ON THE EQUIVALENCE BETWEEN
THE UNIFIED AND STANDARD VERSIONS
OF CONSTRAINT DYNAMICS

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ABSTRACT

The structure of physical operators and states of the unified constraint dynamics is studied. The genuine second–class constraints encoded are shown to be the superselection operators. The unified constrained dynamics is established to be physically–equivalent to the standard BFV–formalism with constraints split.

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1 Introduction

In previous papers [1-3] of the present authors a unified formalism has been suggested for
description of constrained dynamical systems without making use of explicit splitting the
constraints into the first– and second–class ones.

The generating equations, formulated in Ref.[1, 2], yield not only the constraint algebra
entirely, but also determine the fundamental phase variable commutators to satisfy the Jacobi
identity automatically. The latter circumstance is quite nontrivial because of the classical
limit existence requirement.

In recent paper [3] the generating equations have been strengthened in such a way that
their solution under the corresponding boundary conditions becomes determined uniquely up
to a natural canonical arbitrariness. Besides, the unified formulation of constrained dynamics
was given there by direct constructing the Unitarizing Hamiltonian without making use of
splitting the constraints into the first– and second–class ones.

The main purpose of the present paper is to show the formalism, developed in Ref.[1-3], to
be physically–equivalent to the standard BFV–formalism [4-7] with second–class constraints
presented explicitly [8]. In fact, we confine ourselves by the quasiclassical approximation,
having in mind the standard reasoning for extension towards the exact operator formulation.
Besides, we suppose the degenerate metric, defining the classical phase variable bracket, to
be of a constraint rank.

The following remark seems to be relevant here. According to the Dirac quantization
scheme, the second–class constraint operators are considered to equal to zero strongly. If
one constructs the corresponding BRST–BFV operator to be nilpotent, proceeding directly
from the second–class constraints as regarded to vanish strongly, then the nilpotent operator
constructed appears to have nontrivial unphysical cohomologies [9,10]. To eliminate these
cohomologies, it was suggested in paper [10] to introduce a tower of “ghosts for ghosts”.
However, this scheme appears, in fact, to imply the constraint splitting again.

In the scheme of Ref.[1-3], the Heisenberg equations of motion require for the encoded
second–class constraints to be conserved in time but not to vanish strongly. This crucial
circumstance allows one to eliminate the nonphysical cohomologies.

The present consideration is necessarily brief. All statements are formulated without
proofs. The detailed consideration will be published elsewhere.

2 Outline of the unified formalism

In this Section we recall in brief the form of the unified constraint algebra generating equa-
tions as well as the structure of the Unitarizing Hamiltonian.

Let \{\hat{\Gamma}^A\} be a set of the fundamental phase variable operators. For the sake of simplicity
we confine ourselves by the case in which all the operators $\hat{\Gamma}^A$ are supposed to be Bosons. Following the Refs.[1-3], we assign a classical ghost parameter $\Gamma^*_A$ to each initial operator $\hat{\Gamma}^A$. Being the operators $\hat{\Gamma}^A$ bosonic, all the ghost parameters $\Gamma^*_A$ are Fermions.

Let $\hat{\Theta}^\alpha(\hat{\Gamma})$ be a total set of irreducible constraint operators of the theory. We assign a canonical pair of ghost operators $(\hat{C}_\alpha, \hat{\bar{P}}^\alpha)$ to each constraint $\hat{\Theta}^\alpha(\hat{\Gamma})$.

So, we have the following list of values of the Grassmann parity ($\varepsilon$) and ghost number (gh) ascribed to the basic objects introduced:

$$\varepsilon(\hat{\Gamma}^A) = 0, \quad \varepsilon(\Gamma^*_A) = \varepsilon(\hat{C}_\alpha) = \varepsilon(\hat{\bar{P}}^\alpha) = 1,$$

$$\text{gh}(\hat{\Gamma}^A) = 0, \quad \text{gh}(\Gamma^*_A) = \text{gh}(\hat{C}_\alpha) = -\text{gh}(\hat{\bar{P}}^\alpha) = 1.$$  \hspace{1cm} (2.1)

