

# Quantum aspects of 2+1 gravity

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We review and systematize several recent attempts to canonically quantize general relativity in 2+1 dimensions, defined on space–times  $\mathbb{R} \times \Sigma^g$ , where  $\Sigma^g$  is a compact Riemann surface of genus  $g$ . The emphasis is on quantizations of the classical connection formulation, which use Wilson loops as their basic observables, but results from the ADM formulation are also summarized. We evaluate the progress and discuss the possible quantum (in)equivalence of the various approaches. © 1995 American Institute of Physics.

## I. INTRODUCTION

The aim of this paper is to review and systematize attempts to canonically quantize general relativity in 2+1 space–time dimensions. We limit the scope to the case of pure Einstein gravity, possibly with a nonvanishing cosmological constant  $\Lambda$ , on three-dimensional manifolds of the form  $\mathbb{R} \times \Sigma^g$ , where  $\Sigma^g$  is a compact Riemann surface of genus  $g \geq 1$ . The inclusion of point sources or matter, the possibility of topology changes, and lattice approaches will not be discussed. Basic knowledge of the classical theory and its geometric interpretation will be presupposed, and wherever necessary we will refer to the appropriate references for details. Our emphasis will be on a comparison between the different existing quantization methods, and their possible quantum (in)equivalence. This includes issues like the degree of the classical reduction, the choice of the basic variables to be quantized, and the role of the mapping class group (the “large diffeomorphisms”).

If one wants to regard 2+1 gravity as a model for the (3+1)-dimensional theory, with the local physical excitations substituted by a finite number of topological degrees of freedom, it is a somewhat embarrassing fact that many of the quantizations proposed have hardly progressed beyond the genus-one case, where the spatial manifold  $\Sigma$  is a torus  $T^2$ . On the one hand, one finds a rich structure already in this case, but on the other one knows that its mathematical structure is not very representative of (and rather simpler than) the general higher-genus case. This should be kept in mind when drawing conclusions about canonical quantum gravity in general.

We shall distinguish between two different (but classically essentially equivalent) approaches to 2+1 gravity, (i) the geometric description in terms of the Lorentzian three-metric  ${}^{(3)}g_{\mu\nu}$ , and (ii) a gauge-theoretic description in terms of a connection one-form  ${}^{(3)}A_\mu^a$ , taking values in an appropriate gauge algebra. The first one is the three-dimensional analog of the well-known ADM framework, and involves the choice of a parameter representing physical time. In the gauge-theoretic approach, no such choice is required, and one has the option between a Chern–Simons formulation (with a noncompact gauge group) and a closely related three-dimensional version of the Ashtekar formulation of general relativity. In all of these descriptions, 2+1 gravity takes the form of a canonical system with first-class constraints à la Dirac. What distinguishes it from other constrained field theories is the fact that one can solve (part of) the constraints already classically to the extent that the remaining reduced phase space is finite dimensional. This reduces the quantization problem to a quantum-mechanical one. Still, as we will see, there is no unique way of setting up a quantum theory. Different proposals have been made, depending on the classical starting point, and on various requirements one may wish to impose on the quantum theory. Another special feature (and a consequence of the invariance under space–time diffeomorphisms) is the absence of an *a priori* defined time parameter, leading to further ambiguities in the quantization.

In the familiar geometric approach, Einstein gravity with a cosmological constant is defined by the Lagrangian

$$S[g] = \int d^3x \sqrt{-^{(3)}g} (^{(3)}R - 2\Lambda) = \int dt \int_{\Sigma} d^2x (\pi^{ij} \dot{g}_{ij} - N^i \mathcal{H}_i - N \mathcal{H}), \quad (1.1)$$

where a (2+1)-decomposition à la ADM has been performed to arrive at the second expression, and the canonically conjugate variable pairs consist of the spatial metric  $g_{ij}(x)$  and the momentum  $\pi^{ij}(x)$  on the surfaces  $t=\text{const}$  which depends on the extrinsic curvature  $K^{ij}$  via  $\pi^{ij} = \sqrt{^{(2)}g} (K^{ij} - g^{ij}K)$ . (In this paper, we will use  $\mu, \nu, \dots$  for space-time indices,  $i, j, \dots$  on two-dimensional spatial slices, and  $a, b, \dots$  for internal indices.) The spatial diffeomorphism and Hamiltonian constraints are given by

$$\mathcal{H}_i = -2\nabla_j \pi^j_i = 0, \quad (1.2)$$

$$\mathcal{H} = \frac{1}{\sqrt{^{(2)}g}} g_{ij} g_{kl} (\pi^{ik} \pi^{jl} - \pi^{ij} \pi^{kl}) - \sqrt{^{(2)}g} (R - 2\Lambda) = 0$$

(multiplied in the Lagrangian by the shift and lapse functions  $N^i$  and  $N$ ), and thus have the same functional form as in the (3+1)-dimensional case.

## II. QUANTIZATION IN THE CONNECTION FORMULATION

Since it was the gauge-theoretic reformulation of 2+1 gravity and the demonstration of its “exact solubility” by Witten<sup>1</sup> that sparked off much of the recent interest in the theory, we will deal with this case first. It is a remarkable fact that in both three and four space-time dimensions, Einstein–Hilbert gravity may be reformulated on a Yang–Mills phase space. How this leads to the Ashtekar formulation of 2+1 and 3+1 and the Witten formulation of 2+1 gravity has been discussed by Bengtsson.<sup>2</sup>

Witten rewrites the action for 2+1 gravity as that of a Chern–Simons theory for the Poincaré group ISO(2, 1) in three dimensions, the group SO(2, 2) and SO(3, 1) for the cases  $\Lambda=0$ ,  $\Lambda<0$ , and  $\Lambda>0$ , respectively.<sup>1</sup> To this end, one rearranges the dreibein  $e^a_\mu$  and the spin connection  $\omega^a_\mu{}^b$  of the geometric first-order formulation of the Einstein–Hilbert action into a gauge algebra-valued connection form  $A_\mu = e^a_\mu P_a + \omega^a_\mu J_a$ . (Recall that the three-metric in this ansatz is the derived quantity  $g_{\mu\nu} = e^a_\mu e^b_\nu \eta_{ab}$ , where  $\eta$  is the three-dimensional Minkowski metric.) The algebra generators  $P_a$  and  $J_a$  fulfill

$$[J_a, J_b] = \epsilon_{abc} J^c, \quad [J_a, P_b] = \epsilon_{abc} P^c, \quad [P_a, P_b] = \Lambda \epsilon_{abc} J^c, \quad (2.1)$$

where it is understood that internal indices are raised and lowered using  $\eta_{ab}$ . The Lagrangian in this approach becomes

$$\begin{aligned} S'[^{(3)}A] &= \frac{1}{2} \int d^3x \epsilon^{\mu\nu\lambda} \text{Tr} \left( A_\mu \partial_\nu A_\lambda + \frac{1}{3} A_\mu [A_\nu, A_\lambda] \right) \\ &= \int dt \int_{\Sigma} d^2x \epsilon^{ij} (-e_{ia} \dot{\omega}_j^a + \lambda_A F[^{(2)}A]_{ij}^A), \end{aligned} \quad (2.2)$$

where a 2+1 decomposition has been performed in the second step, the index  $A$  labels the six generators  $T_A = (J_a, P_a)$ , and  $F$  is the two-dimensional field strength of the spatial part of the connection  $^{(3)}A$ , taking values in the algebra of the appropriate gauge group  $G$ , where  $G = \text{ISO}(2, 1)$ ,  $\text{SO}(3, 1)$ , and  $\text{SO}(2, 2)$ , depending on the value of the cosmological constant  $\Lambda$ .

