PATH INTEGRAL AND SOLUTIONS OF THE CONSTRAINT EQUATIONS:
THE CASE OF REDUCIBLE GAUGE THEORIES.

Rafael Ferraro\textsuperscript{(a,b)}, Marc Henneaux\textsuperscript{(c,d)} and Marcel Puchin\textsuperscript{(e)}

\textsuperscript{(a)}\textit{Instituto de Astronomía y Física del Espacio,}
\textit{C.C.67, Sucursal 28, 1428 Buenos Aires, Argentina,}

\textsuperscript{(b)}\textit{Facultad de Ciencias Exactas y Naturales, Departamento de Física,}
\textit{Universidad de Buenos Aires, Ciudad Universitaria, Pab.I, 1428 Buenos Aires, Argentina.}

\textsuperscript{(c)}\textit{Faculté des Sciences, Université Libre de Bruxelles, Campus Plaine,}
\textit{C.P.231, B-1050 Bruxelles, Belgium.}

\textsuperscript{(d)}\textit{Centro de Estudios Científicos de Santiago, Casilla 16443, Santiago 9, Chile.}

\textsuperscript{(e)}\textit{Universitatea din Craiova, Str.A.I.Cuza 13, Facultatea de Fizica, 1100 Craiova, Romania.}

\section*{ABSTRACT}

It is shown that the BRST path integral for reducible gauge theories, with appropriate boundary conditions on the ghosts, is a solution of the constraint equations. This is done
by relating the BRST path integral to the kernel of the evolution operator projected on the physical subspace.
In the Dirac quantization of gauge systems, the physical states are annihilated by the constraint operators

\[ G_a(\hat{q}^i, \hat{p}_i) \psi = 0. \] (1)

In the coordinate representation, the constraints, which we shall assume for definiteness to be of even Grassmann parity, read

\[ G_a(q^i, \hbar \frac{\partial}{\partial q^i}) \psi(q^i) = 0. \] (2)

The equations (1-2) are Gauss law constraints in Yang-Mills theory; while in Einstein gravity, they are known as the Wheeler-De Witt equations.

The equations (1-2) are hard to solve in practice for realistic theories. For this reason, it has been tried by several authors to build solutions of (1-2) through path integral methods. This approach has been initiated by Hartle and Hawking [1] in the case of quantum gravity, and has been systematically studied in Ref.[2-4]. In particular it has been shown on general grounds that the path integral with appropriate non minimal sector and appropriate boundary conditions for the ghosts and the antighosts is a solution of the constraint equations in the case of irreducible constraints. The purpose of this letter is to extend the results of [2-4] to the case of reducible constraints. As it is well known, reducible constraints should also be imposed as operator equations on the physical states in the Dirac quantization method, yielding as in (2)

\[ G_{a_0}(q^i, \hbar \frac{\partial}{\partial q^i}) \psi(q^i) = 0, \] (3)

where we have replaced the index \( a \) by \( a_0 \) since there will be further indices. Because of the reducibility of the constraints, however, the ghost spectrum is much bigger and, in order to produce solutions of (3) by path integral methods, one must specify the boundary conditions that the new variables should fulfill. This is done in this letter. We shall follow the method of Ref.[4], where the path integral is related to a definite operator expression, namely the “projected kernel” of the evolution operator. It should be stressed that the problem of the boundary conditions on the ghosts and their conjugate variables is non trivial. Indeed, different boundary conditions generically lead to a path integral expression which may be of interest, but which would not solve the constraint equations.

Our starting point is the concept of “projected kernel” of a gauge invariant operator \( A_0(\hat{q}, \hat{p}) \), \( [A_0, G_{a_0}] = 0 \) (Ref.[4], Chap.13). Let \( |\psi_\alpha > \) be a basis of solutions of the constraint equations (3). The scalar product \( < \psi_\alpha | \psi_\beta > \) is generally infinite because it involves an integration over the gauge orbits. The physical scalar product is obtained by inserting a gauge condition. We denote it by \( (\psi_\alpha | \psi_\beta) \) and we assume that the basis \( |\psi_\alpha > \) is orthonormal, \( (\psi_\alpha | \psi_\beta) = \delta_{\alpha\beta} \).