Next, one introduces the generating operators $\hat{\Omega}, \hat{\Delta}, \hat{\Omega}_\alpha$,

$$\varepsilon(\hat{\Omega}) = \varepsilon(\hat{\Omega}_\alpha) = 1, \quad \varepsilon(\hat{\Delta}) = 0,$$

$$\text{gh}(\hat{\Omega}) = \text{gh}(\hat{\Omega}_\alpha) = 1, \quad \text{gh}(\hat{\Delta}) = 2,$$  \hspace{1cm} (2.2)

to satisfy the equations:

$$(\hbar)^{-1}[\hat{\Omega}, \hat{\Omega}] = \hat{\Delta}, \quad \hat{\Delta} |_{\Gamma^*_A=0} = 0,$$  \hspace{1cm} (2.5)

$$(\hbar)^{-1}[\hat{\Omega}, \hat{\Omega}_\alpha] = 0, \quad (\hbar)^{-1}[\hat{\Omega}_\alpha, \hat{\Omega}_\beta] = 0,$$  \hspace{1cm} (2.6)

together with their compatibility conditions:

$$(\hbar)^{-1}[\hat{\Omega}, \hat{\Delta}] = 0, \quad (\hbar)^{-1}[\hat{\Delta}, \hat{\Omega}_\alpha] = 0.$$  \hspace{1cm} (2.7)

A solution to the generating equations (2.5), (2.6) is sought in the form of a series expansion in powers of the ghost parameters $\Gamma^*_A$ and operators $(\hat{C}_\alpha, \hat{\bar{P}}^\alpha)$. The lowest orders of these expansions are:

$$\hat{\Omega} = \hat{C}_\alpha \hat{\Theta}^\alpha(\hat{\Gamma}) + \Gamma^*_A \hat{\Gamma}^A + \ldots,$$  \hspace{1cm} (2.8)

$$\hat{\Delta} = -2\hat{C}_\alpha \Gamma^*_A \hat{E}^{A\alpha}(\hat{\Gamma}) - \Gamma^*_B \Gamma^*_A \hat{D}^{AB}(\hat{\Gamma}) + \ldots,$$  \hspace{1cm} (2.9)

$$\hat{\Omega}_\alpha = \hat{C}_\beta \hat{\Lambda}^\beta_{\alpha}(\hat{\Gamma}) + \Gamma^*_A \hat{K}^A_{\alpha}(\hat{\Gamma}) + \ldots.$$  \hspace{1cm} (2.10)

The fundamental commutators:
\[ \hat{\Sigma}^{AB} \equiv (\hbar)^{-1} [\hat{\Gamma}^A, \hat{\Gamma}^B] \]  

are expressed in a certain way in terms of coefficient operators of the expansions (2.8), (2.9).

As it has been shown in Ref.[3], a canonical transformation there exists such that transformed operators \( \hat{\Omega}, \hat{\Delta}, \hat{\Omega}_\alpha \) are of the form:

\[ \hat{\Omega} = \hat{C}_\alpha \hat{\Theta}^\alpha(\hat{\Gamma}) + \Gamma^*_A \hat{\Gamma}^A(\hat{\Gamma}), \]

\[ \hat{\Delta} = -2\hat{C}_\alpha \Gamma^*_A (\hbar)^{-1} [\hat{\Gamma}^A, \hat{\Theta}^\alpha] - \Gamma^*_B \Gamma^*_A (\hbar)^{-1} [\hat{\Gamma}^A, \hat{\Gamma}^B], \]

\[ \hat{\Omega}_\alpha = \hat{C}_\alpha, \]

where \( \hat{\Theta}^\alpha \) are Abelian constraints:

\[ (\hbar)^{-1} [\hat{\Theta}^\alpha, \hat{\Theta}^\beta] = 0. \]