(Note that the  $\Lambda$ -dependence is implicit in the commutators  $[A, A]$ .) For  $\Lambda=0$ , this action is equivalent (at least for nondegenerate dreibeins) to the three-dimensional Ashtekar action<sup>3</sup>

$$S''[e, \omega] = \frac{1}{2} \int d^3x \epsilon^{\mu\nu\lambda} e_{\mu a} F_{\nu\lambda}^a = \int dt \int_{\Sigma} d^2x \epsilon^{ij} (-e_{ia} \dot{\omega}_j^a + \mu_a \mathcal{L}_i e_j^a + \nu_a F^{(2)} \omega_{ij}^a), \tag{2.3}$$

where  $F^{(2)}[\omega]$  is now the two-dimensional field strength of the  $SO(2, 1)$ -connection, and  $\mathcal{L}$  its covariant derivative. (This may be considered as a special case of the equivalence between a Palatini action for gauge group  $G$  and a Chern–Simons action for the corresponding inhomogeneous group  $IG$ .<sup>4</sup>) Defining  $E^{bj} = \epsilon^{jk} e_k^b$ , we obtain in both cases a phase space with symplectic structure

$$\{\omega_{ia}(x), E^{jb}(y)\} = \delta_i^j \delta_a^b \delta^2(x-y) \tag{2.4}$$

(that is, six variable pairs per space point). The Hamiltonian densities are linear combinations of six first-class constraints each. In the Chern–Simons formulation, they constrain the spatial component of the field strength of the six-dimensional gauge algebra to vanish, i.e., the corresponding connection  $A$  to be flat, whereas in the Palatini formulation they are made up of three Gauss law constraints for  $SO(2, 1)$  and a flatness condition on the spin connection with its *three*-dimensional gauge algebra. In the first case, the physical, classically reduced phase space can therefore be identified with the space of flat  $G$ -connections modulo  $G$ -gauge transformations. For  $\Lambda=0$ , the alternative Palatini description yields a physical phase space that is a cotangent bundle over the reduced configuration space of flat  $SO(2, 1)$ -connections modulo  $SO(2, 1)$ -gauge transformations. This phase space coincides with the one of the Chern–Simons formulation for  $G=ISO(2, 1)$ , because the group manifold of  $ISO(2, 1)$  is isomorphic to that of  $T^*SO(2, 1)$ . Note that for nonvanishing  $\Lambda$  the reduced phase spaces are *not* cotangent bundles.

For the quantization of these reduced phase spaces it is important to know concrete representations of the abstract quotient spaces, and to decide which classical observables are to be carried over to the quantum theory. Witten in his original paper envisaged a quantization based (for  $g \geq 2$ ) on a set of  $2g$   $G$ -valued holonomy variables  $U_i, V_i, i=1\dots g$ , corresponding to the  $2g$  generators  $\alpha_i, \beta_i, i=1\dots g$  of the homotopy group  $\pi_1(\Sigma)$ ,

$$U_i = P \exp \oint_{\alpha_i} A^F, \quad V_i = P \exp \oint_{\beta_i} A^F, \tag{2.5}$$

where  $A^F$  is a flat  $Lie(G)$ -valued connection one-form on a spatial slice. The  $2g$  holonomies contain all information necessary for reconstructing the moduli space of flat connections modulo gauge, and can be interpreted as gluing data for simply connected patches of Minkowski, de Sitter, and anti-de Sitter space, corresponding to  $G=ISO(2, 1), SO(3, 1)$ , and  $SO(2, 2)$  (for example, see Ref. 5 for how the geometric structure can be reconstructed from the holonomy variables). However, they are not free parameters, but (i) are subject to residual gauge transformations  $U_i \rightarrow g U_i g^{-1}, V_i \rightarrow g V_i g^{-1}$  at some arbitrary common base point of the  $\alpha_i$  and  $\beta_i$  in  $\Sigma$ , and (ii) must obey a relation

$$U_1 V_1 U_1^{-1} V_1^{-1} \dots U_g V_g U_g^{-1} V_g^{-1} = \mathbf{1}, \tag{2.6}$$

which comes from the analogous defining relation among the homotopy generators. In terms of these variables, the counting of physical degrees of freedom is therefore  $2g \times \dim(G) - \dim(G) - \dim(G) = (2g - 2) \times \dim(G)$ .

Thus we may use this parametrization to describe the reduced phase spaces of the Chern–Simons theory for the group  $G$ . In the case of vanishing cosmological constant, we have the

additional option of describing the reduced *configuration* space by holonomy variables, setting  $G=SO(2, 1)$  (because in this case there exists a natural division between coordinates and momenta, a “polarization” of phase space).

Concrete quantization proposals have gone still one step further, and based themselves on a set of explicitly gauge-invariant variables, so-called Wilson loops, obtained from the holonomies by taking traces,

$$T(\gamma)[A^F] := \text{Tr } P \exp \oint_{\gamma} A^F, \quad (2.7)$$

where  $\gamma$  is some element of  $\pi_1(\Sigma)$ . The Wilson loops are well known from their role as observables in usual gauge field theories (see, for example, Ref. 6). However, some problems stand in the way of using them as “coordinates” on the physical moduli spaces. Note first that it is not enough to use only the Wilson loops of the fundamental generators, since this would give us just  $2g$  degrees of freedom [out of the  $(2g-2) \times \dim(G)$  needed]. One thus has to include Wilson loops for more general elements of  $\pi_1(\Sigma)$ . However, these in general are subject to certain algebraic constraints, the so-called Mandelstam identities. In addition, there will be constraints on the Wilson loops coming from the defining relation (2.6).

Another question is that of the completeness of the Wilson loops, i.e., whether they separate all points of the moduli spaces (so that no two physically distinguishable  $A$ -configurations share the same values for all Wilson loops). This is believed to be the case for the compact gauge groups that typically appear in gauge theory, but is not *a priori* clear for the noncompact groups used in gravitational applications, although it has been shown that for gauge group  $SO(2, 1)$  the traced holonomies are essentially complete.<sup>3</sup> A related problem has been investigated in the Ashtekar formulation of 3+1 gravity, where the gauge group is  $G=SL(2, C)$ .<sup>7</sup> There also is the possibility of the existence of inequalities between the Wilson loops, although this is known not to happen for the hyperbolic sector of the  $G=SO(2, 1)$ -moduli space.<sup>8</sup> A problem of incompleteness seems to occur for the case  $\Lambda > 0$ , where the gauge group is  $SO(3, 1)$ , as has been remarked by various authors.<sup>9–12</sup>

Another subtlety arises when one tries to make connection with the geometric formulation in terms of a positive definite spatial metric  $g_{ij}$ . It turns out that the physical phase spaces there correspond to a certain subsector of the moduli spaces introduced above. (Other sectors also contain solutions to Einstein’s equations, but do not have the “correct” signature for  $g_{ij}$ , i.e., in general will allow for closed timelike curves.) For example, for  $G=SO(2, 1)$ , those are the configurations where all holonomies are “hyperbolic,” i.e.,  $\text{Tr } U_i > 2$ ,  $\text{Tr } V_i > 2$  in the two-dimensional representation, which (for  $g \geq 2$ ) together form the so-called Teichmüller space  $\mathcal{T}(\Sigma^g)$  associated with the Riemann surface  $\Sigma^g$ . Since they represent the true gravitational degrees of freedom, the most obvious strategy is to quantize only them, ignoring the remaining sectors of the moduli spaces. For  $g=1$ , which is qualitatively different from the higher-genus case, the different sectors are connected (they are not for  $g > 1$ ), and the structure of the Chern–Simons moduli spaces is non-Hausdorff and rather complicated. For  $\Lambda=0$  this has been investigated in detail by Louko and Marolf,<sup>13</sup> and the analysis has been extended to  $\Lambda \neq 0$  by Ezawa.<sup>14</sup> The former authors have also suggested a unified quantum theory, in which all sectors are quantized together, and which contains operators that map in between the different sectors. From now on, when talking about moduli spaces, we will always mean the appropriate geometrodynamical sector.

