Recent results of BRST quantization on inner product spaces are reviewed. It is shown how relativistic particle models may be quantized with finite norms and that the relation between the operator method and the conventional path integral treatments is nontrivial.
BRST quantization within the operator formulation has the following ingredients: One starts with a nondegenerate state space, $\Omega$, and projects out a physical subspace by the condition $Q\Omega_{ph} = 0$ where $Q$ is the BRST charge operator which must be nilpotent $Q^2 = 0$. Then one makes $\Omega_{ph}$ nondegenerate by dividing out $Q\Omega$ from $\Omega_{ph}$. When this procedure is applied to relativistic particle models one usually derives covariant field equations (see e.g. [1]). However, when one applies the corresponding path integral formalism then one usually derives propagators [2, 3, 4]. It is therefore pertinent to ask how the two formalisms are related. Another question one may ask is whether or not it is possible to check unitarity of relativistic particle models, i.e. is it possible to calculate norms of physical states and check whether or not they are positive? In this talk I shall show that the answer to these questions is satisfactory provided the original state space $\Omega$ is chosen to be an inner product space.

In [5] a general framework for BRST quantization was proposed and analysed. In this framework the original state space $\Omega$ needs not be an inner product space. However, when it is not, one is forced to consider its dual state space $\Omega'$. (One needs the finite bilinear forms, $|'\langle u'|u\rangle| < \infty$ [5].) This framework was inspired by conformal field theory methods in string theory where the zero modes are treated in this asymmetric fashion. The BRST condition leads here to two different physical state spaces: $Q|ph\rangle = 0$ and $\langle ph|Q = 0$ where $|ph\rangle \in \Omega_{ph} \subset \Omega$ and $|ph\rangle' \in \Omega'_{ph} \subset \Omega'$. These conditions may be solved by a bigrading [5, 6] which means that $Q|ph\rangle = 0$ may be replaced by

$$\delta|ph\rangle = d|ph\rangle = 0$$

where

$$Q = \delta + d, \quad \delta^2 = d^2 = [\delta, d]_+ = 0$$

However, $\langle ph|Q = 0$ implies

$$\delta^\dagger|ph\rangle' = d^\dagger|ph\rangle' = 0$$
which means that
\[ \Omega'_{ph} \neq \Omega_{ph} \] (4)
when \( d^\dagger \neq \delta, d \). The problem is therefore how to make sense of \( \Omega'_{ph} \) when (4) is valid.

In several models it was possible to interpret \( \langle ph|ph \rangle \) as a finite wave function [1]. However, no norms are derivable although an effort was made in [3].

The above framework is considerably simplified when \( \Omega \) is required to be an inner product space. In this case \( \Omega' = \Omega \) which implies \( \Omega'_{ph} = \Omega_{ph} \). This simplification is, however, only possible under certain conditions. They are (see e.g. [7])

1. Nontrivial states in \( \Omega_{ph} \) has ghost number zero.

2. The number of (hermitian) gauge generators (constraints) must be even.

3. \( Q \) must be possible to decompose as
\[ Q = \delta + \delta^\dagger, \quad \delta^2 = [\delta, \delta^\dagger]_+ = 0 \] (5)

In order to see what the last two conditions require I consider a general bosonic gauge theory. The BRST charge may then be written in the BFV form [8]
\[ Q = \psi_a \eta^a - \frac{1}{2} i U_{bc}^a \pi_a \eta^b \eta^c - \frac{1}{2} i U_{ab}^b \eta^a + \bar{P}_a \pi^a \] (6)
where \( \psi_a, \ a = 1, \ldots, m \) are hermitian bosonic gauge generators (constraints) satisfying the Lie algebra
\[ [\psi_a, \psi_b]_- = i U_{ab}^c \psi_c \] (7)
where \( U_{ab}^c \) are the structure constants. \( \eta^a \) and \( \bar{\eta}^a \) are ghost and antighost and \( \pi_a \) is the conjugate momentum to the Lagrange multiplier \( u^a \). They are hermitian and satisfy the algebra (nontrivial part)
\[ [\eta^a, P_b]_+ = [\bar{\eta}^a, \bar{P}_b]_+ = \delta^a_b, \quad [\pi_a, u^b]_- = -i \delta^b_a \] (8)