To describe the unified dynamics, one introduces the new operator \( \hat{\Phi} \) to satisfy the equations:

\[ (\hbar)^{-1} [\hat{\Phi}, \hat{\Omega}_\alpha] = \hat{\Omega}_\alpha \]

and boundary condition:

\[ \hat{\Phi}|_{r^* = 0} = \frac{1}{2}(\hat{C}_\alpha \hat{\mathcal{P}}^\alpha - \hat{\mathcal{P}}^\alpha \hat{C}_\alpha) \equiv \hat{G}_0. \]

Then, the truncated operator:

\[ \hat{\Omega}_T \equiv (\hbar)^{-1} [\hat{\Phi}, \hat{\Omega}], \quad \varepsilon(\hat{\Omega}_T) = 1, \quad \text{gh}(\hat{\Omega}_T) = 1, \]

possesses the nilpotency property:

\[ (\hbar)^{-1} [\hat{\Omega}_T, \hat{\Omega}_T] = 0. \]

In the tilded basis one has:

\[ \hat{\Omega}_T = \hat{C}_\alpha \hat{\Theta}^\alpha. \]

Having the operator \( \hat{\Omega}_T \) found, we define the Truncated Hamiltonian \( \hat{H}_T \),

\[ \varepsilon(\hat{H}_T) = \text{gh}(\hat{H}_T) = 0, \]

to satisfy the equations:
\[(i\hbar)^{-1} [\hat{H}_T, \hat{\Omega}_T] = 0, \quad (i\hbar)^{-1} [\hat{\Phi}, \hat{H}_T] = 0 \]  

(2.22)

and boundary condition:

\[ \hat{H}_T = \hat{H}_0 + O(\hat{\mathcal{P}}), \]  

(2.23)

where \( \hat{H}_0 \) is the initial Hamiltonian of the theory. In the tilded basis the operator \( \hat{H}_T \) does not depend on \( \Gamma^*_A \).

As a final step, we introduce canonical pairs of dynamically-active Lagrangian multipliers \((\hat{\lambda}_\alpha, \hat{\pi}^\alpha)\) and antighosts \((\hat{\mathcal{P}}_\alpha, \hat{\bar{C}}^\alpha)\),

\[ \varepsilon(\hat{\lambda}_\alpha) = \varepsilon(\hat{\pi}^\alpha) = 0, \quad \varepsilon(\hat{\mathcal{P}}_\alpha) = \varepsilon(\hat{\bar{C}}^\alpha) = 1, \]  

(2.24)

\[ \text{gh}(\hat{\lambda}_\alpha) = \text{gh}(\hat{\pi}^\alpha) = 0, \quad \text{gh}(\hat{\mathcal{P}}_\alpha) = -\text{gh}(\hat{\bar{C}}^\alpha) = 1, \]  

(2.25)

to define the total BRST-BFV-operator:

\[ \hat{Q} = \hat{\Omega}_T + \hat{\mathcal{P}}_\alpha \hat{\pi}^\alpha \]  

(2.26)

and Unitarizing Hamiltonian:

\[ \hat{H} = \hat{H}_T + (i\hbar)^{-1} [\hat{\Psi}, \hat{Q}], \]  

(2.27)

where \( \hat{\Psi} \) is a gauge Fermion operator,

\[ \varepsilon(\hat{\Psi}) = 1, \quad \text{gh}(\hat{\Psi}) = -1. \]  

(2.28)

Physical operators \( \hat{\mathcal{O}} \) and physical states \(|\text{Phys}\rangle\) are defined in the usual way:

\[ (i\hbar)^{-1} [\hat{\mathcal{O}}, \hat{Q}] = 0, \quad (i\hbar)^{-1} [\hat{\mathcal{O}}, \hat{G}] = 0, \]  

(2.29)

\[ \hat{Q}|\text{Phys}\rangle = 0, \quad \hat{G}|\text{Phys}\rangle = 0, \]  

(2.30)

where \( \hat{G} \) is the total ghost number operator:

\[ \hat{G} \equiv \frac{1}{2} (\hat{\mathcal{C}}_\alpha \hat{\mathcal{P}}^\alpha - \hat{\mathcal{P}}^\alpha \hat{\mathcal{C}}_\alpha) + \frac{1}{2} (\hat{\mathcal{P}}_\alpha \hat{\mathcal{C}}^\alpha - \hat{\mathcal{C}}^\alpha \hat{\mathcal{P}}_\alpha) - i\hbar \Gamma^*_A \frac{\partial}{\partial \Gamma^*_A} \]  

(2.31)

It should be emphasized once again that we have used no constraint splitting in constructing the Unitarizing Hamiltonian.
3 The unified formalism is equivalent to the standard one

In this Section we consider in details the structure of physical operators and states. Then we establish the standard and unified versions of constrained dynamics to be equivalent to each other.