This in turn raises a question on prospective “gauge-theoretic” quantum theories for 2+1 gravity: which feature of the quantum representation reflects the fact that the correct, metric sector is described? More generally, what indicates that the underlying gauge group is  $SO(3, 1)$ ,  $SO(2, 2)$ , rather than the compact group  $SO(4)$ , say? The noncompactness of the gauge group is responsible for a number of subtleties that occur in the existing quantization programs. The two main lines of research are that of Nelson, Regge, and Zertuche on the one hand<sup>15–20</sup> and Ashtekar *et al.*, Smolin,

Marolf, and Loll on the other.<sup>3,21–23</sup> The former are mostly (but not only) concerned with the case  $\Lambda < 0$  and a quantization based on a finite subalgebra of Wilson loops, whereas the latter treat the case  $\Lambda = 0$ , and aim at a rigorous realization of the loop quantization program, first proposed by Rovelli and Smolin for 3+1 gravity.<sup>24</sup> In both cases, explicit descriptions of the quantum theory are available for  $g=1$  and (incompletely) for  $g=2$ . We will describe their main results in turn.

### A. Vanishing cosmological constant

Let us begin with the case of a vanishing cosmological constant,  $\Lambda=0$ . As explained above, it is special because the reduced phase space is of the form of a cotangent bundle  $T^*(\mathcal{A}^F/G)$  over the space of flat  $SO(2, 1)$ -connections modulo gauge. We will from now on work with the defining two-dimensional representation of  $PSU(1, 1) = SU(1, 1)/\mathbb{Z}_2$ , which has been used in most applications. (We in any case are glossing over differences that arise between using  $G$  and some covering group  $\tilde{G}$ .) In close analogy with the 3+1 theory, one introduces as a convenient (over)complete set of phase-space variables the (normalized) Wilson loops<sup>3</sup>

$$T^0(\gamma)[\omega] = \frac{1}{2} \text{Tr} U_\gamma, \quad T^1(\gamma)[\omega, E] = \oint_\gamma d\gamma^i \epsilon_{ij} \text{Tr}(E^j U_\gamma), \quad (2.8)$$

with  $\gamma \in \pi_1(\Sigma^g)$ , and where the canonical pairs  $(\omega, E)$  by slight abuse of notation denote now the coordinates of  $T^*\mathcal{A}^F$ .

The generalized Wilson loops  $T^I$ ,  $I=0,1$ , form a closed Poisson algebra with respect to the canonical structure induced on  $T^*\mathcal{F}(\Sigma^g) \equiv T^*(\mathcal{A}^F/PSU(1, 1))$  from (2.4), given by

$$\begin{aligned} \{T^0(\alpha), T^0(\beta)\} &= 0, \\ \{T^0(\alpha), T^1(\beta)\} &= -\frac{1}{2} \sum_n \Delta_n(\alpha, \beta) (T^0(\alpha \circ_n \beta) - T^0(\alpha \circ_n \beta^{-1})), \\ \{T^1(\alpha), T^1(\beta)\} &= -\frac{1}{2} \sum_n \Delta_n(\alpha, \beta) (T^1(\alpha \circ_n \beta) - T^1(\alpha \circ_n \beta^{-1})), \end{aligned} \quad (2.9)$$

where the sums are over all intersection points  $n$  of the homotopy elements  $\alpha$  and  $\beta$ , with  $\Delta_n(\alpha, \beta) = 1 (= -1)$  if the two tangent vectors  $(\dot{\alpha}, \dot{\beta})$  form a right- (left-)handed zweibein at  $n$ . A similar algebra is derived by Nelson and Regge in Ref. 15. The algebraic Mandelstam constraints mentioned earlier take the form

$$\begin{aligned} T^0(\alpha)T^0(\beta) &= \frac{1}{2}(T^0(\alpha \circ \beta) + T^0(\alpha \circ \beta^{-1})), \\ T^0(\alpha)T^1(\beta) + T^0(\beta)T^1(\alpha) &= \frac{1}{2}(T^1(\alpha \circ \beta) + T^1(\alpha \circ \beta^{-1})) \end{aligned} \quad (2.10)$$

for pairs of intersecting homotopy elements  $\alpha$  and  $\beta$ . Before discussing quantum theories based on the algebraic relations (2.9) and (2.10), we will present an even simpler quantization, defined directly on the reduced phase space  $T^*\mathcal{F}(\Sigma^g)$ , the cotangent bundle over Teichmüller space. Since Teichmüller space is diffeomorphic to the real space  $\mathbb{R}^{6g-6}$ , one can simply choose a set of global coordinates  $x_i$ ,  $i=1 \cdots 6g-6$ , and a corresponding set of momenta  $p_i$ , and quantize à la Schrödinger on  $L^2(\mathbb{R}^{6g-6}, dx)$ . This is straightforward, but not terribly helpful as long as one does not establish an explicit relation between the coordinates  $x$  and the observables introduced earlier. At this point we need a mathematical result by Okai,<sup>25</sup> who established an explicit cross section of the bundle of  $PSU(1, 1)$ -valued holonomies, the space of homomorphisms  $\text{Hom}(\pi_1(\Sigma^g), PSU(1, 1))$  over Teichmüller space. As coordinates on  $\mathcal{F}(\Sigma^g)$  he uses the Fenchel–Nielsen parameters,<sup>26</sup> a set  $(l_i, \tau_i)$ ,  $i=1 \cdots 3g-3$ , of length and angle coordinates associated

with a pants decomposition of the genus- $g$  surface. Using this result, one may write arbitrary Wilson loops as functions  $T(\gamma)[l, \tau]$  of the Fenchel–Nielsen coordinates, and thus make contact with the previous formulation. It also enables one to find an explicit set of independent Wilson loop coordinates on Teichmüller space, i.e., to solve the overcompleteness problem.<sup>21</sup>

If one regards this description of the reduced phase space as the basic one, and the Wilson loop variables as derived quantities, one may still want to represent the Wilson loops as well-defined operators on the Hilbert space  $L^2(\mathcal{F}(\Sigma^g), dl d\tau)$ , and realize some or all of the algebraic relations (2.9) and (2.10) in the quantum theory. Alternatively, one may take the Wilson loop observables as fundamental physical quantities, and try to find self-adjoint representations of the  $T^0$  and  $T^1$ , such that their commutator algebra is isomorphic to the classical Poisson brackets (2.9). *A priori* one expects this latter quantization ansatz to be more general, since it involves the entire representation theory [not just that on  $L^2(\mathcal{F}(\Sigma^g), dl d\tau)$ ] of a rather complicated Poisson algebra, whereas the former is expected to be essentially unique, up to possible factor orderings for the  $\hat{T}^0$  and  $\hat{T}^1$  in terms of the fundamental operators  $\hat{l}$ ,  $\hat{\tau}$ ,  $\hat{p}_l$ , and  $\hat{p}_\tau$ .