We notice that the inclusion of dynamical Lagrange multipliers will always ensure that the second condition above is satisfied, since we have an even number of constraints: \( \psi_a = 0, \pi_a = 0 \). Recently [3, 10], I have shown that it is always possible to
decompose the BRST charge (3) according to (5) which means that also condition three above is satisfied for the BRST charge (3). The general procedure to do this is to look for a unitary transformation which preserves the ghost number:

$$\eta, \mathcal{P}, \bar{\eta}, \bar{\mathcal{P}}, \ldots \rightarrow \eta', \mathcal{P}', \bar{\eta}', \bar{\mathcal{P}}', \ldots$$

Typically it involves the Lagrange multipliers [9] or a gauge fixing operator $\chi^a$ [10]. Define the complex ghosts $c^a$ and $k_a$ by

$$c^a \equiv \frac{1}{2}(\eta^a - i \bar{\mathcal{P}}^a), \quad k_a \equiv \mathcal{P}_a' - i \bar{\eta}'_a$$

They satisfy

$$[k_a, c^b] = \delta^b_a$$

Now if we define $\delta$ by

$$\delta \equiv [c^a k_a, Q]$$

then $Q = \delta + \delta^\dagger$. However, this expression is only nilpotent if also

$$[k_a^+, c^a, \delta] = 0$$

This condition selects possible unitary transformations (3). The solutions found in [9, 10] all had the form

$$\delta = c^a \phi_a = \phi'_a c^a$$

where $\phi_a$ are nonhermitian operators which satisfy the same algebra as $\psi_a$ i.e.

$$[\phi_a, \phi_b] = i U_{ab} \phi_c$$

(In some cases $\phi_a$ may be chosen to be abelian [10].) The BRST condition $Q|\psi_h\rangle = 0$ may now be solved by a bigrading which yields

$$\delta|\psi_h\rangle = 0, \quad \delta^\dagger|\psi_h\rangle = 0$$

One may show that all solutions except possibly some zero norm states also are solutions of

$$c^a|\psi_h\rangle = \phi_a|\psi_h\rangle = 0$$
or
\[ c^{\dagger a}|ph\rangle = \phi^t_a|ph\rangle = 0 \] \hspace{1cm} (18)

Both these sets of equations have nontrivial solutions only for ghost number zero according to the analysis in [11]: The nontrivial solutions of (17) satisfy
\[ c^a|ph\rangle = k_a|ph\rangle = 0 \] \hspace{1cm} (19)

while (18) implies
\[ c^{\dagger a}|ph\rangle = k^{\dagger a}|ph\rangle = 0 \] \hspace{1cm} (20)

both of which requires \( |ph\rangle \) to have ghost number zero. Notice that (17) and (18) are consistency conditions to (19) and (20) since \( \phi_a = [Q,k_a]_+ \) and \( \phi^t_a = [Q,k^{\dagger a}]_+ \).

In [9, 10] general solutions of (17) and (18) were derived all of the form
\[ e^{\alpha[\rho,Q]_+}|\Phi\rangle \] \hspace{1cm} (21)

where \( \alpha \) is a real constant different from zero and where \( |\Phi\rangle \) is a simple BRST invariant state with ghost number zero. Two cases have been found:

1. ) \( \rho = \mathcal{P}_a v^a \) with \( \eta^a|\Phi\rangle = \bar{\eta}^a|\Phi\rangle = \pi_a|\Phi\rangle = 0 \)

2. ) \( \rho = \bar{\eta}_a \chi^a \) with \( \mathcal{P}^a|\Phi\rangle = \bar{\mathcal{P}}^a|\Phi\rangle = (\psi_a + \frac{1}{2}iU^b_{ab})|\Phi\rangle = 0 \)

where \( \chi^a \) is a hermitian gauge fixing operator which must be such that \([\chi^a,\psi_b] \) has an inverse. These general solutions are obtained in a purely algebraic way. They are therefore formal. However, if the quantization is such that \( \Omega \) is an inner product space then they must also belong to an inner product space. Thus, although (21) implies
\[ |ph\rangle = |\Phi\rangle + Q|\cdot\rangle \] \hspace{1cm} (22)

the states \( Q|\cdot\rangle \) cannot be divided out since \( |\Phi\rangle \) is not an inner product state while \( |ph\rangle \) must be. The second case above provides for a way to make the solutions of a
Dirac quantization consistent with an inner product space since $|\Phi\rangle$ is a solution to a Dirac quantization there and $\langle\Phi|e^{a[ρ,Q]}|\Phi\rangle$ is finite.

