On the complex structure in the Gupta–Bleuler quantization method

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Abstract

We examine the general conditions for the existence of the complex structure intrinsic in the Gupta-Bleuler quantization method for the specific case of mixed first and second class fermionic constraints in an arbitrary space-time dimension. The cases $d = 3$ and $10$ are shown to be of prime importance. The explicit solution for $d = 10$ is presented.

1 Introduction

Since its invention in quantum electrodynamics [1], the Gupta-Bleuler method has become a conventional tool when quantizing theories with anomalies and/or second class constraints. In the latter case, it requires the construction of a specific complex structure $J$ on a phase space of a model which allows one to split the original second class constraints into (complex conjugate) holomorphic and antiholomorphic sets [2,3]. The existence of such a $J$ in a neighborhood of a (second class) constraint surface has been proven in Ref. [4]. It was stressed in [4], however, that, generally, this may break manifest covariance in a problem. If the second class constraints are a-priori in the holomorphic representation, the Gupta-Bleuler method was shown to admit an elegant BRST formulation [2,4–6], which

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involves a pair of (Hermitean conjugate) BRST charges (see also related works [7,8]).

As it has already been discussed in Ref. [2], the approach applies also to the specific case of mixed first and second class fermionic constraints

\[
\{ L_\alpha, L_\beta \} = 2i(\Gamma^n p_n)_{\alpha\beta} \equiv \Delta_{\alpha\beta}, \quad \alpha = 1, \ldots, n, \\
\{ L_\alpha, p^2 \} = 0, \\
\text{rank} (\Gamma^n p_n)_{\alpha\beta} = n/2, \\
\tag{1}
\]

with \( p^n \) a light-like vector \( p^2 = 0 \), which is just the case when studying the superparticle [9] and superstring [10] theories in flat superspace \( \mathbb{R}^{1,3} \). Following the procedure, one has to extract first class constraints from the original mixed system \( L_\alpha \approx 0 \) in the covariant (and reducible) way

\[
(\Gamma^n p_n L)_\alpha = 0. \\
\tag{2}
\]

and then solve the equations

\[
J^\alpha_\beta J^\beta_\gamma = -\delta^\alpha_\gamma, \\
J^\alpha_\beta \Delta^\beta_\gamma + \Delta^\alpha_\beta J^\beta_\gamma = 0, \\
\{ L_\alpha, J^\beta_\gamma \} = 0, \quad \{ J^\alpha_\beta, J^\gamma_\delta \} = 0 \\
\{ J^\alpha_\beta, p^2 \} = 0, \\
\tag{3a-3d}
\]

for the complex structure. With \( J \) at hand, the mixed constraints can be split into holomorphic \( L^-_\alpha \equiv p^-_\alpha \beta L_\beta \approx 0 \) and antiholomorphic \( L^-_\alpha \equiv p^-_\alpha \beta L_\beta \approx 0 \) sets

\[
\{ L^-_\alpha, L^-_\beta \} \approx 0, \quad \{ L^+_\alpha, L^+_\beta \} \approx 0, \\
\{ L^+_\alpha, L^-_\beta \} \approx p^+_\alpha \gamma \Delta^\gamma_\beta, \\
(\bar{L}^+_\alpha)^* = L^-_\alpha, \\
\tag{4}
\]

where \( p^\pm_\alpha \beta \equiv \frac{1}{2}(\delta^\alpha_\beta \pm iJ^\alpha_\beta) \). Half of these can further be used to define physical states in a complete Hilbert space

\[
\hat{L}^-_\alpha |\text{phys}\rangle = 0, \quad (\Gamma^n \hat{p}_n \hat{L})_\alpha |\text{phys}\rangle = 0, \quad \hat{p}^2 |\text{phys}\rangle = 0. \\
\tag{5}
\]

Note that only half of the original second class constraints were effectively used in Eq. (5).

\footnote{For simplicity, in what follows we shall discuss the superparticle case only.}
The scheme outlined was shown to admit a remarkably simple solution for the 4d superparticle [5]

\[ J \sim \gamma^5. \]  

(6)

In that case, Eq. (5) was proven to produce a massless irreducible representation of the super Poincaré group of superhelicity 0 (on-shell massless chiral scalar superfield).