First of all, being the fundamental bracket matrix $\Sigma_{AB} \equiv \{\Gamma^A, \Gamma^B\}$ of a constant rank, one can show that classical constraints are representable in the following split form:

$$\Theta^\alpha = M^\alpha_\beta \Theta^{\beta'} + M^\alpha_{\beta''} \Theta^{\beta''}$$

(3.1)

where $\beta = (\beta', \beta'')$, $\beta'' = 1, \ldots, \text{corank } |\Sigma_{AB}|$, the matrix $|M^\alpha_\beta| = |M^\alpha_{\beta'}, M^\alpha_{\beta''}|$ is invertible, and the relations hold:

$$\{\Gamma^A, \Theta^{\alpha''}\} = 0, \quad \{\Theta^{\alpha'}, \Theta^{\beta'}\} = U^{\alpha'}_{\gamma'}^{\beta'} \Theta^{\gamma'}.$$  

(3.2)

Thus $\Theta^{\alpha''}$ is nothing other but a set of genuine second–class constraints, while $\Theta^{\alpha'}$ are first–class constraints.

It is quite natural that the constraint operators $\hat{\Theta}^\alpha$ of the quantum theory admit an ordered counterpart of splitting (3.1), so that one has the corresponding operators $\hat{\Theta}^{\alpha'}$ and $\hat{\Theta}^{\alpha''}$ to be quantum first– and second–class constraints, respectively.

Now, let us turn to the tilded basis mentioned in the previous section. In this basis the theory is determined by the equations:

$$(\hat{i}\hbar)^{-1}[\hat{\Omega}_T, \hat{\Omega}_T] = 0, \quad (\hat{i}\hbar)^{-1}[\hat{H}_T, \hat{\Omega}_T] = 0,$$

(3.3)

where $\hat{\Omega}_T$ and $\hat{H}_T$ do not depend on $\Gamma^*_A$. The same as in the standard case, one can confirm that the solution for $\hat{\Omega}_T$ is determined uniquely, up to a canonical transformation, by the boundary value, while a change of the constraint basis is also induced by a canonical transformation of $\hat{\Omega}_T$. The solution for $\hat{H}_T$ is determined, up to a contribution of the form $(\hat{i}\hbar)^{-1}[\hat{\Psi}, \hat{\Omega}]$, by the boundary value, while a change $\hat{H}_0 \to \hat{H}_0 + O(\hat{\Theta})$ is induced by a transformation of the form $\hat{H}_T \to \hat{H}_T + (\hat{i}\hbar)^{-1}[\hat{\Psi}, \hat{\Omega}_T]$.

Let us suppose that all the required canonical transformations are performed. Besides, let us turn to the new phase variables, $\hat{\Gamma}^A \to \hat{\Gamma}'^A = (\hat{\gamma}'^i, \hat{\Theta}^{\alpha''})$, where the operators $\hat{\gamma}'^i$, in principle, can be considered to form a set of canonical pairs. Then in an appropriate basis the operators $\hat{Q}$ and $\hat{H}$ possess the structure:

$$\hat{Q} = \hat{Q}' + \hat{Q}'', \quad \hat{H} = \hat{H}' + (\hat{i}\hbar)^{-1}[\hat{\Psi}, \hat{Q}].$$