Any gauge invariant operator \( A_0 \) maps the physical subspace on the physical subspace. Its physical matrix elements are \( (\psi_\alpha | A_0 | \psi_\beta) \). The projected kernel of the operator \( A_0 \) in the coordinate representation is defined by

\[ A_0^P(q', q) = \sum_{\alpha, \beta} <q' | \psi_\alpha > (\psi_\alpha | \hat{A}_0 | \psi_\beta) < \psi_\beta | q >, \] (4)
and is manifestly a solution of the constraint equations at \( q' \) and of the complex conjugate at \( q \). The projected kernel contains all the information about the physical matrix elements of \( A_0 \) and fulfills the folding relation

\[
(A_0 B_0)_{P'}(q', q) = \int dq' A_0^{P'}(q'', q) \mu \left( q', \frac{\hbar}{i} \frac{\partial}{\partial q''} \right) B_0^{P'}(q', q),
\]

where \( \mu \) is the insertion needed to regularize the scalar product, \( (\psi_\alpha | \psi_\beta) = <\psi_\alpha | \mu | \psi_\beta > \).

The projected kernel of \( A_0 \) may be related to BRST quantization. To see this, it is necessary to recall some elements of BRST theory. As standard we assume that it has been possible to order the “minimal” BRST charge \( \Omega_{\text{min}} \) in such a way that

\[
(\Omega_{\text{min}}^2) = 0,
\]

and

\[
(\Omega_{\text{min}}^\dagger) = \Omega_{\text{min}}.
\]

The operator \( \Omega_{\text{min}} \) fulfilling (6-7) can be written in \( \eta - P \) order (all ghosts \( \eta^{a_k} \) to the left of their conjugate operators \( P_{a_k} \)) by using repeatedly the ghost (anti)commutation relations, to yield

\[
\Omega = \eta^{a_0} G_{a_0} + \sum_{k=0}^{L-1} \eta^{a_{k+1}} Z_{a_{k+1}}^{a_k} P_{a_k} + \text{“more”},
\]

where “more” contains at least one \( \eta \) and two \( P \)'s, or two \( \eta \)'s and one \( P \). Here, the \( \eta^{a_0} \) are the ghosts of the zeroth generation associated with the constraints, the \( \eta^{a_1} \) are the ghosts of ghosts of the first generation associated with the reducibility identity \( Z_{a_1}^{a_0} G_{a_0} = 0 \), the \( \eta^{a_2} \) are the ghosts of ghosts of the second generation associated with the reducibility identity \( Z_{a_2}^{a_1} Z_{a_1}^{a_0} \approx 0 \), etc. The \( P_{a_k} \) are their conjugate momenta.

As in [4], Section 14.5, we define the constraint operators of the Dirac approach to be the coefficients of the ghost operators in (8) (see also [5, 6]). This is not the only possible definition but it has the following nice properties

(i) In the limit \( \hbar \to 0 \), \( G_{a_0} \) goes over into the original classical constraints.

(ii) \( [G_{a_0}, G_{b_0}] = C_{a_0 b_0}^{c_0} G_{c_0} \) (in that order).

(iii) \( Z_{a_1}^{a_0} G_{a_0} = 0 \) (in that order).

Because of (ii) and (iii), the equations \( G_{a_0} | \psi > = 0 \) are quantum-mechanically reducible and fulfill the necessary integrability condition for having solutions (“no anomaly”).

Given a BRST invariant operator \( A \) of ghost number zero,

\[
[A, \Omega] = 0, \quad A^\dagger = A,
\]
we define the gauge invariant operator $A_0(q,p)$ of which $A$ is the BRST invariant extension, as the ghost-independent part of $A$ in its $\eta-P$ ordering,

$$A = A_0 + \eta^a A_a^b P_b + \ldots$$  \hspace{1cm} (10)

It then follows from (9) that $[A_0, G_{a0}] = A_{a0}^b G_b$ (in that order).