However, the general conditions for the existence of \( J \) in an arbitrary space-time dimension are unknown yet. It is this problem which we address in the present work. As shown below, the cases \( d = 3 \) and 10 are of prime interest. The explicit solution for 10d is the main result of this letter.

In the next section we briefly outline our strategy and specify the dependence of the formalism on the dimension of space-time. Section 3 contains a solution for \( d = 10 \) in the Hamiltonian formalism. As shown below, the possibility to construct a covariant \( J \) satisfying Eq. (3) requires an extension of the original phase space by additional (unphysical) variables. In Sec. 4 we construct the corresponding 10d Lagrangian formulation. Some comments on the problem in 3d superspace are presented in Sec. 5. In particular, we show that the Gupta-Bleuler approach fails in that case. We end the paper with the discussion of open problems in Sec. 6. Some technical details are gathered in Appendix.

2 How to attack the problem?

In short, our proposal is to decompose the tensor \( J_\alpha^\beta \) into irreducible representations of the Lorentz group and then reduce Eq. (3) to equations for those irreps. It should be stressed, that the structure of the equations will essentially depend on the dimension of space-time and the type of spinors existing in a given dimension. If a spinor entering into the superparticle action is of Dirac type, the corresponding constraints will automatically belong to the holomorphic representation. In \( d = 2, 3, 4 \) (mod 8) a spinor can be chosen to be Majorana (and Majorana-Weyl in \( d = 2 \) (mod 8)). For even dimensions it can be decomposed into the direct sum of two (complex conjugate) Weyl spinors, which, again, brings the constraints to the holomorphic representation. Thus, if by analogy with the superstring theory one restricts ourselves with the case \( d < 11 \), there remain only two physically nontrivial possibilities \( d = 3 \) and 10 which we examine below. Note also, that, generally, the possibility to realize a covariant \( J \) satisfying
Eq. (3) will require an extension of the original phase space by additional (unphysical) variables.

3 A solution for $d = 10$

A minimal spinor representation of the Lorentz group in $d = 10$ is realized on Majorana–Weyl spinors (16 real components). Chiral spinors are distinguished by the position of their indices $\psi^\alpha$, $\varphi^\alpha$, $\alpha = 1, \ldots, 16$ (for the $10d$ spinor notation see Ref. [11]). By making use of the Fierz identity\footnote{In what follows, $\Gamma^{ab}, \Gamma^{abcd}, \ldots$ denote the totally antisymmetrized product of the 10d $\Gamma$-matrices (see also Appendix).}

\begin{equation}
32\delta^\alpha_\delta \delta^\beta_\gamma = 2\delta^\alpha_\beta \delta^\delta_\gamma - \Gamma^{ab}_\alpha \beta \Gamma_{ab\gamma} \delta + \frac{2}{4!} \Gamma^{abcd}_\alpha \beta \Gamma_{abcd\gamma} \delta \tag{7}
\end{equation}

one can decompose $J^\alpha_\beta$ into irreducible pieces

\begin{equation}
J^\alpha_\beta = \delta^\alpha_\beta J + \Gamma^{ab}_\alpha \beta J_{ab} + \Gamma^{abcd}_\alpha \beta J_{abcd}. \tag{8}
\end{equation}

A substitution of this into Eq. (3b) yields (see also Appendix)

\begin{equation}
2\Gamma^n (J_{pn} + 2J_{nm}p^m) + 2\Gamma^{abcdn} J_{[abcdp]n} = 0, \tag{9}
\end{equation}

or

\begin{equation}
J_{pn} + 2J_{nm}p^m = 0, \quad J_{[abcdp]n} = 0. \tag{10}
\end{equation}

Analogously, Eq. (3a) reads

\begin{align}
1 + J^2 - 2J_{ab}J^{ab} + b_3J_{abcd}J^{abcd} & = 0, \\
2J_{ab} + a_2J_{abcd}J^{cd} + \frac{1}{2!} \epsilon_{abcdefghij} J^{[cdef} J^{ghij]} & = 0, \tag{11}
\end{align}

\begin{align}
2J_{abcd} + J_{[abcd]}J_{[cd]} + a_1J_{[abcn]}J_{mn}^m + b_2J_{[abmn]}J_{mn}^{cd} & - \\
- \frac{1}{4!} \epsilon_{abcdefghij} (J^{[cdef} J^{ghij]} + b_1J^{[cfgm} J_{m}^{hij]} & = 0,
\end{align}

with $a_1, a_2, b_1, b_2, b_3$ denoting some constants (in what follows we will not need their explicit form) and $\epsilon_{abcdefghij}$ the 10d Levi-Civita tensor.

