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Abstract

The most common spin foam models of gravity are widely believed to be discrete path integral quantizations of the Plebanski action. However, their derivation in present formulations is incomplete and lower dimensional simplex amplitudes are left open to choice. Since the large-spin behavior of these amplitudes determines the convergence properties of the state-sum, this gap has to be closed before any reliable conclusion about finiteness can be reached. It is shown that these amplitudes are directly related to the path integral measure and can in principle be derived from it. This requires a detailed knowledge of the constraint algebra and the corresponding gauge fixing of its first class part which in the case of gravity generates space-time diffeomorphisms. It has been suggested that the discretization of space-time in a spin foam model breaks the diffeomorphism gauge without introducing an explicit gauge fixing. Here we show that minimal requirements of background independence provide non trivial restrictions on the form of an anomaly free measure which imply the existence of divergent amplitudes in the spin foam model that could be interpreted as due to infinite contribution of gauge equivalent configurations.

1 Introduction

In recent years, spin foam models have been established as possible candidates for a quantum theory of gravity (for recent reviews see [1, 2]). They are commonly viewed as covariant (path integral) versions of a canonical quantization and in fact share some features of quantum geometry (though there is no precise relation yet). As a discretized path integral they can be derived from Plebanski’s action [3] which is a formulation of general relativity as a constrained $BF$-theory [4]. Being path integrals of a gauge theory, they have to deal with the anomaly issue: the path integral measure has to be invariant under transformations generated by the constraints. Sometimes it is claimed that a covariant quantization avoids the issue of anomalies which plagued canonical approaches for a long time (but see [5]). However, it is well known that there is also an anomaly problem in path integral quantizations which in spin foam quantizations has just been ignored in most of the existing literature (see, however, [6] for a recent paper which discusses this issue independently in the example of 3-dimensional $BF$-theory). A complete analysis of this issue would require

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an understanding of the continuum limit which has not yet been developed sufficiently. Still, we will see that it is possible to shed light on the problem and to derive conditions for the amplitudes involved in the definition of a spin foam model. To that end we look at the problem from two perspectives. First, we view the discretized version, which is obtained by fixing a space-time triangulation, as a regularization of the path integral, in particular its measure, which would result in the continuum limit. Second, we analyze the restrictions on the spin foam measure imposed by background independence directly at the discrete level.

To explain the part of the anomaly problem studied here we first recall the situation of standard path integrals: There is a prescription which results in a unique measure (up to a constant factor) which is at least formally invariant. This is usually the obvious measure which integrates over all canonical coordinates with constant weight function. After removing the regulator the measure might not be invariant giving rise to quantum anomalies. In any case, it is necessary to use the invariant measure for the regularized version; otherwise the gauge symmetries are broken explicitly and the results are unphysical. Since the formally invariant measure is obvious in most cases, the standard term 'anomaly' only refers to the second issue, namely whether or not the measure will remain invariant after removing the regulator.

In the case of spin foam models, the situation is more involved. First, the constraint algebra is mixed and not closed which will be seen to lead to an additional function in the measure which has not been taken into account previously. Secondly, the space-time discretization obscures the role of the measure and the meaning of invariance in this context. Therefore, even analyzing the formal invariance of a measure requires new techniques which will be provided in the present paper. The usual anomaly problem, which analyzes the invariance after the regulator is removed, will not be touched here since the continuum limit is not understood. We will, however, see that already a formally invariant measure, which is a necessary prerequisite for an anomaly-free continuum measure, puts strong restrictions on amplitudes in the spin foam model. Our definition of a formally invariant discrete measure is that it must descend from the formally invariant continuum measure along the lines of the spin foam discretization. Provided the required calculations are feasible, this will also fix the formally invariant measure uniquely up to a constant factor, which translates in conditions for the spin foam amplitudes.

An immediate question in the context of anomalies is whether or not a spin foam state sum can be finite. As a consequence of an invariant measure, a standard path integral quantization cannot lead to finite results without gauge fixing if gauge orbits do not have finite volume (which is to be expected for gravity; to avoid confusion we emphasize here that we are mainly concerned with the diffeomorphism constraints, not with an $SO(4)$ or $SL(2, \mathbb{C})$ Gauss constraint). When the path integral is discretized (a spin foam quantization involves the discretization in space-time by choosing a fixed triangulation as well as a discretization of some of the fields), diverging integrals are replaced by infinite sums over a discrete set of points along the gauge orbits. If everything is invariant, i.e. the measure as well as the discretized action, the discretized path integral should be infinite, too, if it is not gauge fixed (one just replaces a diverging integral with a diverging sum). However,
performing a space-time discretization which leaves the action invariant is not obvious [7, 8], and one might expect that this can explain the possibility of finite spin foam models [9, 10, 11]. In fact, active space-time diffeomorphisms are clearly broken once a triangulation is introduced, and invoking the usual equivalence of active and passive diffeomorphisms could indicate a breaking of the gauge in the path integral. However, the very triangulation which breaks active diffeomorphisms also breaks the correspondence between active and passive transformations: active diffeomorphisms can no longer be allowed arbitrarily since not all of them fix the triangulation, but passive transformations affect the values of fields in a fixed point (or associated with cells of the discretization) which can still be performed with complete freedom. The passive picture leads directly to the requirement that the discrete amplitudes have to descend from the anomaly free continuum measure. While the only active diffeomorphisms which are allowed within a given discrete model are those which correspond to symmetries of the triangulation with its labeling (and thus do not imply non-trivial restrictions), conditions for the amplitudes can also be obtained in the active picture by requiring background independence: the discretization serves as a background to define the model, and physical results, e.g. the state sum, must not depend on the choice of background. This provides a second strategy to compute the amplitudes motivated by anomaly freedom.

A simple model for full gravity is 2+1 dimensional $BF$-theory. This theory is equivalent to 2+1 dimensional gravity for non-degenerate triads and thus has the same gauge orbit structure at least on the constraint surface. It turns out, and is commonly accepted, that the spin foam amplitude for 2+1 dimensional $BF$-theory is infinite in accordance with the expectation from path integrals. The discrete symmetries of the simplicial action can be explicitly analyzed [6] and directly linked with the triangulation independence of the spin foam model. An interesting case is 2+1 gravity with cosmological constant $\Lambda$. In this case the action can be written as that of a $Spin(4)$ Chern-Simons theory whose level $k$ is given by $k = 4\pi/\sqrt{\Lambda}$ (see [12] and references therein). A path integral quantization of this theory leads to the Turaev-Viro model defined in terms of a quantum group $SU_q(2)$ for $q$ a root of unity related to the cosmological constant by $q = \exp[2\pi i/(k + 2)]$. Transition amplitudes turn out to be finite. Although this is often interpreted as a consequence of the infrared cut-off introduced by the quantum deformation, from our viewpoint this is a consequence of the compactness of the gauge group $Spin(4)$.

It is not clear however how to generalize this intuition to four dimensions. The gauge properties $BF$-theory in four dimensions are very similar to its 3-dimensional relative. In particular, divergences in the path integral [13] can also be traced back to infinite volume factors coming from the topological gauge symmetry. If we concentrate on the spin foam models for four dimensional gravity that are obtained from an implementation of constraints on the $BF$ amplitudes (such as Reisenberger[14] or the Barrett–Crane models [15, 16]) the topological gauge symmetry is manifestly broken by the implementation of the constraints. As a result, it was debated whether the remnant gauge symmetries would produce diverging spin foam amplitudes or rather contain “finite volume” gauge orbits. Here we show that minimal requirements of background independence imply the existence of divergences and rule out the finite normalizations proposed in the literature [9, 11, 17].
The aim of the present paper is to devise methods for checking a spin foam quantization for formal anomalies. Spin foam models are specified by determining vertex and face amplitudes. While there is general agreement on the vertex (4-simplex) amplitude, which can be viewed as representing the exponentiated action, there are no clear-cut arguments as to which lower dimensional simplexes amplitude should be used\(^1\); in the literature, it is largely regarded as being open to choice, maybe constrained by semi-classical issues. This problem is particularly pressing because the question of whether or not a model is finite hinges on the asymptotic behavior of these amplitudes. In fact, we will see that these lower dimensional simplex amplitudes represent a discretized version of some part of the path integral measure and can be derived from it. Choosing different amplitudes is equivalent to inserting an arbitrary function into the path integral; then it is very easy to get a finite model by introducing a suppression of the measure along the orbits. However, such an anomalous model has to be dismissed as unphysical.

Note that our criterion is formulated from the perspective that the fundamental theory is intrinsically discrete. No matter how the approach to a continuum description is performed—via a limit or as a coarse-grained approximation—gauge degrees of freedom have to be removed which is only possible with an invariant measure. It is sometimes argued that finiteness arises because ultraviolet or infrared divergences are regularized by quantum gravitational effects like a minimal length scale (this does in fact occur in canonical quantizations [20]). From our point of view, however, this is not tenable since the anomaly issue is completely unrelated to ultraviolet or infrared divergences.

In the following section, we will introduce a finite dimensional toy model which illustrates the steps of a spin foam quantization mimicking \(BF\)-theory with additional constraints. In Section 3 we discuss the definition of the (formal) path integral for constrained systems. In Section 4 we revisit the spin foam quantization of \(BF\)-theory in three dimensions to introduce notation and review the gauge analysis of the discrete theory performed in [6]. In Section 5 we discuss the definition of the correct path integral measure for 4-dimensional spin foam models defined as constrained BF state sums. We look at the problem from the passive and active diffeomorphism perspectives. In the first case, we reformulate \(BF\)-theory in a way which makes the relation between the path integral measure and the face amplitude obvious. This provides us with a recipe for computing the large spin behavior of amplitudes in spin foams for gravity discussed in Subsection 5.1.2. In Subsection 5.2.1 we discuss the restrictions imposed on the form of the measure by background independence from the active picture. In the case of the Barrett-Crane model, we show that various normalizations proposed in the literature do not satisfy these requirements and should be regarded as anomalous. This includes the finite normalization introduced in [19]. In Subsection 5.2.2 we define a simple model satisfying those anomaly freeness requirements. The latter can be thought of as the spin foam quantization of the Husain–Kuchar model.