A similar issue arises in 3+1 quantum general relativity, where it has been suggested to use a quantization based on  $SL(2, \mathbb{C})$ -Wilson loops.<sup>24</sup> The variables in this case are of course field theoretic and the analog of the loop algebra (2.9) requires a regularization. The question is whether one should abstractly study the representation theory of this loop algebra or consider only special representations that can be obtained (formally) through an integral transform from the connection representation, where  $SL(2, \mathbb{C})$ -connections modulo gauge  $A \in \mathcal{AS}$  are regarded as fundamental. This so-called loop transform<sup>24</sup> has the form

$$\psi(\gamma) := \int_{\mathcal{AS}} d\mu(A) T^0(\gamma)[A] \Psi(A), \quad (2.11)$$

where the Wilson loop functionals  $T^0$  plays the role of an integral kernel, and wave functions  $\psi(\gamma)$  in the loop representation are labeled by spatial closed curves  $\gamma$ . The idea is that once the loop transform has been defined rigorously, one obtains a loop representation that is unitarily equivalent to the connection representation. In contrast with 3+1 dimensions, in 2+1 dimensions this construction can be carried out explicitly. This is particularly useful since in practice it turns out to be difficult to abstractly construct irreducible representations of (2.9), with the operators simultaneously satisfying quantum analogs of the constraints (2.10), constraints arising from (2.6), and other conditions like  $T^0 \geq 1$ .

An early implementation of these ideas can be found in a series of papers,<sup>3</sup> where a rigorous quantum loop representation for  $\Sigma = T^2$  is constructed; however, only the compact sector [where all Wilson loops are bounded by  $|T^0(\gamma)| \leq 1$ ] is quantized. As discussed in detail by Marolf,<sup>22</sup> this construction of the loop representation via the loop transform cannot be carried over unmodified to the physical, noncompact sector. The reason for this is readily illustrated by the explicit form of the transform for the torus case (where the analog of the Teichmüller space is  $\mathbb{R}^2/\mathbb{Z}_2$ , parametrized by  $a_1, a_2$ ), given by

$$\psi(\mathbf{n}) = \langle T^0(\mathbf{n}), \Psi \rangle = \int d\mathbf{a} T^0(\mathbf{n}) \Psi(\mathbf{a}). \quad (2.12)$$

Since the homotopy group  $\pi_1(T^2)$  is abelian, its elements can be labeled by a pair  $\mathbf{n}$  of integers. Since  $T^0(\mathbf{n}) = \cosh \mathbf{n} \cdot \mathbf{a}$ , the integrand of (2.12) diverges rapidly for large  $\mathbf{a}$ . As demonstrated in Ref. 22, the kernel of the transform, i.e., those elements mapped to 0, is in fact dense in  $L^2(\mathbb{R}^2, d\mathbf{a})$ . Nevertheless, one may define loop representations that are isomorphic to the connection representation. This involves the choice of a dense subspace of  $L^2(\mathbb{R}^2, d\mathbf{a})$ , satisfying a number of properties, and as a result of the construction in general contains wave functions that cannot be expressed as functions of homotopy classes. Strictly speaking, these are therefore not “loop representations” in the usual sense.

An alternative way of making the loop transform (2.12) well defined was proposed by Ashtekar and Loll.<sup>23</sup> The basic idea is to employ a nontrivial volume element  $dV = e^{-M(\mathbf{a})} d\mathbf{a}$  in the transform that provides a sufficient damping for large  $\mathbf{a}$ , so as to make it converge for general elements of the connection Hilbert space. That this is a viable procedure was demonstrated in Ref. 23, where the loop representation on  $T^2$  was constructed for a particular choice of the damping factor  $M(\mathbf{a})$ . The choice of a suitable measure is an additional input, and  $M$  has to satisfy a number of conditions in order to make the loop representation well defined. The explicit form of the  $\hat{T}^1$ -operators and their action on loop states  $\psi(\mathbf{n})$  is more complicated than in the case for trivial measure, since it contains a contribution from  $\nabla M$ . Still, by construction this loop representation is isomorphic to the connection representation and, in particular, quantum analogs of (2.9) and (2.10) continue to hold.

Of these two approaches, only the Ashtekar–Loll construction has been extended to the higher-genus case, although not in as much detail as in the torus case. In a first step, let us point out that within the connection representation on  $L^2(\mathcal{S}(\Sigma^g), dl d\tau)$ , one can straightforwardly construct self-adjoint operators  $\hat{T}^0$  and  $\hat{T}^1$  corresponding to (2.8). Using the results of Ref. 25, one can write any Wilson loop  $T^0$  as a function of the Fenchel–Nielsen coordinates  $l_i, \tau_i$ . The corresponding self-adjoint operators act as multiplication operators. To find the momentum operators  $\hat{T}^1$ , one uses the fact that there is a natural symplectic structure on Teichmüller space [although  $\mathcal{S}(\Sigma)$  presently plays the role of a *configuration* space], namely, the Weil–Peterson symplectic form  $\sum_i dl_i \wedge d\tau_i$ , with respect to which the Fenchel–Nielsen coordinates are canonical. As was already discussed in Ref. 3, this structure can be used to obtain an explicit representation for the momentum operators  $\hat{T}^1$ . As an example, consider the Wilson loops of the pair of homotopy generators  $\alpha_1, \beta_1$  for  $g=2$  as functions of the six Fenchel–Nielsen coordinates  $l_{-\infty}, l_0, l_{\infty}, \tau_{-\infty}, \tau_0, \tau_{\infty}$  (see Refs. 21 and 25 for derivation and notation):

$$\hat{T}^0(\alpha_1) = \cosh \frac{l_{-\infty}}{2}, \quad (2.13)$$

$$\hat{T}^0(\beta_1) = \sinh \frac{\tau_{-\infty}}{2} \sinh \frac{\tau_0}{2} + \frac{\cosh l_{-\infty}/2 \cosh l_0/2 + \cosh l_{\infty}/2}{\sinh l_{-\infty}/2 \sinh l_0/2} \cosh \frac{\tau_{-\infty}}{2} \cosh \frac{\tau_0}{2}.$$

