\xi_\mu\left(-\bar{\nabla}^2 - \frac{\bar{R}}{d}\right)\xi^\mu\right.\\ &\left.+\frac{d-1}{d}\sigma\left(-\bar{\nabla}^2\right)\left(-\bar{\nabla}^2 - \frac{\bar{R}}{d-1}\right)\sigma + \left(\frac{1}{d}+\frac{a}{2}\right)h^2\right]\end{aligned} \tag{A.11}$$

in the case of an Einstein type background metric[17]. This yields

$$J = \left(\det_\xi\left(-\bar{\nabla}^2 - \frac{R}{d}\right)\right)^{1/2} \left(\det'_\sigma\left(-\bar{\nabla}^2\right)\right)^{1/2} \left(\det_\sigma\left(-\bar{\nabla}^2 - \frac{R}{d-1}\right)\right)^{1/2} . \tag{A.12}$$

Note, that the prime has the same meaning and physical interpretation as in the previous case: If σ is constant, it does not contribute to $h_{\mu\nu}$.
In both cases, the decomposition of the general gauge field and the York decomposition of the fluctuation field, appropriate rescalings of the fields ϕ, ξ_μ and σ respectively, help us to cancel the non-trivial Jacobians and to achieve, that all modes have the same mass dimension. For the sake of completeness, we present the rescaled versions of the fields:

$$\hat{\phi} = \sqrt{-\nabla^2}\,\phi \tag{A.13}$$

$$\hat{\xi}_\mu = \sqrt{-\bar{\nabla}^2 - \frac{\bar{R}}{d}}\,\xi_\mu \tag{A.14}$$

$$\hat{\sigma} = \sqrt{-\bar{\nabla}^2}\sqrt{-\bar{\nabla}^2 - \frac{\bar{R}}{d-1}}\,\sigma. \tag{A.15}$$

The resulting graviton two-point function, after decomposition of the fluctuation field has the following structure:

$$\Gamma^{(2)}_{hh} = \begin{pmatrix} \Gamma^{(2)}_{h^{\mathrm{TT}}h^{\mathrm{TT}}} & 0 & 0 & 0 \\ 0 & \Gamma^{(2)}_{\xi\xi} & 0 & 0 \\ 0 & 0 & \Gamma^{(2)}_{h^{\mathrm{Tr}}h^{\mathrm{Tr}}} & \Gamma^{(2)}_{h^{\mathrm{Tr}}\sigma} \\ 0 & 0 & \Gamma^{(2)}_{\sigma h^{\mathrm{Tr}}} & \Gamma^{(2)}_{\sigma\sigma} \end{pmatrix} . \tag{A.16}$$

This concludes our discussion of the York decomposition, as a useful tool to simplify calculations in the background field method.

17. A metric is of Einstein type, if $R_{\mu\nu}$ is a constant multiple of $g_{\mu\nu}$, i. e. $R_{\mu\nu} = \frac{1}{d}\mathcal{R}g_{\mu\nu}$.

A.2. Heat-Kernel Techniques

We use heat-kernel techniques to evaluate the r. h. s. of the flow equation (2.25), where we need to compute the functional trace over functions depending on the Laplacian on a curved background. In general, the method can be understood as a curvature expansion on a flat background.
The general formula to compute such traces is given by

$$\mathrm{Tr}\, f(\Delta) = N \sum\!\!\!\!\!\!\int_{\ell} \rho(\ell) f(\lambda(\ell)), \tag{A.17}$$

with some normalization N, the spectral values $\lambda(\ell)$ and their corresponding multiplicities $\rho(\ell)$.
On flat backgrounds, the computation of (A.17) is simply a standard momentum integral, whereas on curved backgrounds, consider for example a four-sphere $\mathbb{S}^4$ with constant background curvature $r = \frac{\bar{\mathcal{R}}}{k^2} > 0$, the spectrum of the Laplacian is discrete and we need to sum over all spectral values.
For our example of $\mathbb{S}^4$, we have

$$\lambda(\ell) = \frac{\ell(3+\ell)}{12} r \qquad \text{and} \qquad \rho(\ell) = \frac{(2\ell+3)(\ell+2)!}{6\ell!}. \tag{A.18}$$

The normalization is then given by the inverse of the four-sphere-volume $(V_{\mathbb{S}^4})^{-1} = \frac{k^4 r^2}{384\pi^2}$. This leads us to the formula for our computation of the r. h. s. of the flow equation on a background with constant positive curvature:

$$\mathrm{Tr}\, f(\Delta) = \frac{k^4 r^2}{384\pi^2} \sum_{\ell=0}^{\infty} \frac{(2\ell+3)(\ell+2)!}{6\ell!} f\left(\frac{\ell(3+\ell)}{12} r\right). \tag{A.19}$$