\(^1\)In the case of the Barrett–Crane model the normalization that yields finite amplitudes is naturally selected in the context of the group field theory (GFT) formulation [18, 19]; however, no clear connection with the formal path integral has not been found yet.
2 A toy model

To illustrate the importance of choosing the correct measure in spin foam models we first discuss a simple toy model with a finite number of degrees of freedom. It incorporates the essential steps of a spin foam quantization of Plebanski’s action for gravity, which are a field discretization and the solution of a constraint for Lagrange multipliers. Being a system with a finite number of degrees of freedom, the continuum limit cannot be modeled. However, as discussed before, the anomaly issue already requires the correct treatment of the regularization before the continuum limit is taken, which will be illustrated here.

2.1 Definition and evaluation

The action of the model is given by

\[ S = \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + \lambda_1 q_1 + \lambda_2 q_2 + \xi (\lambda_1 - \lambda_2)) \]  

which has two constrained degrees of freedom \((q_1, q_2)\), which we assume to live on a circle, with conjugate momenta \((p_1, p_2)\) and three Lagrange multipliers \(\lambda_1, \lambda_2\) and \(\xi\). Compared with BF-theory, \((p_1, p_2, \lambda_1, \lambda_2)\) represents the components of the field \(B\) which contains both physical degrees of freedom and Lagrange multipliers. If we set \(\xi = 0\) resulting in the action

\[ S|_{\xi=0} = \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + \lambda_1 q_1 + \lambda_2 q_2), \]

the theory is constrained completely, i.e., both \(q_1\) and \(q_2\) must be zero. There are no degrees of freedom in this case. With unrestricted \(\xi\), however, the two original Lagrange multipliers are constrained which restores one degree of freedom: \(\lambda_1\) has to equal \(\lambda_2\) and thus only \(q_1 + q_2\) has to be zero whereas the difference is free, which can easily be seen by solving the \(\xi\)-constraint explicitly:

\[ S = \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + \lambda_1 (q_1 + q_2)). \]

This feature mimics the transition from BF-theory to gravity where also additional constraints (the simplicity constraints) reduce the freedom of original Lagrange multipliers of BF-theory and thereby introduce local degrees of freedom.

A spin foam quantization proceeds by quantizing the simple theory whose discretized state sum can be computed explicitly and incorporating the additional constraints at the state sum level. The simple theory (the analog of BF-theory) here is \(S|_{\xi=0}\) with path integral

\[ Z_0 = \int \mathcal{D}^2 q \mathcal{D}^2 p \mathcal{D}^2 \lambda \exp(i S|_{\xi=0}) = \int \mathcal{D}^2 q \mathcal{D}^2 p \exp (i \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2)) \delta(q_1) \delta(q_2) = \int \mathcal{D}^2 p \]  

(2)
with \( D^2 q := D q_1 D q_2 \). The result is certainly infinite since we are dealing with an unfixed gauge theory. In this case a gauge fixing is simple, but we do not do this because we want to understand the role of a field discretization and the multiplier constraints in this respect.

Therefore, we now discretize \( \lambda_1 \) and \( \lambda_2 \) which are analogous to \( B \)-field components (we could also discretize the remaining components \( p_1 \) and \( p_2 \), without changing our results). In analogy to a spin foam quantization we do this by writing the integral representation of the delta function

\[
\delta(q_1) = (2\pi)^{-1} \int d\lambda_1 \exp(i\lambda_1 q_1)
\]
as a sum

\[
\delta(q_1) = (2\pi)^{-1} \sum_{n_1} \exp(in_1 q_1).
\]
The path integral (ignoring constant factors which can be absorbed in the measure) then becomes

\[
Z_0 = \int D^2 q D^2 p \sum_{\{n_1\}, \{n_2\}} \exp(i \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + n_1 q_1 + n_2 q_2))
\]
where the summation index \( \{n\} \) indicates that \( n \) is not a single number but a function of time, and we are summing over the values at fixed times individually (i.e., this is a discrete analog of the path integral).

We will later use this integral to incorporate the additional constraint with multiplier \( \xi \). But first we compute the path integral for \( S \), the analog of gravity, which in this case can also be obtained explicitly:

\[
Z = \int D^2 q D^2 p D^2 \lambda D^2 \xi \exp(iS) = \int D^2 q D^2 p D \lambda_1 \exp(i \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + \lambda_1 q_1 + q_2))
\]
\[
= \int D^2 q D^2 p \exp(i \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2)) \delta(q_1 + q_2) = \int Dq_1 D\Delta p \exp(i \int dt \dot{q}_1 \Delta p) \int Dp'
\]
with \( \Delta p := p_1 - p_2 \) and \( p' = p_1 + p_2 \). Computing the remaining integrations we obtain

\[
Z = \int Dq_1 \delta(\dot{q}_1) \int Dp' = \delta(q_1^{(0)} - q_1^{(1)}) \int Dp'.
\]
Here, \( q_1^{(0)} \) and \( q_1^{(1)} \) represent the initial and the final value of \( q_1 \) which are constrained to be equal but free otherwise. Due to the fact that we have only one remaining gauge symmetry after incorporating the \( \xi \)-constraint, we only have one infinite integral left rather than two in (2).

In finite spin foam models one solves the multiplier constraint \( \lambda_1 - \lambda_2 = 0 \) at the discretized level and, in some cases, obtains a finite result even without fixing the remaining gauge freedom [21]. A spin foam quantization, however, also involves a discretization of space-time which, as already mentioned, is not realized in this finite dimensional model. Still, it is worth checking what effects a field discretization itself can have; effects of the space-time discretization will be discussed later. To do this in our toy model we start from
(3) which contains the integers \( n_1, n_2 \) discretizing the multipliers \( \lambda_1 \) and \( \lambda_2 \). Translating the \( \xi \)-constraint to the discrete level implies \( n_1 = n_2 \) and we must only sum over those pairs of integers fulfilling this condition in order to obtain a quantization for \( S \) with \( \xi \) free (this is analogous to summing only over simple representations in a spin foam quantization):

\[
Z = \int \mathcal{D}^2 q \mathcal{D}^2 p \sum_{\{n_1\}} \exp \left( i \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + n_1 (q_1 + q_2)) \right).
\]

(5)

The result, of course, is the same as in the calculation with continuous \( \lambda_1 \) and, in particular, it is infinite. As anticipated, the field discretization and the spin foam like quantization could not take care of the gauge orbit divergence.

2.2 Modifying the measure

In spin foam quantizations the issue of convergence hinges on the choice of lower dimensional simplex amplitudes, which can be considered as functions of some components of the discretized \( B \)-field. In our model, however, we do not have any free function available since the path integral result is unique. As we will discuss later, the lower dimensional simplex amplitudes of spin foams also are fixed uniquely (up to different discretization choices), but have not been determined yet. To include such a function we write our result in the spin foam form (there is still a \( p \)-integral because we chose not to discretize \( p \))

\[
Z = \int \mathcal{D}^2 p \sum_{\{n_1\}} V(p_1, p_2, n_1)
\]

where

\[
V(p_1, p_2, n_1) := \int \mathcal{D}^2 q \exp \left( i \int dt (\dot{q}_1 p_1 + \dot{q}_2 p_2 + n_1 (q_1 + q_2)) \right)
\]

is the vertex amplitude (analogous to the integration over connections of the discretized \( e^{iS} \)). A model of lower dimensional amplitudes can now be included by simply inserting a new function \( A(n_1) \) into \( Z \) (more generally, \( A \) could also depend on \( p \)):

\[
Z = \int \mathcal{D}^2 p \sum_{\{n_1\}} A(n_1) V(p_1, p_2, n_1).
\]

Our derivation shows that the face amplitude \( A(n_1) \) is fixed and identical to one (or any other non-zero constant), but let us see what a different function would imply. For illustrative purposes, we choose

\[
A(n_1) = (2n_1 + 1)^{-2}
\]

which is finite for all integer \( n_1 \). Now it is easy to see that

\[
Z' = \int \mathcal{D}^2 p \sum_{\{n_1\}} (2n_1 + 1)^{-2} V(p_1, p_2, n_1)
\]
\[
= \int D^2q D^2p \sum_{\{n\}} (2n_1 + 1)^{-2} \exp \left( i \int dt (\dot{q}_1p_1 + \dot{q}_2p_2 + n_1(q_1 + q_2)) \right)
\]
\[
= \int D^2q D^2p \exp \left( i \int dt (\dot{q}_1p_1 + \dot{q}_2p_2 + V_{\text{eff}}(q_1, q_2)) \right)
\]
is finite, where we have the effective potential
\[
V_{\text{eff}}(q_1, q_2) = \log(q_1 + q_2 - \pi) - \frac{1}{2}(q_1 + q_2).
\]
(We used the Fourier series \( \sum_k (2k + 1)^{-2} e^{ik\phi} = -\frac{\pi}{4} (\phi - \pi) e^{-i\phi/2} \) for \( 0 < x < 2\pi \) and extended with \( 2\pi \)-periodicity.) In fact, this is an ordinary path integral for a system of two degrees of freedom in an effective potential \( V_{\text{eff}} \) without constraints. We now have to decide if this finite result makes sense and can tell us anything about the original system. The answer is clearly negative: The role of the effective potential is completely unclear, and it has nothing to do with the original system. Originally, \( q_1 \) and \( -q_2 \) have to equal each other but are free otherwise, whereas in the modified system they are independent but subject to motion in a potential. Furthermore, the kind of modification, e.g. the form of the potential, depends on the face amplitude which has no distinguished form other than \( A(n_1) = 1 \) which follows from the invariant measure. In conclusion, a finite path integral for an unfixed system with constraints cannot be trusted. (It cannot even be regarded as an approximation since the measure is not just a smeared version of a \( \delta \)-function with support on the constraint surface. The effective potential is singular on a submanifold of the configuration space, but this does not happen at the constraint surface \( q_1 + q_2 = 0 \), but at \( q_1 + q_2 = \pi \).) In fact, introducing a non-constant amplitude is nothing but introducing an arbitrary function \( A(\lambda_1) \) into the path integral which breaks the invariance of the measure. (Note that \( \lambda_1 \) serves as a Lagrange multiplier and thus its conjugate momentum \( p_{\lambda_1} \) is implicitly constrained to be zero. The gauge freedom generated by this constraint is broken by introducing an arbitrary function of \( \lambda_1 \) into the measure. Consequently, the multiplier \( \lambda_1 \) is no longer completely free which also affects the remaining gauge freedom.) This explains why we get a finite result with independent \( q_1, q_2 \); and it also demonstrates that here a finite model is anomalous. As discussed in the Introduction, the space-time discretization, which is not modeled here, presents a possible rescue for finite spin foam models. To check this, we need more general methods which will be introduced in what follows.