The corresponding self-adjoint momentum operators are

$$\hat{T}^1(\alpha_1) = -\frac{i\hbar}{2} \sinh \frac{l_{-\infty}}{2} \frac{\partial}{\partial \tau_{-\infty}}, \quad (2.14)$$

$$\begin{aligned} \hat{T}^1(\beta_1) = & -\frac{i\hbar}{2} \frac{\cosh l_{\infty}/2 \cosh l_0/2 + \cosh l_{-\infty}/2}{\sinh^2 l_{-\infty}/2 \sinh l_0/2} \cosh \frac{\tau_{-\infty}}{2} \cosh \frac{\tau_0}{2} \frac{\partial}{\partial \tau_{-\infty}} \\ & -\frac{i\hbar}{2} \frac{\cosh l_{\infty}/2 \cosh l_{-\infty}/2 + \cosh l_0/2}{\sinh l_{-\infty}/2 \sinh^2 l_0/2} \cosh \frac{\tau_{-\infty}}{2} \cosh \frac{\tau_0}{2} \frac{\partial}{\partial \tau_0} \\ & + \frac{i\hbar}{2} \frac{\cosh \tau_{-\infty}/2 \cosh \tau_0/2}{\sinh l_{-\infty}/2 \sinh l_0/2} \sinh \frac{l_{\infty}}{2} \frac{\partial}{\partial \tau_{\infty}} - \frac{i\hbar}{2} \left( \cosh \frac{\tau_{-\infty}}{2} \sinh \frac{\tau_0}{2} \right. \\ & \left. + \frac{\cosh l_{-\infty}/2 \cosh l_0/2 + \cosh l_{\infty}/2}{\sinh l_{-\infty}/2 \sinh l_0/2} \sinh \frac{\tau_{-\infty}}{2} \cosh \frac{\tau_0}{2} \right) \frac{\partial}{\partial l_{-\infty}} \\ & - \frac{i\hbar}{2} \left( \cosh \frac{\tau_0}{2} \sinh \frac{\tau_{-\infty}}{2} + \frac{\cosh l_{-\infty}/2 \cosh l_0/2 + \cosh l_{\infty}/2}{\sinh l_{-\infty}/2 \sinh l_0/2} \sinh \frac{\tau_0}{2} \cosh \frac{\tau_{-\infty}}{2} \right) \frac{\partial}{\partial l_0}, \end{aligned}$$

where we have chosen a factor ordering with the momenta to the right. The functional form of the Wilson loop operators is considerably more complicated than the corresponding expressions in the torus case. In Ref. 21 it is shown that also for the higher-genus case there exist suitable measures that ensure the convergence of the loop transform for a sufficiently big set of connection wave functions. Thus there are no obvious obstacles to quantizing along the lines proposed in Ref. 23, although the details of these loop representations remain to be worked out.

There is a different treatment by Manojlović and Miković in the connection formulation,<sup>27</sup> which is not based on the classical reduction to the reduced phase space, but instead relies on a quantum reduction à la Dirac. For a nonvanishing spatial determinant  ${}^{(2)}g$ , one may rewrite the action (2.3) in such a way that the functional form of the ensuing first-class constraints is exactly analogous to the ones obtained in the Ashtekar formulation in 3+1 dimensions.<sup>28</sup> In particular, in this form the Hamiltonian constraint is quadratic in the momenta  $E$ . For the torus case, one obtains an effective finite-dimensional theory with three Gauss law constraints and one Hamiltonian constraint. It is argued that the quantum theory is given by unitary irreducible representations with zero mass of the Poincaré algebra in three dimensions. Since the states in these representations depend on two real parameters, this suggests that the reduced configuration space of the system (2.3) is  $\mathbb{R}^2$ , which does not quite agree with the usual result. Probably this can be traced to a subtlety in the solution to the Gauss law constraints, which may be given in terms of wave functions of three rotationally invariant parameters  $a_i$ . These are treated as free parameters in Ref. 27, whereas strictly speaking they are subject to a number of inequalities ( $|a_3| \leq a_1 a_2$ ,  $a_1 \geq 0$ ,  $a_2 \geq 0$ ).

## B. Nonvanishing cosmological constant

Let us now turn to the cases with a nonvanishing cosmological constant. As discussed earlier, their physical phase spaces too are given by spaces of flat connections modulo gauge. One may therefore again describe them as suitably regular spaces of homomorphisms of  $\pi_1(\Sigma^g)$  into the gauge groups  $G = \text{SO}(2, 2)$  and  $\text{SO}(3, 1)$ , for  $\Lambda < 0$  and  $\Lambda > 0$ , respectively. However, since those groups do not have a cotangent bundle structure, the holonomies and Wilson loop variables are now necessarily functions on phase space [unlike the Wilson loops  $T^0$  of (2.8), which are functions on configuration space].

In Ref. 16, Nelson, Regge, and Zertuche compute the path-dependent Poisson algebras for the  $G$ -valued phase-space holonomies and, after going to the spinor representations  $\text{SL}(2, \mathbb{R}) \times \text{SL}(2, \mathbb{R})$  and  $\text{SL}(2, \mathbb{C})$ , respectively, the Poisson algebra of the corresponding Wilson loops  $T(\gamma)[A] = \frac{1}{2} \text{Tr } U_\gamma$ . In the former case, one gets two copies  $\{T^+(\gamma)\}$  and  $\{T^-(\gamma)\}$  of  $\text{SL}(2, \mathbb{R})$ -Wilson loops satisfying the Poisson algebra [cf. (2.9)]

$$\begin{aligned} \{T^\pm(\alpha), T^\pm(\beta)\} &= \pm \frac{\Delta}{4} \sqrt{-\Lambda} (T^\pm(\alpha\beta) - T^\pm(\alpha\beta^{-1})), \\ \{T^+(\alpha), T^-(\beta)\} &= 0 \end{aligned} \tag{2.15}$$

for pairs of homotopy elements  $\alpha, \beta$  with a single intersection. For  $\Lambda > 0$ , a similar decomposition is only possible over the complex numbers, and the factor on the nonvanishing right-hand side in (2.15) has to be replaced by  $i\Delta/4\sqrt{\Lambda}$ , which is purely imaginary. They go on to study the representation theory of the “plus sector”  $\{T^+(\gamma)\}$  of the algebra (2.15), restricted to a single “handle,” i.e., to the subgroup of  $\pi_1(\Sigma^g)$  generated by a single pair of generators  $\alpha_i$  and  $\beta_i$ . The motivation for this is the hope that the full quantum theory of a genus- $g$  surface may be obtained by combining several such copies appropriately, although a concrete construction to our knowledge has not yet been given. In any case, the quantization for one handle they propose is based on a

(non-Lie) algebra of a rescaled version of the operators  $\hat{T}^+(\alpha)$ ,  $\hat{T}^+(\beta)$ , and  $\hat{T}^+(\alpha\beta)$ . The physically relevant quantum representations are those where the basic operators are unbounded, and Ref. 16 contains a preliminary discussion of some of their properties.

This investigation for  $\Lambda < 0$  is extended to the case  $g=2$  by Nelson and Regge in Refs. 17–20, where it is proposed to base the quantization on a ring  $\mathcal{R}$  of polynomials of a highly symmetric subset of 15 Wilson loop variables  $\{T^+(\gamma)\}$ .  $\mathcal{R}$  is closed under Poisson brackets and the subset is chosen so that any other traced holonomy may be expressed as a function of this subset via the Mandelstam identities [the first set of relations in (2.10)]. To eliminate the remaining overcompleteness of these observables in the classical theory, they propose a quotient construction in which the physical observables are elements of  $\mathcal{R}/I(\mathcal{R})$ , where  $I(\mathcal{R})$  is an ideal (closed under Poisson brackets), generated by the Mandelstam constraints (here called “rank identities”) and constraints coming from the fundamental relation (2.6), called “trace identities.” The difficulty lies in finding an appropriate basis for this ideal, which should consist of six rank and three trace identities. The ideals for  $g=1,2$  are described in Ref. 19, where it is also argued that a similar quotient space construction should be applied to the higher-genus case. [Note that the results in Ref. 21 may be used to explicitly parametrize the quotient spaces  $\mathcal{R}/I(\mathcal{R})$ .]