This is called spectral sum. For large curvatures r the convergence of the series is rather fast, whereas in the limit $r \to 0$ one finds exponentially slow convergence.
The master equation for heat-kernel computations reads

$$\mathrm{Tr}\, f(\Delta) = \frac{1}{(4\pi)^{\frac{d}{2}}} \left[\mathbf{B}_0(\Delta) Q_2[f(\Delta)] + \mathbf{B}_2(\Delta) Q_1[f(\Delta)]\right] + \mathcal{O}\left(\mathcal{R}^2\right), \tag{A.20}$$

with the heat-kernel coefficients

$$\mathbf{B}_n(\Delta) = \int_x \sqrt{g}\ \mathrm{Tr}\, \mathbf{b}_n(\Delta) \tag{A.21}$$

and

$$Q_n[f(x)] = \frac{1}{\Gamma(n)} \int \mathrm{d}x\, x^{n-1} f(x). \tag{A.22}$$

For computations on $\mathbb{S}^4$, the values for the heat kernel coefficients $\mathbf{B}_n(\Delta)$ are presented in the following:

	TT	TV	S
$\mathrm{Tr}\,\mathbf{b}_0$	5	3	1
$\mathrm{Tr}\,\mathbf{b}_2$	$-\frac{5}{6}\mathcal{R}$	$\frac{1}{4}\mathcal{R}$	$\frac{1}{6}\mathcal{R}$

Table A.1. Heat-kernel coefficients for transverse-traceless tensors (TT), transverse vectors (TV) and scalars (S) for computations on $\mathbb{S}^4$.

The basic idea of the proof of equation (A.17) is based on the Laplace transform

$$f(\Delta) = \int_0^\infty \mathrm{d}s\ \mathrm{e}^{-s\Delta}\,\tilde{f}(s). \tag{A.23}$$

We insert this definition of the Laplace transform into equation (A.17) and find

$$\mathrm{Tr}\, f(\Delta) = \int_0^\infty \mathrm{d}s\ \tilde{f}(s)\,\mathrm{Tr}\,\mathrm{e}^{-s\Delta}\,. \tag{A.24}$$

The trace on the r. h. s. is explicitly the trace of the heat-kernel. We expand this term as follows:

$$\mathrm{Tr}\,\mathrm{e}^{-s\Delta} = \frac{1}{(4\pi)^{\frac{d}{2}}}\sum_{n=0}^{\infty} s^{\frac{n-d}{2}}\mathbf{B}_n(\Delta). \tag{A.25}$$

This is where the heat-kernel coefficients $\mathbf{B}_n$ become important. We proceed by inserting this expanded version of the heat-kernel trace into equation (A.24) and find:

$$\begin{aligned}
\mathrm{Tr}\, f(\Delta) &= \frac{1}{(4\pi)^{\frac{d}{2}}}\sum_{n=0}^{\infty}\mathbf{B}_n(\Delta)\int_0^\infty \mathrm{d}s\ s^{\frac{n-d}{2}}\tilde{f}(s)\\
&= \frac{1}{(4\pi)^{\frac{d}{2}}}\sum_{n=0}^{\infty}\frac{1}{\Gamma\left(\frac{d-k}{2}\right)}\mathbf{B}_n(\Delta)\int_0^\infty \mathrm{d}t\ t^{\frac{d-n}{2}-1}f(t)\\
&= \frac{1}{(4\pi)^{\frac{d}{2}}}\sum_{n=0}^{\infty}\mathbf{B}_n(\Delta)Q_{\frac{d-n}{2}}[f(t)].
\end{aligned} \tag{A.26}$$

This completes the derivation of the master equation (A.20) for heat-kernel computations. Note, that we used the definition of the Q-functionals, given in equation (A.22) and the relation $\int_s s^{-x}\tilde{f}(x) = \frac{1}{\Gamma(x)}\int_z z^{x-1}f(z)$.

When investigating matter fields, such as in chapter 5, we often encounter kinetic operators of the form $\tilde{\Delta} = -\nabla^2\cdot\mathbb{1} + \mathbf{E}$, where $\mathbf{E}$ is a linear map acting on the spacetime and the internal indices of the fields. In this notation, $\mathbb{1}$ has to be understood as the identity

in the respective field space.
If $[\Delta, \mathbf{E}] = 0$[18], we can relate the coefficients of the modified Laplacian $\tilde{\Delta}$ and those of the initially considered operator $-\nabla^2$ via

$$\operatorname{Tr} \mathrm{e}^{-s\left(-\nabla^2+\mathbf{E}\right)} = \frac{1}{(4\pi)^{\frac{d}{2}}} \sum_{k,l=0}^{\infty} \frac{(-1)^\ell}{\ell!} \int_x \sqrt{g}\ \operatorname{Tr} \mathbf{b}_k(\Delta) \mathbf{E}^\ell s^{k+\ell-2}. \tag{A.27}$$

This results in the following, modified values for the coefficients we are interested in:

$$\begin{aligned} \mathbf{b}_0 &= \mathbb{1} \\ \mathbf{b}_2 &= \frac{\mathcal{R}}{6} \cdot \mathbb{1} - \mathbf{E}. \end{aligned} \tag{A.28}$$

For further study and a more general treatment of the modified Laplacians, including higher order coefficients, [5, 17] are recommended.