### 3 General discussion

Since our model incorporated some of the essential steps of a spin foam quantization of gravity, it suggests that the same conclusions regarding the choice of amplitudes hold true in this more complicated case. In this section we discuss the continuous path integral and the correct measure in the presence of second class constraints, which will be necessary to derive the anomaly-free amplitudes.

The characteristic feature of the gravitational action which is commonly used for a spin foam quantization is the presence of a constraint which restricts the allowed values
of Lagrange multipliers appearing in a simpler action. We illustrated this property in the previous toy model where the importance of an invariant path integral measure has been seen explicitly. To find the correct measure it is not sufficient to work solely in a Lagrangian formulation; in particular it is essential to understand the structure of the constraint algebra which can only be achieved in a Hamiltonian analysis. The constraint algebra in this context is always mixed (i.e. neither purely first class nor purely second class) and rather complicated. There are always the usual diffeomorphism constraints of gravity which must form a suitable first class sub-algebra, but in this particular formulation there is also a second class contribution: a constraint which restricts the multipliers of other constraints must be second class. Despite first appearance, even in the toy model the additional constraint is second class. Although the constraints \( C_1 = q_1, C_2 = q_2 \) and \( C_3 = \lambda_1 - \lambda_2 \) Poisson commute, one has to take into account that in this form they are constraints on a non-symplectic Poisson manifold with coordinates \((q_1, p_1; q_2, p_2; \lambda_1, \lambda_2)\) where the standard definitions of Dirac’s classification do not apply (see [22] for a discussion and generalized definitions). One can easily introduce an equivalent constrained system which has constraints on a symplectic manifold by adding the momenta \( \pi_1 \) and \( \pi_2 \) conjugate to the restricted multipliers \( \lambda_1 \) and \( \lambda_2 \), together with the constraints \( C_4 = \pi_1, C_5 = \pi_2 \). The constraints \( C_I, I = 1, \ldots, 5 \) are then defined on a symplectic manifold and now it is obvious that \( C_3 \) does not commute with all constraints. In fact \( C_3 \) and \( C_4 - C_5 \) form a second class sub-algebra, \( \{C_3, C_4 - C_5\} = 2 \), whereas \( C_1, C_2 \) and \( C_4 + C_5 \) are first class.

The presence of second class constraints requires a special treatment when deriving the correct measure. It is not sufficient simply to include an integration over the multipliers since this leaves open an arbitrary function. In the absence of constraints the invariant path integral measure is given by the determinant of the symplectic form which leads to \( \mathcal{D}q\mathcal{D}p \) for canonical coordinates \((q, p)\). A similar treatment is not possible for the multiplier integration since multipliers form a Lagrangian sub-manifold of the extended phase space such that the determinant of their symplectic structure would be zero. The measure is well-defined after solving the second class constraints (and turning the first class constraints into second class ones by fixing the gauge), which leads to the symplectic structure following from the Dirac bracket. For completeness, we will next show how to derive the correct treatment of the multiplier measure by requiring that after solving the constraints in the integral we obtain the determinant of the Dirac symplectic structure [23].

We start with a system with \( 2n \) coordinates \( x_i, i = 1, \ldots, 2n \) on a symplectic phase space \((M, \omega)\) and \( m \) second class constraints \( C_I, I = 1, \ldots, m \). (There might be additional, first class constraints which are not relevant for this section. They can either be gauge fixed and included in the constraints \( C_I \) or be left for later treatment, e.g. factoring out the volume of their gauge orbits. The second possibility is particularly interesting here since a gauge fixing is sometimes claimed to be unnecessary for spin foam models of gravity.) The path integral (where the constraint part has been split off the action \( S = S_0 + \int \sum I \lambda^I C_I \)) then is

\[
Z = \int \mathcal{D}^{2n}x \sqrt{\det \omega} \mathcal{D}^m \lambda \mu(x) \exp \left( i \int \sum I \lambda^I C_I \right) \exp(iS_0)
\]
with a function $\mu(x)$ for the multiplier measure which will be determined shortly. This function must not depend on the $\lambda^I$ because otherwise the multiplier integration would not yield $\delta$-functions of the constraints. With a $\lambda$-independent $\mu$ we can perform the $\lambda$ integrations explicitly and obtain

$$Z = \int \mathcal{D}^{2n}x \sqrt{\det \omega} \mu(x) \prod_I \delta(C_I) \exp(iS_0).$$

To proceed further we transform from the coordinates $x_i$, $i = 1, \ldots, 2n$ to coordinates $(y_\alpha; C_I)$ with $\alpha = 1, \ldots, 2n - m$, $I = 1, \ldots, m$ (assuming that the constraints are regular and irreducible such that they can be used as local coordinates on $M$). To find the Jacobian of this transformation we use the fact that locally the symplectic manifold $(M, \omega)$ can be represented as $(M, \omega) \cong (R, \omega_D) \times_R (P, \Pi^{-1}_P)$ using the following notation. The symplectic manifold $(R, \omega_D)$ is the constraint surface $R \subset M$ defined by $C_I = 0$, $I = 1, \ldots, m$, endowed with the Dirac symplectic structure $\omega_D$. The manifold $P$ is given by the image of a neighborhood of a point in $R$ under the functions $C_I: M \to \mathbb{R}$ (i.e. $P$ is a neighborhood of $0$ in $\mathbb{R}^m$; for our purposes it is sufficient to know $P$ only locally) and coordinatized by $(C_I)$, $I = 1, \ldots, m$. If the constraint algebra is closed, $^2$ $P$ can be defined globally and is a Poisson manifold with Poisson tensor $\Pi_P$ defined by $\Pi_P (dC_I, dC_J) := \{C_I, C_J\}$ where the bracket on the right hand side is computed using the symplectic structure $\omega$ on $M$ [24]. For second class constraints the inverse of $\Pi_P$ exists and $(P, \Pi^{-1}_P)$ is a symplectic manifold. If the constraint algebra is not closed, the Poisson tensor $\Pi_P$ depends not only on the coordinates $C_I$ of $P$, but also on the coordinates of $R$ such that $(P, \Pi_P)$ as a symplectic manifold depends on the point in $R$ chosen for its definition. The right component of the product decomposition of $M$ then depends on a point in the left component, which is indicated by the subscript $R$ of the symbol $\times$. That the decomposition is valid locally can be shown using the methods of [22] where it has been proven for a closed algebra.

Here we use this local decomposition to factor the original symplectic structure $\omega = \omega_D \otimes \Pi^{-1}_P$ which allows us to perform the coordinate transformation in the last integral,

$$Z = \int \mathcal{D}^{2n-m}y \mathcal{D}^mC \sqrt{\det \omega_D/\det(\{C_I, C_J\})} \mu(y, C) \prod_I \delta(C_I) \exp(iS_0)$$

$$= \int \mathcal{D}^{2n-m}y \sqrt{\det \omega_D} \mu(y, 0)/\sqrt{\det(\{C_I, C_J\})} \exp(iS_0).$$

Since $y_\alpha$ are coordinates on the constraint surface, their path integral measure must be given by the Dirac symplectic structure which already appears in the last integral. If there is any other non-constant function besides the exponential of the action, the measure will not be invariant and the path integral will be anomalous. Therefore, the function $\mu(x)$ has to be $\sqrt{\det(\{C_I, C_J\})}$ which fixes the free function in the original integral (strictly speaking, $\mu$ is only fixed on the constraint surface; values away from the surface will not

\[i.e., \text{the right hand side consists of functions of the constraints and thus is constant on the constraint surface}\]
affect the integral). In the presence of second class constraints, therefore, the path integral to start with is

\[ Z = \int \mathcal{D}^{2n}x \mathcal{D}^m \lambda \sqrt{\det \omega} \sqrt{\det \{C_I, C_J\}} \exp \left( i \int \sum_I \lambda^I C_I \right) \exp (iS_0) \]  

which requires a detailed knowledge of the constraint algebra. Note that the determinant of the constraint brackets also appears in this form when first class constraints are gauge fixed a la Faddeev–Popov, where the other half of the second class constraints are gauge fixing conditions \( f_\beta \). In this case the measure contains a function \( \det (\partial f_\alpha / \partial \delta |_{\delta=0}) \) where \( f_\alpha \) is a gauge fixing condition for the first class constraint \( D_\alpha \) and \( \delta_\alpha \) its gauge parameter. Thus, \( \partial f_\alpha / \partial \delta |_{\delta=0} = \{C_{m/2+\alpha}, C_\beta\}|_{\delta=0} \) (we assume that the second class constraints are arranged in such a way that the first \( m/2 \) are the original first class constraints, \( C_I = D_I \) for \( 1 \leq I \leq m/2 \), and the rest are the gauge fixing conditions \( C_I = f_I \) for \( m/2 < I \leq m \)) and

\[ \det (\partial f_I / \partial \delta |_{\delta=0}) = \det \{C_I, C_J\}|_{C=0; 1 \leq I \leq m/2; m/2 < I \leq m} = \sqrt{\det \{C_I, C_J\}}|_{C=0} \cdot \]  

In existing spin foam quantizations the correct factor for the multiplier integration has not been taken care of; instead any multiplier has been associated simply with a measure \( \mathcal{D} \lambda \) without justification. While this simplifies the analysis and avoids a discussion of the constraint algebra, in general it introduces anomalies and is not permissible; it also needs to be included if some second class subalgebra is treated before solving first class constraints (without discussing gauge fixing). Note that the additional factor can be ignored when it is constant on the constraint surface, which is always the case for a closed constraint algebra. In particular, our toy model has a closed algebra and so our treatment was correct even though we ignored the additional factor in the measure. For more complicated systems including gravity in the Plebanski formulation, however, this is not expected to be the case.