In obtaining Eq. (11) the identities [12]

\begin{align}
\Gamma^{(n)a_1 \ldots a_n}_\alpha^\beta = (-1)^{n/2} & \frac{\epsilon^{a_1 \ldots a_n a_{n+1} \ldots a_{10}}_{a_{n+1} \ldots a_{10} \alpha}}{(10-n)!} \Gamma^{(10-n)}_{a_{n+1} \ldots a_{10} \alpha} \quad & n \text{ even} \\
\Gamma^{(n)a_1 \ldots a_n}_\alpha^\beta = (-1)^{(n-1)/2} & \frac{\epsilon^{a_1 \ldots a_n a_{n+1} \ldots a_{10}}_{a_{n+1} \ldots a_{10} \alpha}}{(10-n)!} \Gamma^{(10-n)}_{a_{n+1} \ldots a_{10} \alpha} \quad & n \text{ odd} \tag{12}
\end{align}
have been used.

In the presence of the tensor $J_{abcd}$ Eq. (11) looks rather complicated. It is instructive then to try setting

$$J_{abcd} = 0,$$

which considerably simplifies Eqs. (10) and (11)

$$Jp_n + 2J_{nm}p^m = 0,$$

$$1 + J^2 - 2J_{ab}J^{ab} = 0,$$

$$JJ_{ab} = 0,$$

$$J_{ab}J_{cd} + J_{ac}J_{db} + J_{ad}J_{bc} = 0.$$

It turns out that this system does admit a solution. Actually, from Eq. (14c) it follows that either $J = 0$ or $J_{ab} = 0$ (the choice when they are both equal to zero is in a contradiction with Eq. (14b)). It is straightforward to check that the latter case leads to a contradiction between Eq. (14a) and (14b) (it is enough to consider those equations in the rest frame $p_n = (E, 0, \ldots, 0, E)$). Thus, one has to put

$$J = 0,$$

which brings Eq. (14) to the form

$$J_{nm}p^m = 0,$$

$$1 - 2J_{ab}J^{ab} = 0,$$

$$J_{ab}J_{cd} + J_{ac}J_{db} + J_{ad}J_{bc} = 0.$$

The simplest solution to Eq. (16c) reads

$$J_{ab} = \frac{1}{\alpha}(A_aB_b - A_bB_a),$$

with $A, B$ denoting some vectors and $\alpha$ a scalar. The substitution of this into Eq. (16b) determines $\alpha$:

$$\alpha = \pm 2\sqrt{A^2B^2 - (AB)^2}.$$

In addition, Eq. (16a) requires

$$(Ap) = 0, \quad (Bp) = 0,$$
because, otherwise, it would mean that \( A \) and \( B \) are linearly dependent and, hence, \( J_{ab} = 0 \). One can check, further, that it is impossible to construct two vectors \( A, B \) (satisfying Eq. (19) on the constraint surface) from the phase space variables of the original superparticle model. This suggests an extension of the space by two new variables \( A^n, B^n \). In order for them to be nondynamical, they should be subject to the (first class) constraints

\[
p_{An} = 0, \quad p_{Bn} = 0, \tag{20}
\]

In Eq. (20) \( p_A, p_B \) denote momenta canonically conjugated to \( A, B \) respectively. Following this course, Eq. (19) can further be incorporated into the scheme as gauge fixing conditions for some of the constraints (20). Actually, let us extend the original phase space by one more canonical pair \( (\Lambda^n, P_{\Lambda n}) \). In order to suppress the dynamics in this sector, we impose the constraints [11]

\[
\Lambda^2 = 0, \quad (\Lambda p) = -1, \quad P_{\Lambda n} = 0, \tag{21}
\]

or, equivalently,

\[
\Lambda^2 = 0, \quad (P_{\Lambda p}) = 0, \tag{22a}
\]

\[
(\Lambda p) = -1, \quad (P_{\Lambda \Lambda}) = 0, \tag{22b}
\]

\[
P_{\Lambda n} + \Lambda_n (P_{\Lambda p}) + p_n (P_{\Lambda \Lambda}) = 0. \tag{22c}
\]