Under the unitary transformation generated by

$$C = \frac{1}{2} (\eta^a \epsilon^b_{a0} P_b - P_b \epsilon^b_{a0} \eta^a) + \frac{1}{2} (\eta^a \epsilon^b_{a1} P_{b1} + P_{b1} \epsilon^b_{a1} \eta^a) + \ldots$$  \hspace{1cm} (11)

the constraints and reducibility operators transform as

$$G_{a0} \rightarrow \bar{G}_{a0} = G_{a0} + \epsilon^b_{a0} G_b + \frac{1}{2} [G_{a0}, \epsilon^b_{b0}] - \frac{1}{2} [G_{a0}, \epsilon^b_{b1}] + \ldots$$  \hspace{1cm} (12)

$$Z_{a1} \rightarrow \bar{Z}_{a1} = Z_{a1} - Z_{a1} - \epsilon^b_{a1} Z_{b} + \frac{1}{2} [Z_{a1}, \epsilon^b_{b1}] - \frac{1}{2} [Z_{a1}, \epsilon^b_{b1}] + \ldots$$  \hspace{1cm} (13)

The physical states $|\psi>$ of the Dirac approach, solutions of (3), transform thus as

$$|\psi> \rightarrow |\bar{\psi}> = |\psi> - \frac{1}{2} \epsilon^a_{a0} |\psi> + \frac{1}{2} \epsilon^a_{a1} |\psi> - \ldots$$  \hspace{1cm} (14)

This leaves the physical scalar product unchanged (although $<\psi|\chi>$ may be modified), because the regulator $\mu$ in $(\psi|\chi) = <\psi|\mu|\chi>$ transforms as $\mu \rightarrow \bar{\mu} = \mu + \frac{1}{2} (\epsilon \mu + \mu \epsilon)$ with $\epsilon = \epsilon_{a0} - \epsilon_{a1} + \ldots$ \[4\]. In Ref.\[7\], it has been shown that the above definition of quantum constraints and transformation laws are quite natural for constraints linear in the momenta, and amount to regarding the wave functions as densities of weight 1/2 under changes of coordinates in $(q,\eta)$ space.

The physical states of the BRST method are the states annihilated by the BRST charge. It is easy to verify that the solutions of the constraints (3), tensored by the ghost state annihilated by all the ghost momenta is a solution of $\Omega^{min}|\Psi> = 0$

$$|\Psi_\alpha> = |\psi_\alpha> \otimes |\chi> \quad G_{a0} |\psi_\alpha> = 0 \quad P_{a1} |\chi> = 0 \quad \Rightarrow \quad \Omega^{min}|\Psi_\alpha> = 0.$$  \hspace{1cm} (15)

Furthermore, these states cannot be BRST-exact since they do not involve the ghosts. For each solution of (3) there is thus one and only one equivalence class of BRST physical states.

The states (15) are convenient for making the comparison between the Dirac and BRST quantization methods. However, it is as difficult to construct them as it is difficult to find the physical states of the Dirac method, since in both cases one needs to solve
explicitly the constraint equations (3). For this reason it is useful to replace the states (15) by states that are dual to them, namely the states annihilated by the ghosts themselves,

$$\eta^{ak}|\Psi> = 0 \quad \Rightarrow \quad \Omega^{\text{min}}|\Psi> = 0.$$ (16)

These states are solutions of the BRST physical state condition no matter what their q-dependence is, since each term in the BRST charge $\Omega^{\text{min}}$ contains one more $\eta$ than there are $P$’s and therefore annihilates $|\Psi>$. They are in general not independent in cohomology, however, since some of them are BRST exact. The cohomological classes that they define are actually dual to the cohomological classes defined by (15) (Ref.[4],§14.5.3; see also [8,J9]). So, one may say that while the constraints $G_{a0}|\psi> = 0$ are directly enforced for the states (15), they are indirectly enforced for the states (16), through the passage to the BRST cohomology.

The states (16) do not have, in general, ghost number zero, $\mathcal{G}|\Psi> \neq 0$. This leads to ill-defined scalar products. For this reason, it is convenient to extend further the ghost spectrum, by introducing new variables, and to tensor the states (16) with appropriate states of the new sector, in such a way that the product states have total ghost number zero [4, 8, 9]. The new variables are called ”non minimal” variables.