In this section we have recalled the correct choice of measure in the continuum path integral in the presence of second class constraints. In the remainder we will see how it can be built into the discrete spin foam version and how it affects the amplitudes.

4 BF-theory

BF-theory is important in what follows because it appears as an intermediate step in the definition of the gravitational spin foam model. In this section we use it also to illustrate the possible role played by the space-time discretization in the context of anomalies.

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3In a covariant formulation one usually chooses a gauge fixing functional depending on all components of the fields, which would require the use of the extended phase space in a canonical picture. For pure first class constraints it is expected that a covariant gauge fixing would be better suited to the spin foam approach; an example can be found in [6] and in Section 4.
The action of \( BF \)-theory is given by

\[
S(B, A) = \int_{\mathcal{M}} \text{Tr}(B \wedge F(A)),
\]

where the field \( A \) corresponds to a connection on a principal bundle with compact structure group \( G \) (which will later be taken to be \( SU(2) \), which gives 3-dimensional Riemannian gravity) and the field \( B \) is a Lie algebra valued 1-form. The local symmetries of the action correspond to the internal gauge transformations

\[
\delta B = [B, \omega], \quad \delta A = d_A \omega,
\]

for \( \omega \) a Lie algebra valued scalar field where \( d_A \) denotes the covariant exterior derivative, and ‘triad translations’

\[
\delta B = d_A \eta, \quad \delta A = 0,
\]

where \( \eta \) is a Lie algebra valued function. The first invariance is manifest from the form of the action, while the second is a consequence of the Bianchi identity, \( d_A F = 0 \). If one writes the theory in the Hamiltonian formulation, one observes that the previous symmetries are gauge symmetries in the Dirac sense, i.e., they are generated by the Poisson bracket with the corresponding first class constraints; there are no second class constraints.

Moreover the number of constraints equals the number of configuration variables of the phase space, which implies the theory can only have global degrees of freedom. This can be checked directly by writing down the equations of motion

\[
F(A) = d_A A = 0, \quad d_A B = 0.
\]

The first equation is solved by flat connections which are locally gauge. The solutions of the second equation are also locally gauge, once the flatness condition \( (F(A) = 0) \) holds, as any closed form is locally exact.\( ^4 \)

### 4.1 Derivation of the spin foam model

To fix our notation, we will now discuss the spin foam quantization of three-dimensional gravity, where \( \mathcal{M} \) is a three-dimensional manifold and \( G = SU(2) \), and later mention necessary changes for four dimensions. The quantization of \( BF \)-theory is done by replacing the manifold \( \mathcal{M} \) with an arbitrary cellular decomposition \( \mathcal{C} \). We also need the notion of the associated dual 2-complex of \( \Delta \) denoted by \( \mathcal{J}_\Delta \). The dual 2-complex \( \mathcal{J}_\Delta \) is a combinatorial object defined by a set of vertices \( v \in \mathcal{J}_\Delta \) (dual to 3-cells in \( \Delta \)) edges \( e \in \mathcal{J}_\Delta \) (dual to 2-cells in \( \Delta \)) and faces \( f \in \mathcal{J}_\Delta \) (dual to 1-cells in \( \Delta \)).

\( ^4 \) As is well known, one can easily check that the infinitesimal diffeomorphism gauge action \( \delta B = \mathcal{L}_v B \), and \( \delta A = \mathcal{L}_v A \), where \( \mathcal{L}_v \) is the Lie derivative in the \( v \) direction, is a combination of (10) and (11) for \( \omega = v^a A_a \) and \( \eta_b = v^a B_{ab} \), respectively, acting on the space of solutions, i.e. when (12) holds.
The fields $B$ and $A$ have support on these discrete structures. The $su(2)$-valued 1-form field $B$ is represented by the assignment of a $B \in su(2)$ to each 1-cell in $\Delta$. The connection field $A$ is represented by the assignment of group elements $g_e \in SU(2)$ to each edge in $J_{\Delta}$.

The action of the simplicial theory is given by

$$S = \sum_{f \in J_{\Delta}} B_{\ell_f} U_f,$$

where $B_{\ell_f}$ is the Lie algebra element associated to the 1-simplex $\ell_f \in \Delta$, dual to the face $f \in J_{\Delta}$, and $U_f = g_e^1 g_e^2 \cdots g_e^N$ is the discrete holonomy around $f \in J_{\Delta}$. Since $\ell_f$ is in one-to-one correspondence with $f \in J_{\Delta}$, from now on we denote $B_{\ell_f}$ simply as $B_f$.

The partition function is defined as

$$Z(\Delta) = \int \prod_{f \in J_{\Delta}} dB_f \prod_{e \in J_{\Delta}} dg_e \, e^{Tr[B_f U_f]},$$

where now $dB_f$ is the regular Lebesgue measure on $su(2) \cong \mathbb{R}^3$, $dg_e$ corresponds to the invariant measure on $SU(2)$.

Integrating over $B_f$, we obtain

$$Z(\Delta) = \int \prod_{e \in J_{\Delta}} dg_e \prod_{f \in J_{\Delta}} \delta(g_e^1 \cdots g_e^N),$$

where $\delta$ corresponds to the delta distribution defined on $L^2(SU(2))$.

The integration over the discrete connection $(\prod_e dg_e)$ can be performed if one expands the delta function in the previous equation using harmonic analysis on the group. In the case of compact groups this is known as Peter–Weyl theorem, which asserts that any function $f \in L^2(SU(2))$ can be written as a sum over matrix elements of unitary irreducible representations of $SU(2)$. Using the Peter–Weyl decomposition, the $\delta$-distribution becomes

$$\delta(g) = \sum_{j \in \text{irrep}(SU(2))} \Delta_j \, \text{Tr}[\rho_j(g)],$$

where $\Delta_j$ denotes the dimension of the unitary representation $j$, and $\rho_j(g)$ is the corresponding representation matrix. Using equation (16), the partition function (15) becomes

$$Z(\Delta) = \sum_{C_f: \{f\} \to \{j\}} \int \prod_{e \in J_{\Delta}} dg_e \prod_{f \in J_{\Delta}} \Delta_j \, \text{Tr}[\rho_{j_f}(g_e^1 \cdots g_e^N)],$$

where $C_f: \{f\} \to \{j\}$ represents the assignment of irreducible representations to faces in the dual 2-complex $J_{\Delta}$. Each particular assignment is referred to as a coloring. The summation is then over colored 2-complexes (spin foams).

If the $SU(2)$ group element $g_e$ corresponds to an $n$-valent edge $e \in J_{\Delta}$, i.e., an edge bounding $n$ faces, there are $n$ representation matrices evaluated on $g_e$ in (15). The relevant integral is

$$\int dg \, \rho_{j_1}(g) \otimes \rho_{j_2}(g) \otimes \cdots \otimes \rho_{j_n}(g) = \sum_l C_{j_1 j_2 \cdots j_n}^{l_1 l_2 \cdots l_n} C_{j_1 j_2 \cdots j_n}^{l_1 l_2 \cdots l_n},$$

13
i.e., the projector onto $\text{Inv} [\rho_{j_1} \otimes \rho_{j_2} \otimes \cdots \otimes \rho_{j_n}]$. On the RHS we have chosen an orthonormal basis of invariant vectors (intertwiners) to express the projector. Integrating over the connection (15) becomes

$$Z(\Delta) = \sum_{c_f : \{f \} \to \{j\}} \sum_{c_v : \{e\} \to \{v\}} \prod_{f \in J_\Delta} \Delta_{j_f} \prod_{v \in J_\Delta} A_v(t_v, j_v),$$  

(19)

where $A_v(t_v, j_v)$ is given by the appropriate trace of the intertwiners $t_v$ corresponding to the edges bounded by the vertex and $j_v$ are the corresponding representations. This amplitude is given in terms of $SU(2)$ $3Nj$-symbols. When $\Delta$ is a simplicial complex then all the edges in $J_\Delta$ are 3-valent and vertices are 4-valent. Consequently, there are 3 representation matrices for all edges in (18) and the corresponding amplitude is given by the contraction of the corresponding four 3-valent intertwiners, i.e., a $6j$-symbol. In that case the partition function takes the familiar Ponzano–Regge form

$$Z(\Delta) = \sum_{c_f : \{f \} \to \{j\}} \prod_{f \in J_\Delta} \Delta_{j_f} \prod_{v \in J_\Delta} A_v(t_v, j_v),$$  

(20)

where the tetrahedron corresponds to the graphical representation of the $6j$-symbol.

In the next section we will analyze the case of gravity in four dimensions. The models of interest are defined in terms of constrained 4-dimensional $BF$-theory which for Euclidean signature has the gauge group $SO(4) \cong SU(2) \times SU(2)$. The discretization of $BF$-theory in four dimensions is analogous to that of 3-dimensional $BF$-theory. The main difference is that the $B$ field is now a Lie algebra valued 2-form and so is discretized by the assignments of Lie algebra elements $\{B_{t_f}\}$ to the 2-dimensional surfaces defined by the triangles $t_f \in \Delta$. Triangles are in one-to-one correspondence with faces $f \in J_\Delta$. The connection is discretized in precisely the same way as in three dimensions; namely, by assigning group elements $g_e$ to edges $e \in J_\Delta$. Upon integration over $\{B_{t_f}\}$ and the $\{g_e\}$ the amplitudes can be expressed as a spin foam sum similar to (20), where the $6j$-symbol is replaced by a $15j$-symbol represented by a 4-simplex (see [1] for details and references). Since any $SO(4)$-representation can be decomposed into a product of to $SU(2)$-representations, the face amplitude is now $\Delta_{\{j_1, j_2\}} = (2j_1 + 1)(2j_2 + 1)$.