The constraints (22a), (22b) are second class, while there are only eight linearly independent first class ones in Eq. (22c) (one can find two identities for the constraint set (22)). Note that the total number of constraints (22) is sufficient to suppress just one canonical pair of variables.

Let us now impose two gauge fixing conditions to the first class constraints (20)

\[
p_{An} = 0, \quad (Ap) = 0, \tag{23a}
\]

\[
p_{Bn} = 0, \quad (Bp) = 0, \tag{23b}
\]

and make use of \( \Lambda \) to split these into first and second class

\[
(Ap) = 0, \quad (P_{\Lambda p}) = 0, \tag{24a}
\]

\[
p_{An} + p_n (P_{\Lambda p}) = 0, \tag{24b}
\]

\[
(Bp) = 0, \quad (P_{\Lambda \Lambda}) = 0, \tag{25a}
\]

\[
p_{Bn} + p_n (P_{\Lambda \Lambda}) = 0. \tag{25b}
\]
The constraints (24a) ((25a)) are second class, whereas the first class ones (24b) ((25b)) contain only nine linearly independent components. The total number of constraints is sufficient to suppress the dynamics in the sector \((A, p_A), (B, p_B)\). In other words, the variables \((A, p_A), (B, p_B)\) are unphysical. Note also, that Eqs. (3c) and (3d) automatically holds when extending the space in this way.\footnote{In passing a to quantum description one has to fix completely the gauge freedom in the sector \((A, p_A), (B, p_B)\) because, otherwise, the vanishing of the first class constraints (24b) and (25b) on physical states would be incompatible with the prescription (5). In order to maintain the manifest covariance in the problem, it is sufficient to introduce eight sectors like Eq. (21), when imposing a gauge choice. The corresponding Lagrangian formulation will be presented elsewhere.}

Thus, in the enlarged phase space Eqs. (8),(13),(15),(17) and (18) realize the needed complex structure.

### 4 A 10d Lagrangian formulation

A Lagrangian formulation which reproduces Eqs. (21) and (23) when passing to the Hamiltonian formalism reads

\[
S = \int d\tau \frac{1}{2e} \left( \dot{x}^n - i\theta \Gamma^n \dot{\theta} - \omega \Lambda^n - \omega_1 A^n - \omega_2 B^n \right)^2 - \omega - \Phi \Lambda^2. \tag{26}
\]

As compared to the Casalbuoni-Brink–Schwarz model, this Lagrangian involves a set of auxiliary variables \((\omega, \omega_1, \omega_2, \Phi, \Lambda^n, A^n, B^n)\).

Moving to the Hamiltonian formalism one finds the primary constraints (we denote as \((p_e, p, p_\theta, p_\omega, p_{\omega_1}, p_{\omega_2}, p_\Phi, p_\Lambda, p_A, p_B)\) the momenta canonically conjugated to the variables \((e, x, \theta, \omega, \omega_1, \omega_2, \Phi, \Lambda, A, B)\), respectively)

\[
\begin{align*}
    p_e &= 0, & p_\theta + i\theta \Gamma^n p_n &= 0, & p_\Phi &= 0, \\
    p_\omega &= 0, & p_{\omega_1} &= 0, & p_{\omega_2} &= 0, \\
    p_\Lambda^n &= 0, & p_A^n &= 0, & p_B^n &= 0,
\end{align*}
\tag{27}
\]

and the relation determining \(\dot{x}^n\) as a function of some of the remaining variables

\[
\dot{x}_n = e p_n + i\theta \Gamma_n \dot{\theta} + \omega \Lambda_n + \omega_1 A_n + \omega_2 B_n. \tag{28}
\]