An issue we have not touched upon so far is the role of the large diffeomorphisms  $\text{Diff } \Sigma / \text{Diff}_0 \Sigma$ , i.e., those that do not lie in the component  $\text{Diff}_0 \Sigma$  connected to the identity. They form the so-called mapping class group, also called the Teichmüller modular group, whose generators are the Dehn twists. The question is whether one should regard them as gauge degrees of freedom, to be factored out like the connected diffeomorphisms. The canonical Dirac treatment of constraints only requires invariance under the action of the connected component of a gauge group. For 2+1 gravity, there is a whole spectrum of proposals on how the large diffeomorphisms should be treated in both the classical and the quantum theory, which goes from ignoring them altogether over implementing them as unitary symmetries to requiring strict invariance, even in the quantum theory (see also the discussion in Ref. 29). In principle such a controversy should be settled by physical arguments, but this presents a problem for a theory like three-dimensional gravity which is largely unphysical. We therefore do not expect that this issue has a definite resolution, and what remains to be understood is which approaches to the large diffeomorphisms are feasible in practice.

Note that none of the Wilson loop variables introduced so far are invariant with respect to large diffeomorphisms. However, they carry (more or less complicated) actions of the mapping class group. For  $g=2$ , Nelson and Regge investigate a canonical (nonlinear) action of the Dehn twists on the algebra  $\mathcal{R}$  introduced above.<sup>17</sup> They also study the center of this algebra with respect to the Dehn twists, i.e., its invariant elements. For genus  $g$ , they find  $g+1$  such central elements, out of which two remain linearly independent once the rank identities are taken into account.<sup>19</sup> However, it is not explained whether or how this construction intertwines with the quotient construction of  $\mathcal{R}/I(\mathcal{R})$  to yield classical observables that are invariant under the mapping class group.

It turns out to be rather nontrivial to quantize the algebraic structure of the classical algebra  $\mathcal{R}$  of Wilson loop variables, and the quotient construction for the physical observables. A quantum analog of the classical Poisson algebra of the polynomials in the 15 chosen  $T^+$ 's is given in Ref. 18. Since the algebra is polynomial, this involves a particular choice of operator ordering. The commutation relations involve a complex constant  $K$  that depends on the cosmological constant  $\Lambda$  and goes to 1 as  $\hbar \rightarrow 0$ . In this limit, the classical Poisson brackets are recovered by substituting

$$\frac{1}{K-1} [\hat{T}^+(\gamma_1), \hat{T}^+(\gamma_2)] \rightarrow \{T^+(\gamma_1), T^+(\gamma_2)\}. \quad (2.16)$$

Similarly, a  $K$ -dependent quantum action of the Dehn twists on the quantized algebra  $\hat{\mathcal{R}}$  can be defined. This framework seems suggestive of a quantum theory defined on a Hilbert space  $L^2(\mathbb{R}^{15})$ ,

where the quantum counterparts of the classical constraints remain to be imposed to project out the physical wave functions. However, it seems to be difficult to carry out this program explicitly, as well as to implement modular invariance at the quantum level. These observations are in line with remarks made earlier in the context of the  $(\Lambda=0)$  case. It therefore may not come as a total surprise that in their most recent paper on the  $g=2, \Lambda<0$  quantum gravity, Nelson/Regge propose a quantization based on a reduced set of six variables, three angles  $\varphi_a$ , and conjugate momenta  $p_a, a=1\cdots 3$ , in terms of which all of the 15 Wilson loop variables can be expressed.<sup>20</sup> [Recall that the dimension of the physical phase space for  $g=2$  is 12, and that we have split it into two  $SL(2, \mathbb{R})$ -sectors.] The quantum operators  $\hat{T}^+(\alpha)$  are functions of the basic operators  $\hat{\varphi}_a$  and  $\hat{p}_a$ , and depend on a complex parameter  $\Theta$ , where  $e^{i\Theta}=K$  (and the classical limit therefore corresponds to  $\Theta=0$ ). For example, the form one finds for the classical Wilson loops of the fundamental homotopy generators  $\alpha_1$  and  $\beta_1$  is

$$T^+(\alpha_1) = \frac{\cos \varphi_2}{\cos \Theta/2}, \tag{2.17}$$

$$T^+(\beta_1) = \frac{1}{2 \cos \Theta/2} \sum_{n,m=\pm 1} \times \frac{\sin(\Theta/4 + (n\varphi_1 + m\varphi_2 + \varphi_3)/2) \sin(\Theta/4 + (n\varphi_1 + m\varphi_2 - \varphi_3)/2)}{\sin n\varphi_1 \sin m\varphi_2} \times e^{-i(3/2n\varphi_1 + 3/2m\varphi_2 + \Theta(np_1 + mp_2))}.$$

One observes that, in contrast with (2.14), the conjugate momentum operators  $\hat{p}_a = -i\partial/\partial\varphi_a$  will not appear linearly but exponentially in the corresponding quantum operator  $\hat{T}^+(\beta_1)$ . In this “reduced phase-space quantization,” the trace and rank identities are fulfilled both classically and quantum mechanically. What is slightly puzzling about this approach is the fact that one seems to end up with three free canonical coordinate pairs, although one knows that the underlying moduli space of  $SO(2, 2)$ -[or  $SL(2, \mathbb{R}) \times SL(2, \mathbb{R})$ -] connections is not a cotangent bundle.

The Hilbert space proposed in Ref. 20 is an  $L^2$ -space on a suitable domain  $D^3 \subset \mathbb{R}^3$  based on three real parameters  $z_a = \cos \varphi_a$ . The inner product on this Hilbert space is to be determined by requiring the basic operators to be self-adjoint. The physically interesting case is the one where the angles  $\varphi_a$  are imaginary and therefore  $z_a \geq 1$ . An appropriate scalar product for this case remains to be found.

Another quantization method for  $\Lambda < 0$  has been suggested by Ezawa,<sup>14</sup> who aims at constructing a unified quantization for all sectors of the reduced Chern–Simons moduli space, not just that corresponding to the geometrodynamical solutions. He adopts a “brute-force” approach (i.e., without any physical justification) to make this space into a cotangent bundle, to which then a geometric quantization procedure (in the sense of Kostant and Souriau) may be applied.

### C. Summary

This concludes our discussion of quantizations in the gauge-theoretic approach to 2+1 gravity. The most promising approaches seem to be those that are closest to a Schrödinger-type quantization on the reduced physical phase space. For  $\Lambda=0$ , the construction of a well-defined quantum theory is straightforward, and in addition one can show that the physically interesting Wilson loop observables can be defined as self-adjoint operators in this representation. The existence of appropriate transforms ensures that there are well-defined loop representations, whose wave functions are labeled by homotopy classes. This shows that the noncompactness of the gauge

group does not present any problems in principle to constructing such representations. However, since they are by construction unitarily equivalent to the reduced phase-space quantization, they do not yield *a priori* any additional physical information.

The situation for  $\Lambda \neq 0$  is not as straightforward, since the lack of a cotangent bundle structure prevents an analogous construction of the reduced phase-space quantization. Given such a quantum theory, for example, as the result of a geometric quantization based on a complex polarization on the reduced phase space,<sup>1,10</sup> one could again attempt to define the Wilson loop operators on its Hilbert space. As we have seen, the existing quantization proposals for  $\Lambda \neq 0$  are based on algebras of Wilson loops, and the details beyond genus-1 become rather involved. Furthermore, it is in general difficult to find quantum analogs of the constraints satisfied by the classical Wilson loops. Although the most recent proposal of Nelson and Regge for  $g=2$ ,  $\Lambda < 0$  is based on wave functions of three variables (and therefore looks like a “reduced quantization”),<sup>20</sup> it is not clear how these relate to an explicit parametrization of the reduced phase space.