### 4.2 Gauge fixing for 3d gravity

The expression (19) is generically divergent. A reason for this divergence is the (non-compact) gauge freedom (11) [6]. That implies that some of the $B$ integrations in (14) –or equivalently some of the $\delta$-functions in (15)– are redundant. The explicit form of the discrete gauge symmetry of simplicial 3-dimensional gravity is described in detail in [6].
In addition to the standard $SU(2)$ gauge invariance of the action corresponding to (10), there is a discrete analog of (11). Namely, the action (13) is invariant under the transformation

$$\delta B_{\ell_f} = \eta_v - [\Omega^v_{\ell_f}, \eta_v] \quad \text{if} \quad v \subset \ell_f$$
$$\delta B_{\ell_f} = 0 \quad \text{if} \quad v \not\subset \ell_f$$

(21)

where $\eta_v$ is a Lie algebra element associated to the vertex $v \in \Delta$ and $\Omega^v_{\ell_f}$ is also in the Lie algebra and can be explicitly given in terms of the logarithm of the elements $\{U_f\}$ for $f' \neq f$ and contained in the set of faces that form the dual bubble in $\mathcal{J}_\Delta$ around the vertex $v$ [6]. The previous transformation is a symmetry of the action due to the discrete version of the Bianchi identity stating that the (ordered) product of $U_f$ around a bubble is equal to the identity.

Assuming $\Delta$ is path connected we can set $B = 0$ along a contractible (within $\Delta$) path $L$ containing all vertices in $\Delta$ using the gauge freedom (21). This fixes this gauge freedom completely. As shown in [6] this gauge fixing contributes to the measure with the Fadeev–Popov determinant of the type appearing in (8)

$$\det(\partial f_I / \partial \delta_J |_{\delta=0}) = (1 + |\Omega^v_{\ell_f}|^2)$$

(22)

which only depends on the connection $\{g_v\}$. Integration over the $B$ field produces the curvature delta functions as in (15) which in turn imply $\Omega^v_{\ell_f} = 0$ and hence trivial Fadeev–Popov factors.

The effect is very simple: we have to drop out all the $\delta$-functions in (15) corresponding to faces $f \in \mathcal{J}_\Delta$ dual to the 1-simplices in $L$ which are redundant leading to divergences in the non gauge fixed formulation. The gauge fixing is in this way analogous to changing the discretization. In this precise sense we find a connection between discretization independence of the partition function and the gauge freedom (21).

5 Spin foam measure

As we have seen in the previous section, the well-known amplitudes of the spin foam model of $BF$-theory can be seen as emerging from a concrete transition from the continuum measure of a path integral. In gravity models which are defined as constrained $BF$-theories (of the type of Barrett and Crane’s), this direct construction of the measure is not available. This is so because the simplicity constraints are imposed on $BF$ amplitudes after the integration over the (discrete) $B$-field has been performed. At this stage, one is already dealing with a discrete state sum—where $B$ configurations are replaced by irreducible unitary representations—and the connection with the formal continuous measure is lost.

In this section, we will first present a method to impose the simplicity constraints (or, more generally, second class constraints in a spin foam quantization) in such a way that there is always a clear connection to the continuum measure. It is based on a passive
interpretation where the (diffeomorphism) gauge transformations do not change the discretization, but rather values of the fields. Explicit calculations in the case of gravity, however, so far look complicated and we will not pursue a calculation of the amplitudes here. Nevertheless, one can hope that at least their asymptotic behavior for large labels can be found easily, which would allow us to see whether or not the state sum will be finite.

Independently, we can use the active picture where diffeomorphisms change the discretization. There are now moves which lead to a different discretization but a configuration which has to be considered as physically equivalent to the original one. Anomaly freeness, or background independence, then requires that the amplitudes do not change under those moves, which as we will see in Subsection 5.2.1 imposes non trivial restrictions on the measure. In Subsection 5.2.2 we present a toy model of an anomaly free spin foam.

5.1 Passive picture

As recalled in Section 4.1, in the standard derivation of the spin foam models for $BF$-theory there is a clear-cut relationship between the path integral measure and spin foam amplitudes. This relationship comes directly from the result of integrating the $B$-field in (14) which results in (15) and the use of the Peter–Weyl decomposition of the $\delta$-distribution on the group (harmonic analysis on the group). We can relate the $B$ configurations with the spin labels appearing in the state sum in a way that can be useful for exploring the definition of the spin foam measure in the more complicated case of gravity. In the next subsection we will re-derive part of the spin foam quantization of $BF$-theory with different methods which also show how the labels $j$ arise from a discretization of $B$. This will allow us to propose a recipe for the construction of the correct measure in the case of gravity in Subsection 5.1.2. The implementation of this general construction requires the knowledge of the constraint structure of the discrete theory that is not available at present.

5.1.1 $BF$-theory in polar coordinates

As recalled in Section 4.1, a main ingredient of the spin foam quantization of $BF$-theory in three or four dimensions is the formula

$$(2\pi)^{-1} \int dB \exp(2i \text{tr}(Bg)) = \delta(g) = \sum_j (2j + 1) \text{tr} \rho_j(g)$$

(23)

where $g \in SU(2)$, $B \in su(2)$, which follows from the Peter–Weyl theorem. By using this formula, the continuous values of the field $B$ are replaced by discrete values $j$, but the exact correspondence remains unclear. In particular, there are three independent components in $B$, but only one discrete label $j$. The following calculations will now show that $j$ is the discretized radial component of $B$ in polar coordinates whereas the angular components are integrated out.

To this end we write the $su(2)$ element

$$B = r n^i \tau_i \quad \text{with} \quad n = (\sin \vartheta \cos \varphi, \sin \vartheta \sin \varphi, \cos \vartheta)$$

16
in polar coordinates \((r, \theta, \varphi)\) where \(\tau_j = i/2\sigma_j\) with Pauli matrices \(\sigma_j\). To simplify the calculation we also choose the \(SU(2)\) element in the gauge \(g = \exp c\tau_3 = \cos \frac{c}{2} + 2\tau_3 \sin \frac{c}{2}\) without loss of generality (thanks to gauge invariance of the trace in \((23)\)). This leads to

\[
\delta(g) = (2\pi)^{-1} \int dr d\theta d\varphi r^2 \sin \theta \exp(-2i r \cos \theta \sin \frac{c}{2})
\]

\[
= \int dr \frac{r}{2\sin \frac{c}{2}} \left( \exp(-2i r \sin \frac{c}{2}) - \exp(2i r \sin \frac{c}{2}) \right)
\]

whereas the discrete form is

\[
\delta(g) = \sum_j (2j + 1) \frac{\sin(j + \frac{1}{2}) c}{\sin \frac{c}{2}} = \sum_j \frac{2j + 1}{2i \sin \frac{c}{2}} \left( \exp(i(2j + 1) \frac{c}{2}) - \exp(-i(2j + 1) \frac{c}{2}) \right).
\]

A comparison shows that one obtains the discrete version of the \(\delta\)-function by integrating out the angles \(\theta\) and \(\varphi\), and replacing the continuous \(r\) with the discrete label \(j + \frac{1}{2}\). (One also has to replace \(\sin \frac{c}{2}\) with \(\frac{c}{2}\), but since both expressions represent a \(\delta\)-function in \(c\), they can be regarded as identical.) This clarifies the relation between continuous values for \(B\) and the discrete \(j\).

The calculation can also be used to show that the standard \(BF\)-face amplitude agrees with the one derived from the invariant path integral measure and thus is anomaly-free. For Euclidean gravity, \(SO(4)\) \(BF\)-theory is used which can be written as a state sum with \(SU(2)\)-valued variables after using \(SO(4) \cong SU(2) \times SU(2)\). For each copy of \(SU(2)\) on every face with spin \(j_f\) one has a factor contributing to the face amplitude by \(2j_f + 1\) which agrees with the last result in \((24)\) where a single factor \(r\) remains in the measure after integrating over \(\theta\). Discretizing \(r\) then yields the correct face amplitude (see Section 4.1).

5.1.2 A general recipe

For theories more complicated than \(BF\)-theory, but still written as constrained \(BF\)-theory, we have to combine the result of Section 3, which tells us the correct continuum measure in the presence of second class constraints, with a transition from the continuum measure to discrete amplitudes. We also have to expect additional functions which result from inserting \(\delta\)-functions imposing the constraints into the path integral. The constraint algebra tells us what function we have to include to obtain the correct continuum measure as in Eq. \((7)\), and there can be additional functions coming from Jacobians if the constraints do not directly restrict the integration variables but a more complicated functional of them.

A possible strategy is to use the formulation of \(BF\)-theory in polar \(B\)-coordinates where we have seen that the spin foam model with the correct face amplitude arises from integrating out the \(B\)-angles and discretizing the radial \(B\)-coordinate. Additional constraints restricting the \(BF\)-theory then arise from the action via integrating over the corresponding Lagrange multipliers which results in \(\delta\)-functions inserted into the path integral. We have seen that the multiplier integration requires a special measure which can be computed if
the constraint algebra is known. Next, we write both this function and the \( \delta \)-functions in the polar \( B \)-coordinates, integrate over the \( B \)-angles and discretize the radial \( B \)-coordinate just as we did in \( BF \)-theory without any additional constraints. In general, this will result in an additional functional besides the vertex amplitude in the spin foam model which depends on the spins \( j \).

In practice, the calculation, in particular the computation of the constraint algebra and the integration over the angular \( B \)-coordinates, will be complicated (several faces are coupled by the simplicity constraint). To illustrate the type of integrals involved we have included an example of constrained \( BF \)-theory in the appendix. It is also not easy to decide which part of the constraint algebra one has to consider. All constraints taken together form a mixed system containing first class constraints. Usually, as in Section 3, one would have to gauge fix the first class constraints resulting in a pure second class algebra. However, for gravity a gauge fixing is not known, and it is hoped that spin foam models can avoid the need to introduce an explicit gauge fixing by using a space-time discretization. In the following subsection we will see however that simple considerations of background independence with respect to active diffeomorphisms imply the existence of bubble divergences which are naturally associated to infinite gauge volume contributions. Diffeomorphisms in fact seem not to be completely gauge fixed by the discretization and the situation might be similar to that of \( BF \)-theory reviewed in the previous sections. One then would have to find a second class sub-algebra of the full constraint system plus the appropriate gauge fixing constraint in order to compute the correct measure.