The canonical Hamiltonian looks like

\[
H = (p_\theta + i\theta \Gamma^n p_n) \lambda_\theta + p_e \lambda_e + p_\omega \lambda_\omega + p_{\omega_1} \lambda_{\omega_1} + p_{\omega_2} \lambda_{\omega_2} + p_\Phi \lambda_\Phi + p_\Lambda \lambda_\Lambda + p_A \lambda_A + p_B \lambda_B + e \frac{p^2}{2} + \omega (1 + \Lambda p) + \omega_1 (pA) + \omega_2 (pB) + \Phi \Lambda^2, \tag{29}
\]

where the \(\lambda\)'s denote Lagrange multipliers enforcing the primary constraints.
The preservation in time of the primary constraints implies the secondary ones
\[
\Lambda^2 = 0, \quad \Lambda p + 1 = 0, \quad p^2 = 0, \quad (30a)
\]
\[
(pA) = 0, \quad (pB) = 0, \quad (30b)
\]
\[
\omega_1 p^n = 0, \quad \omega_2 p^n = 0, \quad (30c)
\]
\[
\omega p^n + 2\Phi \Lambda^n = 0, \quad (30d)
\]
and determines half of the \(\lambda_\theta\)
\[
\Gamma^n p_n \lambda_\theta = 0. \quad (31)
\]
Consider now Eq. (30d). Multiplying it by \(\Lambda^n, p^n\) and taking into account Eq. (30a) one gets
\[
\omega = 0, \quad \Phi = 0. \quad (32)
\]
Analogously, Eq. (30c) provides us with
\[
\omega_1 = 0, \quad \omega_2 = 0. \quad (33)
\]
The consistency check for the secondary constraints yields
\[
p\lambda_A = 0, \quad \Lambda \lambda_A = 0,
p\lambda_A = 0, \quad p\lambda_B = 0,
\lambda_\omega = 0, \quad \lambda_\Phi = 0,
\lambda_\omega_1 = 0, \quad \lambda_\omega_2 = 0, \quad (34)
\]
and no tertiary constraints appear.

Thus, the complete constraint system can be written in the form
\[
p_e = 0, \quad p_\theta + i\theta \Gamma^n p_n = 0, \quad p^2 = 0, \quad (35a)
\]
\[
\omega = 0, \quad p_\omega = 0,
\omega_1 = 0, \quad p_\omega_1 = 0,
\omega_2 = 0, \quad p_\omega_2 = 0,
\Phi = 0, \quad p_\Phi = 0, \quad (35b)
\]
\[
\Lambda^2 = 0, \quad (\Lambda p) + 1 = 0, \quad p_\Lambda^n = 0, \quad (35c)
\]
\[
(pA) = 0, \quad p_A^n = 0, \quad (35d)
\]
\[
(pB) = 0, \quad p_B^n = 0. \quad (35e)
\]
The constraints (35a) are just those of the Casalbuoni-Brink–Schwarz superparticle. The constraints (35b) are second class and can be omitted.
after introducing the associated Dirac bracket, which leaves us with the needed constraints (21), (23) in Eqs. (35c)-(35e).

It is straightforward to check, further, that in the light-cone gauge

\[ \Gamma^+ \theta = 0, \quad x^+ = \tau p^+, \]
\[ e = 1, \quad \Lambda^i = 0, \]
\[ A^2 = 1, \quad A^i = 0, \]
\[ B^2 = 1, \quad B^i = 0, \]

the physical sector of the theory described by the action (26) coincides with the one of the Casalbuoni-Brink–Schwarz superparticle [9]. This proves the physical equivalence of the models.

5 A comment on the problem in 3d superspace

The minimal spinor representation of the Lorentz group in \( R^{3/2} \) superspace is realized on Majorana spinors \( \theta^\alpha, \alpha = 1, 2 \), \( (\theta^\alpha)^* = \theta^\alpha \). In that case, the fermionic constraints \( L_\alpha \approx 0, \alpha = 1, 2 \) involve only one second class constraint which, evidently, can not be separated into holomorphic and antiholomorphic sets. It is straightforward to check that Eqs. (3), being reduced to the 3d superspace (with \( \Delta = 2iC\Gamma^n p_n \) and \( C \) the charge conjugation matrix), imply

\[ p_n = 0, \]

and, hence, possesses no physically sensible solution.