(3.4)

where the nilpotent operator $\hat{Q}'$ is constructed in the usual way, proceeding form the first–class constraints $\hat{\Theta}^{\alpha'}$, and depends on the variables $\hat{X}' \equiv (\hat{\gamma}'^i, \hat{C}_{\alpha'}, \hat{P}_{\alpha'}, \hat{P}_{\alpha'}, \hat{C}_{\alpha'}, \hat{\lambda}_{\alpha'}, \hat{\pi}_{\alpha'})$, while the operator $\hat{Q}''$ is of the form:
\[ \hat{Q}'' = \hat{C}_\alpha'' \hat{\Theta}'' + \hat{P}_\alpha'' \hat{\pi}''; \]  
(3.5)

\( \hat{H}' \) is the standard Unitarizing Hamiltonian of the BFV–formalism for the variables \( \hat{X}' \), first–class constraints \( \hat{\Theta}' \), nilpotent operator \( \hat{Q}' \) and boundary value

\[ H_0(\gamma) \equiv H_0|_{\Theta''=0}. \]  
(3.6)

Besides, the definition of the variables \( \hat{X}' \) implies natural splitting of the form:

\[ \hat{C}_\alpha = (\hat{C}_\alpha', \hat{C}_\alpha''), \ldots \]  
(3.7)

for the ghosts \( \hat{C}_\alpha \), \( \hat{P}_\alpha \) antighosts \( \hat{P}_\alpha \), \( \hat{C}_\alpha \) and Lagrangian multipliers \( \hat{\lambda}_\alpha \), \( \hat{\pi}_\alpha \).

Further, as the second–class constraints \( \hat{\Theta}'' \) commute with all the variables, the Heisenberg equations of motion, as applied to the \( \hat{\Theta}'' \), yield:

\[ d_t \hat{\Theta}'' = (i\hbar)^{-1}[\hat{\Theta}''', \hat{H}] = 0, \]  
(3.8)

and hence the conservation law holds:

\[ \hat{\Theta}'' = \text{const.} \]  
(3.9)

However, this constant is not to vanish certainly, but rather one should considered \( \hat{\Theta}'' \) to be the superselection operators [11]. The mentioned circumstance appears to be of crucial importance for establishing the sructure of physical space.

Let us consider the structure of physical operators. We suppose the operators \( \hat{O} \) to be formal polynomials in all the variables. One should consider all the operator equations, (2.29) for instance, to be operator equalities valid for arbitrary values of \( \hat{\Theta}'' \). Then one can establish the physical operators \( \hat{O} \) to have the structure:

\[ \hat{O} = \hat{O}' + (i\hbar)^{-1}[\hat{\Theta}'', \hat{Q}], \]  
(3.10)

where \( \hat{A} \) is an arbitrary operator, while the operator \( \hat{O}' \) depends on the variables \( \hat{X}' \) only, and satisfies the equation

\[ (i\hbar)^{-1}[\hat{O}', \hat{Q}'] = 0. \]  
(3.11)

together with the condition

\[ \text{gh}(\hat{O}') = 0. \]  
(3.12)

Thus we conclude that the class of physical operators \( \hat{O} \) is equivalent to the class of physical operators \( \hat{O}' \) of the BFV–formalism for the variables \( \hat{X}' \).
Let us consider the structure of physical subspace. We suppose the total space of states to be generated by applying the Heisenberg–field formal polynomials to the cyclic vector (vacuum), the same as it holds in theories without the superselection operators. This space is invariant under the action of physical operators, the Hamiltonian or $S$-matrix for instance. Besides, we suppose the vacuum $|0\rangle$ to be a physical state.

Let:

$$|\text{Phys}\rangle = \hat{P}(\hat{X}, \Gamma^*)|0\rangle,$$

$$\hat{Q}|\text{Phys}\rangle = 0, \quad \hat{G}|\text{Phys}\rangle = 0,$$

where $\hat{X}$ denotes the total set of variables of the extended phase space operators, and $\hat{P}$ is a formal polynomial.

It follows from (3.13), (3.14) that:

$$(i\hbar)^{-1}[\hat{P}, \hat{Q}] = 0, \quad (i\hbar)^{-1}[\hat{P}, \hat{G}] = 0.$$

i.e. $\hat{P}$ is a physical operator.