18. In the case of $[\Delta, \mathbf{E}] \neq 0$, there would be additional terms including (higher order) commutators of Δ and $\mathbf{E}$ due to the Baker-Campbell-Hausdorff formula.

Additional calculations

For the sake of completeness, we present some auxiliary calculations and important steps, that were used to obtain the results presented in the scope of this work, but were in general too long or unsuitable to be included in the main part.

Fermion part

In the part on fermions, we are confronted with the Dirac operator $\not{\nabla} = \gamma^\mu \nabla_\mu$, that involves a contraction with gamma matrices. Here, we present the proof of an identity we used in the derivation of the fermion two-point function:

$$
\begin{aligned}
\not{\nabla}^2 &= \nabla_\mu \nabla_\nu \gamma^\mu \gamma^\nu \\
&= \frac{1}{2} \nabla_\mu \nabla_\nu \left(\gamma^\mu \gamma^\nu + \gamma^\nu \gamma^\mu \right) \\
&= g^{\mu\nu} \nabla_\mu \nabla_\nu \\
&= \nabla^2
\end{aligned}
\tag{B.1}
$$

The third line follows from the definition of the Clifford algebra $\{\gamma^\mu, \gamma^\nu\} = 2g^{\mu\nu}\mathbb{1}$.

We solved the functional trace for the fermions once again using the heat-kernel techniques introduced in appendix A, but in contrast to the other calculations, we could not use the rather simple expressions for the threshold functions evaluated for the Litim cutoff (2.16). We expressed the fermion regulator in terms of the scalar regulator which means, that the Q-functionals looked different and therefore we had to solve these "new" integrals analytically. We start by simplifying the trace expression as much as possible:

$$
\begin{aligned}
-\operatorname{Tr}\left[G_{k,\bar{\psi}\psi}\, \partial_t R_{k,D} \right] &= -\operatorname{Tr}\left[\frac{Z_D \cdot \bar{\Delta}_{(1/2)} \cdot \left(\frac{\partial_t r_{k,S}}{2\sqrt{1+r_{k,S}}} - \eta_D \left(\sqrt{1+r_{k,S}} - 1 \right) \right)}{Z_D \cdot \bar{\Delta}_{(1/2)} \sqrt{1+r_{k,S}}} \mathbb{1}_D \right] \\
&= -N_D \operatorname{Tr}\left[\frac{\theta\,(1-\chi)}{\chi \left(r_{k,S}(\chi) + 1 \right)} + \eta_D \left(\frac{1}{\sqrt{r_{k,S}(\chi) + 1}} - 1 \right) \right],
\end{aligned}
\tag{B.2}
$$

We substituted $\chi := \bar{\Delta}_{(1/2)}/k^2$. Since the trace is linear, we can separate both terms and compute them independently.

First term:

$$-N_D \operatorname{Tr}\left[\frac{\theta(1-\chi)}{\chi\Big(1+r_{k,S}(\chi)\Big)}\right] \overset{(A.20)}{=} -N_D \frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[\operatorname{Tr}\mathbf{b}_0 \int_0^\infty \mathrm{d}\chi\, \chi \frac{\theta(1-\chi)}{\chi\Big(1+r_{k,S}(\chi)\Big)}\right.$$
$$\left. + \operatorname{Tr}\mathbf{b}_2 \int_0^\infty \mathrm{d}\chi\, \frac{\theta(1-\chi)}{\chi\Big(1+r_{k,S}(\chi)\Big)}\right]. \tag{B.3}$$

We directly set $\Gamma(n) = 1$, since n equals 1 or 2 in both cases. With the definition of the Heaviside function, we can set the upper integration bound for the $\mathrm{d}\chi$ integration to 1. We use the heat-kernel coefficients for the spin-$\frac{1}{2}$ Laplacian given in (5.41) and the definition of the Litim shape function (2.16) and obtain:

$$\begin{aligned} -N_D \operatorname{Tr}\left[\frac{\theta(1-\chi)}{\chi\Big(1+r_{k,S}(\chi)\Big)}\right] &= -N_D \frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[4\int_0^1 \mathrm{d}\chi\,\chi + \frac{5}{12}\bar{\mathcal{R}}\int_0^1 \mathrm{d}\chi\, 1\right] \\ &= -N_D \frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[2 + \frac{5}{12}\bar{\mathcal{R}}\right]. \end{aligned} \tag{B.4}$$