In the reducible case, the non minimal variables are given by blocks $[ (\lambda^s_{ka}, b^k_{sa}), (\bar{C}^s_{ka}, \rho^k_{sa}) ]$, $0 \leq k \leq L$, $k \leq s \leq L$, where $\lambda^0_{ka} \equiv \lambda^a_0$ are the Lagrange multipliers for the constraints and

$$[b^k_{s'}b^s, \lambda^s_{ka}] = -\delta^k_{s'} \delta^s_s \delta^a_{a'}, = [\rho^s_{a'} \bar{C}^s_{ka}, \bar{C}^s_{ka'}],$$ (17)

$$\epsilon(b^k_{sa}) = \epsilon(\lambda^s_{ka}) = s - k \pmod{2},$$ (18)

$$\epsilon(\rho^s_{ka}) = \epsilon(\bar{C}^s_{ka}) = s - k + 1 \pmod{2},$$ (19)

$$gh(b^k_{sa}) = k - s = -gh(\lambda^s_{ka}),$$ (20)

$$gh(\rho^k_{sa}) = -k + s + 1 = -gh(\bar{C}^s_{ka}),$$ (21)

(see Ref.[10], and [4] Section 19.2.4).

The BRST generator is $\Omega = \Omega^{\text{min}} + \Omega^{\text{non min}}$ where

$$\Omega^{\text{non min}} = \sum_{k=0}^{L} \sum_{s=k}^{L} b^k_{ka} \rho^k_{sa}.$$ (22)

The constant in the ghost number operator $\mathcal{G} = \mathcal{G}^{\text{min}} + \mathcal{G}^{\text{non min}}$ is adjusted so that $\mathcal{G}$ is antihermitian. With this choice, $\mathcal{G}$ is a sum over all conjugate pairs of terms of the form $i\eta\mathcal{P} + (-)^{\epsilon_\eta}(g/2)$, where $\epsilon_\eta$ is the parity of the conjugate pair $(\eta, \mathcal{P})$ in question.

We are now in the position to state the main results of our letter. There are:
Theorem 1. The states $|q, \text{ghosts}\rangle$ defined by

$$
|q, \text{ghosts}\rangle \equiv |q^i > \prod_{s=0}^{L_i} |\eta^{as} = 0 > \prod_{k=0}^{L_i} |\lambda^s_{a_s} = 0, \rho^s_{a_s} = 0 > \prod_{k=0}^{L_i} |b^s_{a_s} = 0, \bar{C}^s_{a_s} = 0 >,
$$

(23)

are (i) BRST invariant,

$$
\Omega |q, \text{ghosts}\rangle = 0,
$$

(24)

and (ii) of ghost number zero,

$$
\mathcal{G} |q, \text{ghosts}\rangle = 0.
$$

(25)

Theorem 2. The matrix elements $< q', \text{ghosts} | A | q, \text{ghosts} \rangle$ of the BRST invariant operator $A$ between the states of Theorem 1 are, up to inessential numerical factors that we shall not write, just equal the projected kernel (4) of the operator $A_0$, of which $A$ is the BRST invariant extension,

$$
< q', \text{ghosts} | A | q, \text{ghosts} \rangle \sim A_0^P(q', q).
$$

(26)

In particular, if one takes for $A$ the evolution operator $U(t' - t) = \exp -iH(t' - t)$, where $H$ is the BRST invariant extension of the Hamiltonian, one gets, using the fact that the projected kernel is annihilated by the constraints:

Theorem 3. The path integral

$$
U_0^P(q', t'; q, t) = \int \mathcal{D}q^i \mathcal{D}p_i \mathcal{D}\eta^{as} \mathcal{D}\lambda^s_{a_s} \mathcal{D}(\text{non minimal sector}) \exp i \int_t^{t'} dt (p\dot{q} + \mathcal{P}_{a_s} \eta^{as} + ....... - H),
$$

(27)

with boundary conditions

$$
q^i(t') = q^i, \quad q^i(t) = q^i,
$$

(28)

$$
\eta^{as}(t') = 0, \quad \eta^{as}(t) = 0,
$$

(29)

$$
\lambda^s_{a_s}(t') = 0, \quad \lambda^s_{a_s}(t) = 0 \quad k \text{ odd},
$$

(30)

$$
b^s_{a_s}(t') = 0, \quad b^s_{a_s}(t) = 0 \quad k \text{ even},
$$

(31)
\[ C^s_{k a_s}(t) = 0, \quad k \text{ even}, \]  \hspace{1cm} \tag{32}

\[ \rho^k_{s a_s}(t) = 0, \quad \rho^k_{s a_s}(t) = 0 \quad k \text{ odd}, \]  \hspace{1cm} \tag{33}

is a solution of the constraint equations.