Further, one can show that if the equations

$$\hat{Q}|\text{Phys}\rangle = 0, \quad \hat{\Theta}^{\alpha''}|\text{Phys}\rangle = \xi^{\alpha''}|\text{Phys}\rangle \neq 0$$

hold, where $\xi^{\alpha''}$ are some numbers, then

$$|\text{Phys}\rangle = \hat{Q}|\xi\rangle$$

for some vector $|\xi\rangle$.

Thus the second–class constraints vanish effectively in the physical subspace whose structure and dynamics coincide exactly with the ones of the BFV–formalism for the basic operators $\hat{\gamma}^i$, first–class constraints $\hat{\Theta}^{\alpha^i}$ and Hamiltonian $\hat{H}_0(\hat{\gamma})$.

## 4 Equivalence between the formalisms from the functional–integral viewpoint

In this section we elucidate the equivalence between the unified and standard versions of constrained dynamics from the functional–integral viewpoint.

Let us consider the generating functional of quantum Green’s functions:

$$Z(J) = \langle 0|T\exp\{\frac{i}{\hbar} \int J_a \hat{X}^a dt\}|0\rangle,$$

where $J_a(t)$ is the external source, and

$$\hat{X}^a(t) \equiv (\hat{\Gamma}^A, \hat{C}_\alpha, \hat{P}^\alpha; \hat{P}_\alpha, \hat{\bar{C}}^\alpha; \hat{\lambda}_\alpha, \hat{\pi}^\alpha)$$
are the Heisenberg operators,

\[ d_t \hat{X}^a = (\hbar)^{-1}[\hat{X}^a, \hat{H}] \]  

As is usual, the following equations hold for the generating functional (4.1):

\[ \left( \dot{X}^a - \{X^a, H\} + J_b \{X^a, X^b\} \right) |_{X=-\hbar \delta/\delta J} Z = 0. \]  

Let us seek for a solution to these equations by making use of the functional Fourier transform:

\[ Z(J) = \int DX \exp \left\{ \frac{i}{\hbar} \int J_a X^a dt \right\} \tilde{Z}(X). \]  

It follows from (4.4) that:

\[ \tilde{Z}(X) = [\prod_t \prod_{\alpha''} \delta(\dot{\Theta}^{\alpha''}(t))] \tilde{Z}_1(X), \]  

where \( \Theta^{\alpha''} \) are the second–class constraints, commuting strongly with all the phase variables.

Thus the functional integral (4.5) appears to be concentrated on the time–independent values

\[ \Theta^{\alpha''} = \text{const}. \]  

At the present stage we have to formulate an important assumption. Namely, we suppose the representation (4.5) to be valid for the concrete superselection sector. As for the functional integral itself, the above assumption means that the integration is concentrated on such trajectories that the functions \( \Theta^{\alpha''}(X(t)) \) take asymptotically just the values corresponding to the superselection sector chosen.

It has been established in the previous Section that the superselection sectors with nonzero values of \( \Theta^{\alpha''} \) are physically–trivial. The only nontrivial sector is the one with \( \Theta^{\alpha''} = 0 \). In this sector one has:

\[ \prod_t \delta(\dot{\Theta}^{\alpha''}) = (\text{Det} d/dt)^{-1} \prod_t \delta(\Theta^{\alpha''}). \]  

Performing a transformation to the representation, used in the previous Section, and choosing in (3.4) a gauge Fermion to be of the form

\[ \Psi = \mathcal{P}^{\alpha''} \lambda_{\alpha''} \]  

one obtains:

\[ ^1 \text{The same as in the standard case, one can confirm that there exists a change of integration variables that transforms the integrals with different } \Psi \text{ into each other.} \]
\[ Z(J) = \int DX' \exp \left\{ \frac{i}{\hbar} \int \frac{1}{2} X'(\{X', X'\}')^{-1} \dot{X}' - H' + J'X' \right\} dt \], \quad (4.10) 

where external sources are introduced directly to the variables \( X' \), the variables \( \gamma' \) are canonical, and \( \{ , \}' \) means the canonical brackets for all the variables \( X' \)

The expressions (4.10) coincides with the standard BFV–formalism prescription.

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