Second term:

$$\begin{aligned} -N_D \operatorname{Tr}\left[\eta_D\left(\frac{1}{\sqrt{1+r_{k,S}(\chi)}} - 1\right)\right] &= -N_D \frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[4\eta_D \int_0^1 \mathrm{d}\chi\, \chi(\chi-1)\right. \\ &\qquad \left. + \frac{5\eta_D}{12}\bar{\mathcal{R}}\int_0^1 \mathrm{d}\chi\,(\chi-1)\right] \\ &= -N_D \frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[\eta_D\left(-\frac{2}{3} - \frac{5}{24}\bar{\mathcal{R}}\right)\right]. \end{aligned} \tag{B.5}$$

Altogether we find the following result for the functional trace of the fermion contribution:

$$-\operatorname{Tr}\left[G_{k,\bar{\psi}\psi}\partial_t R_{k,D}\right] = -N_D \frac{1}{(4\pi)^2}\int_x \sqrt{\bar{g}}\left[2\left(1-\frac{\eta_D}{3}\right) + \frac{5}{12}\bar{\mathcal{R}}\left(1-\frac{\eta_D}{2}\right)\right]. \tag{B.6}$$

Gauge field part

During the computation of the gauge field contribution to the running of G and Λ, we encounter the following term, which can be simplified a lot after a few manipulations:

$$
\begin{aligned}
\int g^{\mu\nu} g^{\kappa\lambda} F_{\mu\kappa} F_{\nu\lambda} = \int F_{\mu}{}^{\lambda} F^{\mu}{}_{\lambda} &= \int F_{\mu\lambda} F^{\mu\lambda} \\
&= \int \left(\partial_\mu A_\lambda - \partial_\lambda A_\mu\right) F^{\mu\lambda} + \mathcal{O}\left(A^3\right) \\
&\overset{(\star)}{=} \int 2\partial_\mu A_\lambda F^{\mu\lambda} \\
&= \int 2\partial_\mu A_\lambda \left(\partial^\mu A^\lambda - \partial^\lambda A^\mu\right) \\
&= \int 2\left(\partial_\mu A_\lambda \partial^\mu A^\lambda - \partial_\mu A_\lambda \partial^\lambda A^\mu\right) \\
&\overset{(\dagger)}{=} -\int 2\left(A_\lambda \partial^2 A^\lambda - A_\lambda \partial_\mu \partial^\lambda A^\mu\right) \\
&= \int 2A_\lambda \left[\partial^\mu \partial^\lambda - g^{\mu\lambda}\partial^2\right] A_\mu
\end{aligned}
\tag{B.7}
$$

For the first non-trivial step ($\star$) we use that $2\partial_\mu A_\lambda = \partial_{(\mu} A_{\lambda)} + \partial_{[\mu} A_{\lambda]}$, where $(\cdots)$ and $[\cdots]$ denote symmetrization and antisymmetrization w. r. t. the respective indices. The symmetric part vanishes due to the fact, that $F^{\mu\lambda}$ is antisymmetric under $\mu \rightleftharpoons \lambda$. This allows us to write $2\partial_\mu A_\lambda F^{\mu\lambda} = \partial_{[\mu} A_{\lambda]} F^{\mu\lambda} = \left(\partial_\mu A_\lambda - \partial_\lambda A_\mu\right) F^{\mu\lambda}$. The second non-trivial step (†) results from integrating by parts and assuming vanishing boundary terms.

Later on in the gauge field calculation, after specifying the gauge parameter $\xi = 1$, we encounter a commutator of covariant derivatives acting on A_μ. With the definition of the curvature tensor, given in equation (3.10), we find

$$
\begin{aligned}
\left[\nabla^\mu, \nabla^\lambda\right] A_\mu &= R_{\mu}{}^{\rho\mu\lambda} A_\rho \\
&= R^{\rho\lambda} A_\rho
\end{aligned}
\tag{B.8}
$$

Now, one simply has to rename the dummy indices $\rho \rightleftharpoons \mu$ to find the wanted expression. We encounter the same term also in the computation of the Faddeev-Popov ghosts for the graviton sector.

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List of Figures

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Declaration of Authorship

I hereby certify that this thesis has been composed by me and is based on my own work, unless stated otherwise.

Heidelberg, 8th of July 2019

Mathieu Kaltschmidt