In the irreducible case, this theorem reproduces the result of Ref.[3]. But Theorem 2 gives a more precise information about the nature of the path integral since it relates it explicitly to the projected kernel of the operator formalism.

**Proof of theorem 1:** The first assertion is obvious: the states (23) are manifestly annihilated by the minimal BRST charge as well as by each term of \( \Omega^{\text{non min}} \). To prove the second assertion, one check that the ghost number adds up to zero at each reducibility level. Consider the level \( s \), \( 0 \leq s \leq L \), i.e., all the variables parametrized by \( a_s \). Assume for definiteness that \( s \) is even so that the ghosts \( \eta^{a_s} \) are fermionic. [If \( s \) is odd, then all signs in the following discussion must be changed. This does not alter the conclusion that the total ghost number at level \( s \) adds up to zero.] With \( s \) even, \( b^k_{s a_s} \) and \( \lambda^k_{s a_s} \) (respectively \( \rho^k_{s a_s} \) and \( \bar{C}^k_{s a_s} \)) are bosonic (respectively fermionic) for \( k \) even, while they are fermionic (respectively bosonic) for \( k \) odd. We now count the ghost number: the state \( | \eta^{a_s} = 0 > \) of the minimal sector brings in \( (s + 1)m_s/2 \). The state \( | b^k_{s a_s} > = 0 > | \bar{C}^k_{s a_s} = 0 > (k \text{ even } \leq s) \) brings in \( [(s - k) + (k - s - 1)]m_s/2 = -m_s/2 \). The state \( | \lambda^k_{s a_s} > = 0 > | \rho^k_{s a_s} > = 0 > (k \text{ odd } \leq s) \) brings in \( [(s - k) + (k - s - 1)]m_s/2 = -m_s/2 \). Hence, the contribution of the non minimal sector, obtained by summing over \( k \) \( (0 \leq k \leq s) \), is just \( -(s + 1)m_s/2 \). This compensates the ghost number contribution from the minimal sector, yielding a ghost number equal to zero for the states \( |q, \text{ ghosts} > \).

**Proof of theorem 2:** The proof of theorem 2 proceeds in two steps. (i) first one verifies than the right and left hand sides of Eq.(26) transform in the same way under redefinitions of the constraints;

(ii) second, one verifies Eq.(26) in the representation where the constraints are abelian and the reducibility functions equal to 0 or 1.

The verification of (i) is direct and based on (11-14) above. To check (ii), one observes that for abelian constraints, “more” in eq.(8) turns out to be zero, and \( \Omega^{\text{min}} \) reduces to

\[ \Omega^{\text{min}} = \eta^{A_0} G_{A_0} + \sum_{s=0}^{L-1} \eta^{A_{s+1}} \mathcal{P}_{\alpha s}. \]  \hspace{1cm} \tag{34}

Here, (i) the \( G_{A_0} \) are a set of independent (abelian) constraints,

\[ G_{a_0} = (G_{A_0}, G_{\alpha_0} \equiv 0); \]  \hspace{1cm} \tag{35}

and (ii) the reducibility matrices are chosen to be

\[ Z_{A_{s+1}}^{\alpha s} = \sim \delta_{A_{s+1}}, \]  \hspace{1cm} \tag{36}
the other components being zero. We have split the indices \(a_s\) as

\[
a_s = (A_s, \alpha_s), \quad A_s = 1, \ldots, m_s - r_s, \quad \alpha_s = 1, \ldots, r_s.
\]

(37)

where \(r_s = \text{rank } Z^a_{a_s-1}\). There are no indices \(\alpha_L\) since the last \(Z^a_{a_L-1}\) are independent, and there are as many indices \(A_{s+1}\) as there are indices \(\alpha_s\).