Incorporating invariance under the mapping class group seems problematic in all of the connection approaches, because the natural classical observables, the Wilson loops, are not modular invariant. If we regard the large diffeomorphisms as gauge degrees of freedom, the only true classical observables are the Casimir operators of Nelson/Regge, which are far from forming a complete set. This has to do with the fact that the action of the Dehn twists on the reduced phase spaces is rather complicated, so that even in the torus case there is no simply defined “fundamental region” for the modular action. As discussed by Peldàn,<sup>29</sup> this also leads to problems if one tries to find finite-dimensional representations of the modular group on connection wave functions (invariance being just a special case).

### III. QUANTIZATION IN THE GEOMETRIC FORMULATION

The classical starting point for a quantization within the geometric formulation is the Lagrangian (1.1). Also, in this description one can reduce the degrees of freedom to a finite number, as was already noted by Martinec in 1984,<sup>30</sup> and later rediscovered by Hosoya and Nakao<sup>31</sup> and Moncrief.<sup>32</sup> We will discuss in the following the case where  $\Lambda=0$ , although most of the classical treatment is readily extended to include a cosmological constant.<sup>11,33</sup> For general genus, the three constraints  $\mathcal{H}=0$  and  $\mathcal{H}_i=0$  can be decoupled and subsequently solved if one adopts the York time gauge, where  $\Sigma^g$  is taken to be a constant mean curvature surface,<sup>31,32</sup> i.e., the “time” parameter  $\tau$  is given by

$$\tau = \frac{1}{\sqrt{(2)g}} g_{ij} \pi^{ij}. \quad (3.1)$$

For  $g > 1$ , the two-metric  $\gamma_{ij}$  can be uniquely decomposed into a conformal factor  $e^{2\lambda(x)}$  and a metric  $h_{ij}$  of constant scalar curvature  $-1$ . The space  $\mathcal{M}_{-1}$  of such metrics is infinite dimensional, but one may show that the quotient by the diffeomorphisms,  $\mathcal{M}_{-1}/\text{Diff}_0$ , is well defined and diffeomorphic to the Teichmüller space  $\mathcal{T}(\Sigma^g)$ .<sup>32</sup> The physical phase space is given by the cotangent bundle  $T^*\mathcal{T}(\Sigma^g)$ , with global coordinates  $(m_a, p^a)$ ,  $a = 1 \cdots 6g - 6$ . The Hamiltonian constraint  $\mathcal{H}=0$  determines the conformal factor  $\lambda$  as a  $\tau$ -dependent function on this cotangent bundle. The action (1.1) in the reduced variables becomes

$$\int d\tau \left( p^a \frac{dm_a}{d\tau} - H(m, p, \tau) \right), \quad (3.2)$$

where  $H = \int_{\Sigma^g} d^2x \sqrt{(2)g} e^{2\lambda(m, p, \tau)}$  is the Hamiltonian associated with the York time slicing. This Hamiltonian measures the area of the spatial surfaces of constant  $\tau$ , and generates the time evolution on  $T^*\mathcal{T}(\Sigma^g)$ . Unfortunately, the solution for  $\lambda$  is known only implicitly, as a solution of a differential equation,<sup>32</sup> which is problematic since  $H$  depends on it explicitly. Consequently, the

classical and quantum theories have been studied in detail only for the torus case, where the Hamiltonian  $H$  is known as a function of the basic variables. We will concentrate on this case in the following, starting with  $\Lambda=0$ . There are two canonical pairs with  $\{m_a, p^b\} = \delta_a^b$  and the Hamiltonian is given by

$$H(m, p, \tau) = \frac{1}{\tau} \sqrt{m_2^2((p^1)^2 + (p^2)^2)}. \tag{3.3}$$

Alternatively, if one wants to get rid of the square root, one may employ a different gauge condition. For example, Martinec chooses the two-metric to be spatially constant,  $g_{ij}(x, t) = g_{ij}(t)$ ,<sup>30</sup> and Hosoya/Nakao the lapse function to be spatially constant,  $N(x, t) = N(t)$ .<sup>31</sup> This leads to a reduced action of the form

$$S = \int dt \left( p^a \frac{dm_a}{dt} + \tau \frac{dv}{dt} - N \mathcal{H}(m, p, v, \tau) \right), \tag{3.4}$$

with a Hamiltonian constraint  $\mathcal{H}$  quadratic in both the momenta  $p^a$  and the variable  $\tau$  canonically conjugate to the volume variable  $v$ . Going to the quantum theory, in the first approach one looks for solutions of the Schrödinger equation

$$i \frac{\partial \psi}{\partial \tau} = \hat{H} \psi(m, \tau), \tag{3.5}$$

whereas in the second one tries to solve the quantum Hamiltonian constraint  $\hat{\mathcal{H}}\psi=0$ , which takes the form of a Klein–Gordon equation. The latter form is more convenient because of the absence of the square root. As wave functions one may take either the “volume representation” on states  $\psi(m, v)$  or the “time representation” on  $\psi(m, \tau)$ . Martinec chooses the former since he is also interested in the case with a nonvanishing cosmological constant (for which the Hamiltonian constraint contains a term proportional to  $\Lambda v$ ), and gives the general form of the solution.<sup>30</sup> Hosoya and Nakao<sup>34</sup> impose the Hamiltonian constraint on wave functions  $\psi(m, s)$ , where  $s = \ln v$ , but in addition insist that physical states should be invariant under large diffeomorphisms, which in the torus case are elements of the group  $SL(2, \mathbb{Z})$ . They propose to superimpose (noninvariant) solutions to  $\hat{\mathcal{H}}\psi=0$  in order to arrive at  $SL(2, \mathbb{Z})$ -invariant wave functions. However, this construction remains somewhat implicit, because it involves Maass forms, which have been the subject of much study, but are not all known explicitly as functions of the modular parameters  $m_a$ . [The Maass forms are the modular-invariant eigenfunctions of the Laplacian on Teichmüller space which come from the quantization of the term under the square root in (3.3). Its spectrum has a continuous and a discrete part and Puzio has recently suggested that both should be taken into account when constructing physical wave functions.<sup>35</sup>]

These ideas are taken up by Carlip,<sup>36</sup> who argues that (with an appropriate operator ordering) the Schrödinger equation (3.5) should be regarded as the positive square root of the Wheeler–DeWitt equation  $\hat{\mathcal{H}}\psi(m, \tau)=0$ . (How one may rigorously make sense of the square-root operation at the quantum level is discussed in Ref. 35.) He also attempts to define a Hilbert space, i.e., an inner product on the solution space, an issue not considered by previous authors. He proposes  $\int d^2m/m_2^2$  as a modular-invariant inner product, where the integration domain is taken to be a fundamental region for the modular group in the upper half-plane for the complex variable  $m := m_1 + im_2$ . The  $\tau$  in  $\psi(m, \tau)$  is therefore considered as an external parameter. With this scalar product, the momentum operators  $\hat{p}^a = -i\partial/\partial m_a$  are not self-adjoint, although the Laplacian appearing in the Hamiltonian is. This leads one to consider a different representation for the basic operator  $\hat{p}_2$ , namely,  $\hat{p}_2 = -i\partial/\partial m_2 + i/m_2$ , which makes it self-adjoint, provided the wave functions obey appropriate falloff conditions at the integration boundaries for  $m_1$  and  $m_2$ . In order to reproduce the correct modular transformation behavior for the operators in the connection formu-

lation (formally related to the  $\hat{m}_i, \hat{p}_i$ ; see below), the wave functions  $\psi$  then have to transform as forms of modular weight 1/2. This construction is extended in Ref. 37 to representations on spaces of weight functions of arbitrary modular weight, and leads to a one-parameter family of quantum Hamiltonians, which according to Refs. 12 and 37 give rise to inequivalent quantum theories.