Thus, the 3d superparticle yields an example when the Gupta-Bleuler approach fails.

6 Discussion

We conclude with some remarks and open problems.

a) When analyzing the problem in 10d superspace the condition \( J_{abcd} = 0 \) has been chosen by hands. It still remains to be understood, whether there are other solutions to Eq. (3) and, if it is the case, how they are related to each other.

b) (A comment on quantization) The complex structure constructed involves \( \frac{1}{\sqrt{A^2 B^2 - (AB)^2}} \) which, generally, leads to a nonlocal operator at the

\[ \text{Here we assume that } p^i \neq 0. \]

\[ \text{To see this, it is sufficient to decompose the } J_{\alpha \beta} \text{ with respect to the complete basis } \{1, (\Gamma^n)^T\} \text{ in the space of } 2 \times 2 \text{ complex matrices } J_{\alpha \beta} = J_\delta \delta_{\alpha \beta} + J_\gamma \Gamma^n \gamma_{\alpha \beta} \text{ and plug this into Eqs. (3a), (3b).} \]
quantum level. Note, however, that this can be avoided in the covariant gauge $A^2 = 1, B^2 = 1, (AB) = 0$. Since the commutation relations in the sector $(x, p), (\theta, p_\theta)$ are canonical, it suffices to realize the quantum brackets in the sector $(A, p_A), (B, p_B), (\Lambda, p_\Lambda)$. This work is currently in progress.

c) There are two comments on a possible generalization to superstrings. First of all, since the super Virasoro constraints weakly commute with the fermionic ones, Eqs. (1)-(3) will take a more complicated form. Secondly, in a recent work [13] Berkovits proved that the generalization of Eq. (21) to the superstring case, which was previously used in [11,14,15], is not harmless. There remains one physical zero mode in the sector of the additional variables. Because of this reason, it is not obvious to us how to construct the corresponding Lagrangian formulation, when generalizing the above construction to superstrings.

Appendix

Our convention for antisymmetrization of indices are as follows

\[ A_{[ab]} = \frac{1}{2} (A_a B_b - A_b B_a), \]
\[ A_{[abc]} = \frac{1}{3} (A_{[bc]} + A_{b[c}] + A_{c[ab]}), \]
\[ A_{[abcd]} = \frac{1}{4} (A_{[bcde]} + A_{b[cde]} + A_{c[deab]} + A_{d[eab]} + A_{d[eabc]} + A_{e[abcd]}), \]
\[ A_{[abcde]} = \frac{1}{5} (A_{[bcdef]} + A_{b[cdef]} + A_{c[deab]} + A_{d[eabc]} + A_{e[abcd]}), \]
\[ A_{[abcdef]} = \frac{1}{6} (A_{[bcdef]} - A_{b[cdefa]} + A_{c[deab]} - A_{d[eabc]} + A_{e[abedf]} - A_{f[abcde]}), \]

and so on.

In particular, given two totally antisymmetric tensors $J_{ab}$ and $J_{abcd}$ one has

\[ J_{[ab]} J_{cd} = \frac{1}{3} (J_{ab} J_{cd} + J_{ac} J_{db} + J_{ad} J_{bc}); \]
\[ J_{[ab]} J_{cdem} = \frac{1}{15} (J_{ab} J_{cdem} + J_{ac} J_{demb} + J_{ad} J_{embc} + J_{ac} J_{mbcd} + J_{am} J_{bcde} - \]

\footnote{For issues on the covariant quantization of the Green-Schwarz strings see e.g. [16,17].}
\[-J_{de}J_{mabc} - J_{dm}J_{abce} - J_{be}J_{dema} - J_{bd}J_{emac} - J_{be}J_{macd} - J_{bm}J_{acde} + \\
+ J_{em}J_{abcd} + J_{cd}J_{emab} + J_{ce}J_{mabd} + J_{cm}J_{abde};
\]

\[J_{[abc}m_J^{m]d} = \frac{1}{4}(J_{abcm}J_{m}^{m}d - J_{dabm}J_{m}^{m}c + J_{cdam}J_{m}^{m}b - J_{cdm}J_{m}^{m}a).\]

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