It seems necessary to look for a description of gauge symmetries and constraint algebra that would be directly defined at the discrete level. In the previous section we have reviewed how this can be done in the case of \( BF \)-theory. Indeed, in order to regularize the path integral Freidel and Louapre had to introduce a gauge fixing of the topological symmetry (21) which in turn modifies the measure by the Fadeev-Popov determinant (22). In the case of this topological theory the modification is trivial and the factor is independent of the fields reducing simply to unity. In the case of gravity one should expect a non trivial dependence on spins consistent with a theory with local excitations. It is clear from this example that even when the gauge symmetries of the discrete action retain some similarities with the continuum ones their action can be only interpreted at the discrete level. An equivalent analysis in the case of 4-dimensional Plebanski theory seems necessary in order to implement the general prescription of Subsection 5.1.2 for the construction of the measure and hence settle the issue of lower dimensional simplex amplitudes. In such a context, the first class part of the constraint algebra might not even be first class in the continuum sense as results of Gambini and Pullin show [7, 8]. In this case a direct application of our recipe (or a slight generalization) should be feasible.

### 5.2 Active picture

Our aim is to derive conditions for the correct path integral measure for gravity in the context of spin foam models. In this section we will show how the requirement of background independence and diffeomorphism invariance from the active point of view imposes restric-
tions on the spin foam amplitudes with important consequences. The discussion presented here is general to any spin foam model of a diffeomorphism invariant theory. However, we will focus the attention on the Barrett–Crane model for quantum gravity. We shall perform our derivation in the context of the formulation of simplicial quantum gravity presented in [3]. In this formulation the Barrett–Crane model arises as the simplicial counterpart of Plebanski’s formulation of gravity.

5.2.1 Discretization consistency

The requirements imposed here on the spin foam amplitudes can be viewed as a 4-dimensional generalization of the notion of cylindrical consistency and diffeomorphism invariance in the canonical formulation of loop quantum gravity [25]. According to background independence, the cellular decomposition used to represent the space-time manifold does not carry any physical information. Gravitational degrees of freedom are encoded in the labeling of faces in the dual 2-complex with irreducible representations of the corresponding internal gauge group: a spin foam. As a consequence there remains some redundant information in a spin foam defined on a particular 2-complex which links the ‘physical’ configuration with the discretization on which it has been defined. The background independent information is encoded in the appropriate equivalence classes of spin foams. These equivalence classes have to be introduced if one wants to think of spin foams as morphisms in the spin network category [26]. Elements of a given equivalence class of spin foams can be related by the following moves:

1. (Piecewise linear) maps preserving the cell-complex and its coloring
2. Subdivision
3. Orientation reversal

A detailed definition can be found in [26]. These moves can be interpreted as the counterpart of diffeomorphisms in simplicial quantum gravity (with perhaps the addition of more equivalence relations if the remnant of diffeomorphism invariance is larger as in $BF$-theory).

The situation is analogous to that of the canonical formulation of loop quantum gravity where one essentially solves the diffeomorphism constraint by considering equivalence classes of spin networks under 3-diffeomorphisms. The spin foam equivalence classes correspond to the 4-dimensional generalization of this idea. In the following we will see how these considerations restrict to some extent the freedom in the definition of the spin foam measure.

The natural amplitude for faces in spin foam models is given by the Plancherel measure arising in the harmonic analysis on the corresponding internal gauge group (e.g., $Spin(4)$ and $SL(2,C)$ for the Riemannian and Lorentzian models respectively). This can be derived in various ways and it is directly linked with the notion of locality of spin foams [14]. Namely that degrees of freedom communicate along faces by boundary data given by the
Spin(4) and \(SL(2, C)\) connection respectively. If the face amplitude \(A_f(j)\) is given by the Plancherel measure (e.g., \(A_f(j) = (2j + 1)^2\) in the case of the Riemannian Barrett–Crane model) then the arbitrary subdivision of a face \(f \in J_\Delta\) does not change the amplitude. Amplitudes are invariant under subdivision of faces which produces combinatorially equivalent spin foams. Assigning different amplitudes to these would correspond to an anomaly (in the sense that members of the same equivalence class would be associated with different amplitudes). Recall that in background independent spin foam models the 2-complex has no geometrical meaning whatsoever and two spin foams for which a face has been subdivided in this way correspond to physically equivalent configurations. This requirement already rules out the normalizations of the model proposed in [17] to improve convergence properties of the finite model introduced in [19].

The previous analysis raises the question of whether we can find more stringent conditions on spin foam amplitudes by solely imposing background independence. In fact this is possible. We can obtain restrictions on both the face and edge amplitudes by imposing further necessary conditions for the anomaly freeness of the spin foam model. The full action of diffeomorphisms in the context of spin foams is not well understood. We have seen that the non trivial action of gauge transformation in \(BF\)-theory suggests that even in a simplicial theory a non trivial remnant of the diffeomorphisms might survive. The precise nature of this remnant is not understood but there are some basic symmetries that have to be represented there. These correspond to the discrete symmetries of the spin foams which are imposed by the requirement of background independence. As discussed above, two spin foams are to be thought of as physically equivalent whenever we can ‘deform’ one into the other by the action of the simplicial analog of a diffeomorphism, i.e. they are in the same equivalence class defined by Baez in [26]. A necessary consistency condition is that amplitudes should be well defined on the equivalence class of spin foams. Two equivalent spin foams must have the same amplitude as a necessary condition for anomaly freeness.

The simplest consistency condition can be obtained by the requirement that the bubble spin foams of Figures 1 and 2 have the same amplitude. The figure represents spin foams where most of the faces are labeled by the trivial representation except for the shown bubbles, labeled by the simple representation \(\rho = j \otimes j^*\). These spin foams are clearly equivalent according to the definition of equivalence in [26]. This is also in agreement with our intuitive notion of background independence (see [1] for more discussion): given that the underlying two complex does not carry any geometrical information and that geometry degrees of freedom are represented by the labeling of faces with spins and its intrinsic combinatorics there is no way to distinguish between these bubble spin foam excitations. Their apparent difference is linked to the fiducial background 2-complex used to represent them and should not play any physical role (this is true only if, as assumed, all outside faces are labeled by the trivial representation; otherwise they will not be related by the moves defined above).

Let us compute these amplitudes in the case of the Barrett–Crane model.\(^5\) The Barrett–Crane model is a definition of the 4-simplex or vertex amplitude for quantum gravity up to

\(^5\)The same analysis can be performed with any other model.
an overall factor. If we normalize the vertex amplitude in some arbitrary way we can shift the ambiguity to the value of the edge amplitude. The normalization we choose is that for which the Barrett–Crane intertwiner is a norm-one-vector in the Hilbert space where \( \text{Inv}[\rho_1 \otimes \rho_2 \otimes \rho_3 \otimes \rho_4] \) acts. This normalization is naturally obtained in the implementation of Plebanski’s constraints on the BF partition function [3].

\[ B_1(j) = (2j + 1)^{-1} A_e(j)^6 A_f(j)^4, \]  

(25)

where \( A_e(j) \) and \( A_f(j) \) are the so far undetermined edge and face amplitudes. We have given already an argument for the value of \( A_f(j) \) which we will re-derive here from the anomaly freeness condition. In the case of the bubble diagram on the right of Figure 1 we have six vertices corresponding to the same 10j-symbol as before, so the amplitude for the prism bubble \( B_2(j) \) is

\[ B_2(j) = (2j + 1)^{-6} A_e(j)^9 A_f(j)^5. \]  

(26)

Finally in the case of the bubble spin foam on the right of Figure 2 we have four vertices of the previous type plus the vertex on the top whose 10j-symbol is illustrated in the center diagram of Figure 3 which evaluates to \((2j + 1)^{-2}\). With this the amplitude of the pyramid bubble \( B_3(j) \) is

\[ B_3(j) = (2j + 1)^{-6} A_e(j)^8 A_f(j)^5. \]  

(27)

The requirement \( B_1(j) = B_2(j) = B_3(j) \) fixes the values of \( A_e(j) \) and \( A_f(j) \) uniquely to

\[ A_e(j) = 1 \quad \text{and} \quad A_f(j) = (2j + 1)^2. \]  

(28)

The previous bubble spin foams are particularly easy to compute and helpful to explain the intuitive idea behind our consistency requirement. There is a more general statement of
this property to be satisfied by any anomaly free spin foam. Namely, spin foam amplitudes are required to be invariant under the arbitrary subdivision of their faces. Equivalently, if we deform (by a piece-wise linear homeomorphism$^6$) a colored face by coloring with the same spin adjacent faces in the 2-complex (previously labeled by the trivial representation) the amplitude should remain invariant. We see that with the normalization found above the Barrett–Crane model satisfies this necessary condition for anomaly freeness as the amplitude of the composite face $A_{j}^{\text{com}}(j)$ can be easily shown to be given by

$$A_{j}^{\text{com}}(j) = (2j + 1)^{2n_e - 2n_e + 2n_f} = (2j + 1)^{2x} = A_{j}(j),$$

where $n_e$ and $n_f$ is the number of internal edges and faces of the composite face. We see that (28) yields an invariant amplitude.

The analysis here fixes the value of the face amplitude to be given by the Plancherel measure of the corresponding gauge group. It is important to notice however that the edge

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$^6$For an extensive analysis of role of piece-wise linear homeomorphisms as opposed to diffeomorphism as basic symmetry of quantum gravity see [27].
amplitude found here is only valid for the degenerate situations in which we have vertex configurations of the form illustrated in Figure 3. These situations are degenerate in the sense that the simplicity constraints of that reduced $BF$-theory to gravity are trivial. In the general situation one expects the value of the edge amplitude to differ from the trivial value obtained here. It is precisely here where the appropriate Fadeev–Popov factors advocated in Sections 3 and 5.1.2 will play an important role.