The quantum equivalence between the reduced connection formulation and the reduced geometric formulation for  $\Sigma=T^2$  is considered in Ref. 38. The aim is to construct a quantum analog of the time-dependent canonical transformation between the corresponding classical theories<sup>38</sup> (see also Ref. 39). However, since the canonical transformation contains inverses of the momenta, the corresponding operators are only formally defined. Still, a formal operator problem arises, and it is shown that this ambiguity is reduced when demanding the quantum operators to have the same modular transformation behavior as their classical counterparts. (An analogous procedure is used in Ref. 12 to relate the Nelson–Regge torus quantum theory for  $\Lambda < 0$  with the ADM quantization.) A similar line of argument is followed by Anderson,<sup>40</sup> who defines another quantum canonical transformation between the reduced holonomy and the geometric approach. This is defined rigorously, but requires a peculiar “operator-valued measure density” in the definition of the scalar product for the connection representation, which is different from  $\int da$ , but likewise modular invariant. The author alludes to the fact that the use of such measure densities leads to a new ambiguity in the quantization. It would be interesting to understand the physical significance of this generalized quantum structure.

For completeness, let us mention that there is a proposal to construct modular-invariant wave functions in the connection representation for the torus via an integral transform from the geometric approach in terms of Maass form wave functions.<sup>36,37</sup> However, as has been pointed out elsewhere,<sup>13,29</sup> this construction is rather subtle and not well defined in the form proposed there. This presumably affects also the analogous quantum constructions for  $\Lambda \neq 0$  in Ref. 11.

Carlip considers yet another quantization approach, in the form of a more complicated Wheeler–DeWitt equation, obtained when no time slicing is imposed classically, i.e., the function  $\lambda$  in the conformal factor  $e^{2\lambda}$  is left arbitrary.<sup>41</sup> As a result, the Wheeler–DeWitt equation is a nonlocal equation on wave functions  $\Psi(m, \lambda)$ , containing functional derivatives with respect to  $\lambda(x)$ , and too complicated to be solved directly. To enable comparison with the gauge-fixed ADM wave functions, he introduces a formal functional Fourier transform  $\Psi(m, \lambda) \mapsto \tilde{\Psi}(m, \tau)$ , but the resulting equation for  $\tilde{\Psi}(m, \tau)$  is not in any obvious way equivalent to the ones discussed in the fully reduced ADM formulation. Attempting to gauge-fix à la Faddeev/Popov to obtain a scalar product for wave functions in the York time gauge leads to a highly complicated, operator-valued Faddeev–Popov determinant, whose structure is known only perturbatively.

A similar quantization is suggested by Visser in Ref. 42, who also formulates the Wheeler–DeWitt equation on the superspace of the infinite-dimensional space of conformal factors times the finite-dimensional moduli space. Since he freezes out the spatial diffeomorphisms, he obtains an equation without any nonlocal terms (which in principle should be present<sup>43</sup>). It is argued that the Hamiltonian constraint splits into independent constraints on the conformal mode and the modular parameters, and that the latter should contain, besides the Laplacian, also a term proportional to the Ricci scalar on moduli space; however, no solutions are given. This ends our overview of the geometric formulations of 2+1 quantum gravity.

#### IV. CONCLUSIONS

Let us try to draw some conclusions from the quantization attempts described above. From the point of view of the generic higher-genus case, the quantization on the reduced connection phase space—at least for  $\Lambda=0$ —is the one furthest developed and potentially most promising, and allows the Wilson loop observables to be defined as well-defined quantum operators. On the other hand, looking for abstract representations of algebras of Wilson loops not based on a Schrödinger-type quantization of the reduced phase space seems to be much harder. For the case of  $\Lambda \neq 0$ , further illumination is needed of how the noncotangent bundle structure of the moduli spaces is

reflected in the quantum theory. In fact, the case  $\Lambda > 0$  has hardly been explored (see, however, the discussion in Ref. 10, which contains some suggestions of how the gravitational Hilbert space may be related to that of a Chern–Simons theory with complex gauge group).

If one is not content with such a quantization of “frozen dynamics” in terms of time-independent constants of motion, one has to consider an ADM-type quantization, which at this moment does not seem feasible beyond  $g=1$ . In this approach—although *a priori* undesirable—one is in practice restricted to a particular choice of gauge-fixing, the York time gauge. Quantization of this gauge degree of freedom poses difficulties that so far have not been overcome. In the ADM treatment of the torus case, in contrast with the connection formulation, incorporation of modular invariance does not present any obvious problems, although the resulting Hilbert spaces are not known in a very explicit way.

Since in the best-explored case of  $\Sigma=T^2$ , the relation between the classical formulations in terms of Teichmüller parameters and holonomies, and reduced ADM variables is well understood,<sup>38,39</sup> it is natural to search for a corresponding relation in the quantum theory. It is relatively easy to establish a formal correspondence between operators in the various quantizations (by “putting hats on everything”), but it seems difficult to make these constructions rigorous. This is not particularly surprising, and in essence a consequence of the Groenewold–Van Hove theorem. If one has a quantum theory in which a complete set of basic operators is represented by self-adjoint operators, it is in general *not* possible to represent another quantity, which is classically a nonpolynomial function of those basic variables, as a self-adjoint operator on the same Hilbert space. This makes it hard to relate the quantum theories of the metric and the connection approaches.

This may seem an unsatisfactory state of affairs but, on the other hand, it is well known that not every classical equivalence can be elevated to a quantum equivalence. Moreover, for phase spaces not of the form of an  $\mathbb{R}^{2n}$ , the quantization is typically nonunique, even if one starts from a single classical description. For more physical theories, one may of course decide that one quantum theory rather than another is correct, because it is in better agreement with physical observations, but this road is not available to us in the case of 2+1 gravity.

It has been suggested to resolve the ambiguity in the choice of a time slicing in the ADM quantum theory by declaring the timeless connection quantization as fundamental,<sup>12</sup> which is an interesting idea. From what we have said above it follows that (for  $g=1$ ) one should expect to recover the ADM quantum theory at most in some appropriate perturbative or semiclassical sense. In turn, one may in the same limit try to *define* a (perturbative) quantum theory in the geometric formulation for  $g \geq 2$  via a connection quantization.

When attempting to generalize any of the above conclusions to 3+1 canonical gravity, one should keep in mind that its structural resemblance is greatest with that of 2+1 gravity for  $g \geq 2$ , as explained by Moncrief.<sup>44</sup> Recall that also in 3+1 dimensions one has the choice between a geometric formulation in terms of the four-metric  $g_{\mu\nu}$  and a connection formulation in terms of the Ashtekar connection  $A_{\mu}^a$ , and that also in this case the quantization of the latter has progressed much further than that of the traditional ADM approach. One may therefore again be tempted to regard this approach as fundamental as far as the quantum theory is concerned. However, note that in this case the Wilson loops can be used to label solutions to both the spatial diffeomorphism and Gauss law constraints, but not to the quantum Hamiltonian. The solution to the “problem of time” suggested above for the 2+1 theory cannot therefore be carried over immediately to the (3+1)-dimensional context.

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