Now, the form of the full BRST charge in the abelian representation is exactly the same as that of the BRST charge for an irreducible gauge system described by the following set of independent constraints:

(i) \(G_{A_0} = 0, \mathcal{P}_{\alpha_0} = 0, \ldots, \mathcal{P}_{\alpha_{L-1}} = 0\), with non minimal sector \([b^0_{0A_0}], [b^0_1A_1], \ldots, [b^0_{LAL}]\)

(we use the “block notation” \([b]\) of [4] to denote all the variables of a basic building block \((b, \rho, \lambda, \bar{C})\) of the non minimal sector);

(ii) \(b^{k-1}_{sA_s} = 0 (k \text{ odd})\) with non minimal sector \([b^{k-1}_{s-1\alpha_{s-1}}]\);

(iii) \(b^{k+1}_{s\alpha_{s+1}} = 0 (k \text{ odd}, k \leq L - 1)\) with non minimal sector \([b^{k+1}_{s+1A_{s+1}}]\).

The boundary conditions on the ghosts and non minimal variables are furthermore exactly the same as those that one would write down by considering the theory as an irreducible gauge theory with constraints (i)-(iii).

Therefore, if \(A_0(q,p)\) is an observable - which we may assume to commute strongly with the abelian constraints, \([A_0, G_{A_0}] = 0\) - , the analysis of the irreducible case applied to the system with constraints (i)-(iii) implies that

\[
< q', \text{ghosts} | A | q, \text{ghosts} > =
\]

\[
< q', \eta^{\alpha_s} = 0, \lambda^{s_{\alpha_s}} = 0 | A_0(q,p) \prod \delta(G_{A_0}) \delta(\mathcal{P}_{\alpha_i}) \delta(b^j_{s_{\alpha_s}}) | q, \eta^{\alpha_{\rho_{\lambda\alpha_{s}}}} = 0, \lambda^{s_{\alpha_s}} = 0 >
\]

(38)

up to inessential numerical factors (the index \(k\) of \(\lambda\) in (38) is odd) [4, chapter 16]. Since

\[
< u = 0 | \delta(p_u) | u = 0 >= (2\pi)^{-1} (u \text{ even}) \text{ or } 1 (u \text{ odd}),
\]

one finally gets

\[
< q', \text{ghosts} | A | q, \text{ghosts} > =< q' | A_0(q,p) \prod \delta(G_{A_0}) | q >
\]

(39)

up to unwritten irrelevant factors. The right hand side of Eq. (39) is the projected kernel for the case of \(m\) independent abelian constraints [4]. This achieves the proof of Theorem 2.

Once Theorem 2 is demonstrated, the proof of Theorem 3 is immediate : one simply takes for \(A\) the evolution operator and reexpresses \(< q', \text{ghosts} | U(t' - t) | q, \text{ghosts} >\) as a path integral.

We have thus established in this letter that the BRST path integral (27) for the evolution operator, with the boundary conditions (28)-(33), is a solution of the constraint equations. We have actually done more than merely verifying that (27) solves the constraints (zero is also a solution of the constraints). Namely, we have explicitly related the path integral to the projected kernel of the evolution operator. Thus, the path integral is non trivial and contains all the information about the evolution of the gauge system.
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FOOTNOTES

\#1 We stress that in the Dirac quantization method for constrained systems, all the constraints are imposed as conditions on the physical states. There exists an alternative approach, based in the Fock representation, in which only the holomorphic part of the constraints is enforced on the physical states. This approach is devoid of the scalar product difficulties and is implemented by means of different boundary conditions on the ghosts [4], chapter 13.

\#2 It is sometimes stated in the literature that the equations (3) are inconsistent when the constraints are reducible, and that the constraints should therefore not be imposed on the physical states. This is of course incorrect. In the absence of anomalies, the equations (3) are consistent both in the irreducible and the reducible cases. For example, the constraints for an abelian 2-form gauge field theory read $\pi^{ij,\gamma} \psi = 0$, i.e. $(\delta \psi / \delta A_{ij})_{,j} = 0$. These equations are not independent and consistent. Their general solution is $\psi = \psi[A^L]$, where $A^L$ is defined through $A_{ij} = \epsilon_{ijk} \partial^k A^L + \partial_i \theta_j - \partial_j \theta_i$ (in three spatial dimensions).
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