The requirements studied here should be met by any theory admitting a spin foam quantization. In particular it is easy to see that $BF$-theory in any dimension would satisfy them. This is however to some extent trivial as $BF$-theory is topological and has a finite number of degrees of freedom. The following is a simple example of an (in this sense) anomaly free spin foam model for a theory with infinitely many degrees of freedom.

5.2.2 An anomaly free toy model: a trivial example

In this section we define a spin foam model satisfying the above minimal requirements of anomaly freeness whose vertex amplitude is in a certain sense the simplest compatible with these requirements. The model is tailored to produce a physically interesting model that can be thought of as a spin foam quantization of the Husain–Kuchar model [28].

Let us start by assuming that the model is defined on a simplicial decomposition. In this case all vertices in the dual 2-complex are 5-valent and boundary graphs have 4-valent nodes. The vertex amplitude of the model is given by the $BF$-theory vertex amplitude, namely the corresponding 15$\jmath$-symbol constructed with normalized intertwiners (where $\jmath$ are unitary irreducible representations of the corresponding compact group $G$) whenever at least three links forming a triangle in the graphical representation of the 15$\jmath$-symbol vanish, and zero otherwise. In other words, the amplitude for the creation of new 4-valent nodes is zero and transition amplitudes can be non-vanishing only when they involve an initial and final spin-network with the same number of 4-valent nodes (with the addition of disconnected Wilson loops with no intersections). However, lower valence nodes can be created and have a well defined amplitude.

This model can produce infinite transition amplitudes whenever a ‘vacuum’ bubble of the kind represented in Figures 1 or 2 is created. Vacuum bubbles appear when we have disconnected Wilson loops in intermediate states. Notice that more general bubbles – namely those surrounded by faces labeled with non trivial representations– are not allowed by the dynamics encoded in the vertex amplitude. Therefore, divergences are of a trivial nature corresponding to an overall infinite factor which can be regularized by factoring out the amplitudes for closed loops with no intersections which as in $BF$-theory are physically equivalent to a $c$-number (e.g., if $G = SU(2)$ this $c$-number is $2j+1$). The model is anomaly free in the sense above if we set the face amplitude as the corresponding Plancharel measure of the Lie group of interest to the power of the Euler characteristic of the face. For exterior faces one has to modify the amplitude in the usual way by adapting the definition of the Euler characteristic by counting by 1/2 edges with one endpoint on the boundary and 0 for edges with two endpoints on the boundary and external faces. The (regularized) transition amplitude between two spin network states is one if the two spin networks are contained.
in the same (piecewise linear homeomorphism) equivalence class and zero otherwise.

The model can be generalized to the case of arbitrary cellular decompositions. This allows for the computation of arbitrary (piecewise linear graph based) spin network to spin network transition amplitudes. Namely, we can construct the generalized projection operator $P$ to the physical Hilbert space of the theory. A key ingredient in this definition is that amplitudes are independent of the cellular decomposition as a simple consequence of the topological invariance of $BF$-theory. The continuum limit can be defined as in [29]. In that limit, the physical Hilbert space defined by $P$ is much larger than that of BF theory. It corresponds to the Hilbert space of spin network states modulo the 3-diffeomorphisms and therefore to the physical Hilbert space of the Husain–Kuchar model.

![Diagram](image)

Figure 4: Diagrammatic representation of two distinct Wilson loops that will be physically equivalent in our toy model. In general any two spin network states differing by a piecewise linear homeomorphism will be physically equivalent, i.e., their difference will be in the kernel of $P$. In the continuum limit one can extend the equivalence to smooth graphs as the one shown on the right.

Amplitudes in this model are crossing symmetric in the sense of [30, 31]. We are imposing 3-diffeomorphism invariance in any arbitrary slicing. The purpose of this simple example is to show how the requirement of 3-diffeomorphism invariance is directly related to the values of the face and vertex amplitudes in the special configurations studied in the previous subsection. In other words, the non triviality of gravity transition amplitudes (or more precisely its generalized projection operator $P$) should be encoded in the details of the vertex amplitude for the configurations that have been avoided de facto in this model and otherwise agree with it.

6 Discussion

The aim of this paper is to point out the need to check the anomaly problem for spin foam models which in this context means that lower dimensional simplex amplitudes are not free but essentially fixed. This is illustrated by the toy model of Section 2 which demonstrates that without the correct amplitudes non-physical results arise.

In order to find an anomaly-free formulation one has to study the constraint algebra which in spin foam models for gravity, formulated as a restricted $BF$-theory, always involves second class constraints. For a complicated constraint algebra this requires a non-trivial function in the measure which has been overlooked before. Furthermore, the constraints,
imposed by $\delta$-functions, lead to additional functions in the measure which all contribute to the face and edge amplitudes.

We have proposed a way to address the issue of computing the measure in a simplicial theory starting from the formal path integral measure in the continuum. To get the face and edge amplitude of the discrete spin foam model we have seen that it is helpful to introduce polar coordinates for the $B$-field since it provides a direct link to the spin parameters of the state sum. Explicit calculations require several integrations which, if not possible to be done explicitly, are well-suited for a stationary phase approximation. One has to be careful, however, because the integrand can be singular. A detailed study is necessary in order to be able to judge if amplitudes of a particular model would be finite. We hope to come back to this issue for the case of gravity in a future publication.

Even though we have not explicitly computed the anomaly free measure for spin foam models of gravity we have pointed out that some minimal requirements can be imposed that severely restrict the value of the face and edge amplitudes. The restrictions imposed by anomaly freeness in the way presented here imply that bubble divergences in spin foam models for gravity are linked with the gauge action of diffeomorphisms. We have shown that these restrictions rule out some proposals in the literature. In particular, without modification of the singular edge amplitude (associated with edges bounded by less than four faces), the finite normalizations for both the Riemannian and Lorentzian Barrett–Crane model proposed in [19, 32] are to be regarded as formulations where the diffeomorphism gauge symmetry has been broken by an anomalous path integral measure. Other anomalous formulations are: the new normalization of the Barrett–Crane model proposed in [17] to improve the convergence properties of the previous model and Model A in [33]. This seems to severely limit the physical relevance of such proposals. Model B in [33] is the only normalization of the Barrett–Crane model which satisfies the minimal requirements of anomaly freeness presented here and is the one naturally arising in the quantization of Plebanski’s $Spin(4)$ formulation of gravity presented in [3]. However, a non trivial modification of the edge amplitude for generic edges should appear due to the contribution of the simplicity constraints as argued in Section 3.

The value of the correct edge amplitude can be determined if we understand the canonical algebra of simplicity constraints so that the appropriate determinant as in (7) is included. This is a complicated issue as it might require the understanding of the canonical formulation of the simplicial model. Perhaps the ideas of Gambini and Pullin in the context of their consistent discretization formulation might shed some new light on this issue. Unfortunately the canonical formulation of discrete theories seems rather complicated in the case of Plebanski’s formulation at this stage.

An interesting alternative procedure to deal with constrained systems is the projection-

\footnote{This (anomalous) finite normalization of the Barrett–Crane model was naturally obtained in the context of the group field theory (GFT) formulation of the model. One could modify the amplitudes of singular edges and make the model anomaly free in the sense described here. In this way there will be divergences but of a rather simple kind (only isolated vacuum bubble will diverge and could be easily renormalized). However, the naturality argument in relation to a GFT formulation would not stand in this case.}
operator approach [34, 35, 36]. Since this method allows to deal with first and second class constraints on an equal footing, it may be possible to sidestep some of the difficulties mentioned above.

The well understood results in 3-dimensional gravity [6] and those of Section 5.2.1 suggest that the discretization does not completely fix the diffeomorphism gauge transformations as it is usually assumed. In the case of the Barrett–Crane model (any other model could be analyzed in this way) we have shown that the minimal requirements for anomaly freeness already imply the presence of certain bubble divergences that have no physical content. For instance, even in our toy (Husain–Kuchar) spin foam model divergences can not be avoided without some extra manipulation. In this simple case one can regularize the model by dividing out the trivial bubble divergences.

These divergences will be present in any spin foam model for a background independent theory and the way to deal with them is by appropriate gauge fixing conditions. This is in fact possible in the case of three dimensional gravity. In four dimensions one would need to understand in a precise manner the action of 4-diffeomorphisms in the context of the simplicial models. Our minimal requirements of anomaly freeness of Section 5.2.1 are closely related to the action of 3-diffeomorphisms in loop quantum gravity (it is tempting to think that due to the fact that our requirements hold for any ‘slicing’ of the spin foam one is imposing a ‘bit’ of 4-diffeos in this picture) and are imposed by the requirement of background independence. In this sense, what is left to understand is the old question: where and what, is the remnant gauge transformation encoded in the action of the Hamiltonian constraint in the canonical framework?

Although one would need to understand the technical details of the generalization to 4-dimensional $BF$-theory, the divergence structure of the Crane-Yetter model suggests that the Freidel–Louapre analysis should go through in this case. Going to the Barrett–Crane model we constrain the $B$ field to be given by a simple bivector field derived from a tetrad. At the level of the spin foam model, $Spin(4)$ representations are constrained to simple representations of the type $j \otimes j^*$. The essential questions are: what part of ‘translational’ gauge freedom (21) in the $B$’s of $Spin(4)$ $BF$-theory remains after the implementation of the simplicity constraints? Namely: Is this gauge symmetry remnant fully encoded in the equivalence class of spin foams defined by Baez?, or is there also a symmetry that can change the values of the representation labels as in $BF$-theory?

In the first scenario the gauge divergence structure of the model does not seem problematic. We have seen in Section 3 that in addition to the Fadeev–Popov factor coming from gauge fixing first class constraints the measure should be modified when implementing the simplicity constraints. From the analysis of Section 5.2.1 we conclude that this modification should involve the value of the edge amplitude for non-singular edges. The implementation of the simplicity constraints could modify the amplitudes in a way that would make non-singular bubble amplitudes finite. One could think of the non trivial damping edge amplitude of the finite model [9] as arising in this way for non-singular edges. These factors should not arise for singular edges bounding only two non trivial faces as our consistency argument fixes the edge amplitude to unity in these cases. In this scenario the divergence structure of the (non-gauge fixed) amplitudes would be similar to
that of our toy model of Section 5.2.1: only the vacuum bubbles as the ones represented in Figures 1 or 2 would be the divergent contributions to the amplitudes and can be easily regularized.

In the second scenario, it is appealing to think that the sum over representations ‘flowing’ inside a bubble is un-physical and its contribution correspond to the diffeomorphism gauge volume. Gauge fixing will correspond to dropping these redundant sums and replace them by the appropriate Fadeev–Popov determinants of Section 3. In the case of $BF$-theory topological invariance follows from the triviality of these Fadeev–Popov determinants (they are all equal to one due to the presence of the $\delta$-functions that set the curvature to zero [6]). In the case of gravity they should depend in a non-trivial fashion on the spins as the correct model must contain local excitations.

Further possibilities to explore are the following ones: Second class constraints require a special function to be inserted into the path integral measure which can be derived from the constraint algebra. This requires a Hamiltonian analysis of the full constraint system and thus will in general lead to a breaking of manifest covariance, which is usually regarded as a key advantage of the spin foam approach. (To see this in the case of the Plebanski action it is enough to note that only timelike components of $B$ serve as Lagrange multipliers which are restricted by the simplicity constraints and thus have momenta constrained to be zero.) Such a phenomenon is not special to gravitational systems and has been observed before (see, e.g., [37] for a discussion and examples). Usually, one can see that the non-covariance of this function is canceled by other non-covariant terms which arise in the regularization process. For spin foam models, however, the discretized amplitude obtained by ignoring the additional function is covariant such that it would be impossible to cancel the non-covariance in the additional function. This observation suggests that non-trivial effects have to happen in the continuum limit which will eventually cancel non-covariant contributions. A computation of the constraint algebra and the non-covariant terms in the measure may lead to important clues as to what has to occur in the continuum limit.

We have proposed concrete ways to address the issue of the definition of the correct physical measure for spin foam models of quantum gravity. The meaning of diffeomorphisms in the context of spin foam models for 4-dimensional quantum gravity is a pressing subject which demands a detailed analysis. It is our view that no meaningful progress in understanding other important questions in the approach such as the constructions of observables, continuum limit, etc.--can be achieved without giving a definite answer to the issues that we are stressing here.

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Appendix: A degenerate sector of gravity

In this appendix we illustrate how the study of the integrals that define the amplitude of spin foam models obtained by discretizing a continuous action can be used to find the correct normalization of the spin foam measure (in the large spin limit). The example presented here is a little bit more involved than the one of Subsection 5.1.1 and is closely related to the Barrett-Crane model. Even though the analog of simplicity constraints is present in this example, it should be pointed out that our analysis here has an important limitation: we are not including the determinant factors appearing in (7) produced by the implementation of second class constraints in the path integral. This example is however meant as a computation that illustrates the difficulties involved in the analysis. We have kept this calculation as an appendix hoping that the technique can be useful for further developments.

A degenerate sector of Euclidean gravity can be obtained by introducing the additional constraint [38]

\[ B^L_i = V^j_i B^R_j \]  

(30)

into SO(4) BF-theory where \( B^L \) and \( B^R \) are \( su(2) \) valued components of the \( so(4) \) valued \( B \) according to the decomposition \( SO(4) \cong SU(2) \times SU(2) \), and \( V \) is an \( SO(3) \)-matrix. The new action then is

\[ S = \int (B^L_i \wedge F^L_i + B^R_i \wedge F^R_i + \lambda^i \wedge (B^L_i - V^j_i B^R_j)) \]

where \( F^L_i \) and \( F^R_i \) denote the curvatures of the left and right component of the connection, respectively. In this form, \( \lambda^i_{ab} \) and \( V^j_i \) appear as new Lagrange multipliers in addition to the multipliers \( B^L_{0a} \) and \( B^R_{0a} \) of the original BF-theory. Only the multipliers \( B^L_{0a} \) and some components of \( \lambda^i_{ab} \) are restricted by the new constraints and in order to obtain a constrained system on a symplectic phase space we have to add their momenta to the canonical variables together with constraints requiring them to be zero. The algebra of all the constraints will determine the measure we have to use in the path integral. For a discussion of the finiteness of the resulting model it is most interesting to see whether or not the additional factor depends on components of \( B \) with a non-trivial scaling behavior since this would affect the large-\( j \) behavior of face amplitudes. This has to be expected here because varying \( V^j_i \) yields a constraint which restricts components of the multipliers \( \lambda^i_{ab} \) and is linear in \( B \). Since we had to add momenta of these components of \( \lambda^i_{ab} \), the constraint algebra will contribute positive powers of \( B \) to the measure which would enhance a divergence.

Here, however, we do not enter a detailed discussion of the constraints; rather we will ignore the additional factor and derive the large-\( j \) behavior of the face amplitude for the naive spin foam quantization with a trivial multiplier measure. While the resulting face...
amplitude would not be correct, our aim here is solely to compare this calculation with the result of [3] where this factor has been ignored, too.

To include the additional constraint into a spin foam quantization we impose the constraint face-wise such that any face \( f \) carries a matrix \( V_f \). This gives the state sum

\[
Z = \int \prod_e d^3 g_e^L d^3 g_e^R \prod_f d^3 B_f^R d^3 V_f \exp \left( i \text{tr} \left( B_f^R (U_f^R + U_f^L V_f) \right) \right)
\]

\[
= \int \prod_e d^3 g_e^L d^3 g_e^R \prod_f d^3 B_f^R d^3 B_f^R d^3 V_f d^3 \lambda_f \exp \left( i \text{tr} \left( B_f^R U_f^R + \frac{1}{2} \lambda_f (B_f^L - V_f B_f^R) \right) \right)
\]

where \( \lambda_f \) are Lagrange multipliers and the edge holonomies \( (g_e^L, g_e^R) \) form the holonomies \( U_f^L \) and \( U_f^R \) along closed loops. Integrating over \( \lambda_f \) yields \( \delta \)-functions which will be solved after introducing polar coordinates \( (r_f^L/R, \vartheta_f^L/R, \varphi_f^L/R) \) for \( B_f^L/R \) and Euler angles \( (\psi, \theta, \phi) \) for \( V \). The \( \delta \)-functions then imply \( r_f^L = r_f^R \) for all faces \( f \), and \( \vartheta^L \) and \( \varphi^L \) will be given as functions \( \vartheta^L = F(\vartheta^R, \psi - \varphi^R, \theta) \) according to

\[
\cos \vartheta^L = \sin \theta \sin \vartheta^R \sin (\psi - \varphi^R) + \cos \theta \cos \vartheta^R
\]

and \( \varphi^L = G(\vartheta^R, \varphi^R, \psi, \theta, \phi) \). Choosing again a gauge for \( U_f^{L/R} \) without loss of generality, we obtain

\[
Z = \int \prod_e d^3 g_e^L d^3 g_e^R \prod_f d^3 B_f^R d^3 B_f^R d^3 V_f \delta^3(B_f^L - V_f B_f^R) \exp \left( i \text{tr} \left( B_f^R U_f^R + B_f^L U_f^L \right) \right)
\]

\[
= \int \prod_e d^3 g_e^L d^3 g_e^R \prod_f d r_f^L d \vartheta_f^L d \varphi_f^L (r_f^L)^2 \sin \vartheta_f^L d \vartheta_f^R d \varphi_f^R (r_f^R)^2 \sin \vartheta_f^R d \psi_f d \vartheta_f d \varphi_f
\times \left( -i r_f^L \cos \varphi_f^L \sin \frac{\vartheta_f^L}{2} + r_f^R \cos \varphi_f^R \sin \frac{\vartheta_f^R}{2} \right)
\times \exp \left( -i \left( r_f^L \cos \varphi_f^L \sin \frac{\vartheta_f^L}{2} - r_f^R \cos \varphi_f^R \sin \frac{\vartheta_f^R}{2} \right) \right)
\times \exp \left( -i \left( \frac{r_f^L}{2} \sin \frac{\vartheta_f^L}{2} - \frac{r_f^R}{2} \sin \frac{\vartheta_f^R}{2} \right) \right)
= \int \prod_e d^3 g_e^L d^3 g_e^R \prod_f d r_f^R d \vartheta_f^R d \varphi_f^R (r_f^R)^2 \sin \vartheta_f^R d \psi_f d \vartheta_f d \varphi_f
\times \exp \left( -i \left( \frac{r_f^R}{2} \sin \frac{\vartheta_f^R}{2} - \frac{r_f^R}{2} \sin \frac{\vartheta_f^R}{2} \right) \right)
\times \exp \left( -i \left( \frac{r_f^R}{2} \sin \frac{\vartheta_f^R}{2} - \cos \frac{\vartheta_f^R}{2} \right) \right).
\]

We now have to perform the angle integrations, which is effectively a four-dimensional integral since the \( \varphi \)- and \( \psi + \varphi^R \)-integrations are trivial, in order to find the face amplitude after discretizing \( r^R \). This integral can be approximated by using the saddle point approximation method, but one has to be careful because the integrand has zeroes due to the sines.

Disregarding the zeroes, one would expect the asymptotic behavior of the angle integral to be \( (r^R)^{-2} \) because every one-dimensional integration contributes \( r^{-\frac{1}{2}} \) in the standard stationary phase result. This factor would cancel the \( (r^R)^2 \) from the measure and thus yield 1 as the large-\( j \) behavior of the face amplitude. In [3] a face amplitude 1 has been
obtained for this model by different means, however for a different discretization of the action which is not accessible to our methods. Nevertheless, there is agreement provided that the modified discretization in [3] corresponds to the standard stationary phase result. An analysis similar to that of Section 5.2.1 would lead to the conclusion that such a face amplitude is in conflict with background independence. This should not be surprising as there is no reason for amplitudes to be anomaly free when we do not include the correct measure in the presence of second class constraints. At present we do not see whether doing so in this case would lead to amplitudes in agreement with background independence.

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