In order to find out what quantization rules have to be used in order for (21) to belong to an inner product space I consider first the trivial case when the gauge symmetry is abelian and $\psi_a = p_a$ where $p_a$ is a conjugate momentum operator to some coordinate operator $x^a$ \[12\]. The BRST charge is here

$$Q = p_a \eta^a + \bar{P}^a \pi_a = \delta + \delta^\dagger$$  \tag{23}$$

where $\delta = c^\dagger a \phi_a$ where in turn

$$c^a = \frac{1}{2}(\eta^a - i\bar{P}^a), \quad \phi_a = p_a - i\pi_a$$  \tag{24}$$

(Actually this case has been solved in its full generality in \[13\].) If e.g. the conditions (16) allow for (17) then the solutions are

$$|ph\rangle = c^\dagger a |\cdot\rangle, \quad \phi^\dagger_a |\cdot\rangle$$  \tag{25}$$

which all are zero norm states except for the vacuum state $|0\rangle$ which satisfies

$$k_a |0\rangle = \xi^a |0\rangle = 0$$  \tag{26}$$

where

$$k_a = P_a - i\bar{\eta}_a, \quad \xi^a = \frac{1}{2}(ix^a - v^a), \quad [\xi^a, \phi^\dagger_b]_- = \delta^a_b$$  \tag{27}$$

The diagonal basis of this Fock space is spanned by

$$a_a = \frac{1}{2}(\xi^a + \phi_a), \quad b_a = \frac{1}{2}(\xi^a - \phi_a)$$  \tag{28}$$

where $a_a^\dagger$ spans positive metric states and $b_a^\dagger$ indefinite ones. Obviously the quantization should be such that the unphysical states have a basis consisting of half of positive metric states and half of indefinite ones.

In order to understand what this result means for spectral bases let me consider the hermitian coordinate and momentum operators $Q$ and $P$ satisfying

$$[Q, P]_- = i$$  \tag{29}$$
Real spectral decompositions requires positive definite states since

$$Q|q\rangle = q|q\rangle, \quad \langle q|q'\rangle = \delta(q - q')$$  \hspace{1cm} (30)

implies

$$\langle \phi|\phi \rangle = \int dq|\phi(q)|^2 > 0$$  \hspace{1cm} (31)

There is also the possibility to use imaginary spectral decompositions as Pauli showed in [14]:

$$Q|iq\rangle = iq|q\rangle, \quad \langle iq|iq'\rangle = \delta(q - q'), \quad \langle -iq| = (|iq\rangle)^\dagger$$  \hspace{1cm} (32)

which requires indefinite metric states

$$\langle \phi|\phi \rangle = \int dq\phi^*(-q)\phi(q) = \pm C, \quad C > 0$$  \hspace{1cm} (33)

where the sign depends on the parity $\phi(-q) = \pm \phi(q)$.

Consider now the general solution (21) with $\rho = \mathcal{P}_a v^a$ for the trivial case (23) i.e.

$$|ph\rangle = e^\alpha[p\mathcal{Q}]|\Phi\rangle = e^\alpha(p_a v^a + i\mathcal{P}_a v^a)|\Phi\rangle$$  \hspace{1cm} (34)

It has the norm

$$\langle ph|ph \rangle = \langle \Phi|e^{2\alpha[p\mathcal{Q}]}|\Phi\rangle = \langle \phi| \pi\langle 0|e^{2\alpha p_a v^a}|0\rangle \pi|\phi \rangle$$  \hspace{1cm} (35)

The condition that this is a finite expression leads to the quantization rule:

**Rule 1:** Lagrange multipliers must be quantized with opposite metric states to the unphysical variables which the gauge generator $\psi_a$ eliminates.

For (35) it implies

$$\langle ph|ph \rangle = \int d^n p d^n v e^{i\alpha p_a v^a} |\phi(p)|^2 = \frac{1}{\alpha^m}|\phi(0)|^2$$  \hspace{1cm} (36)
It also implies that there must be equally many positive as indefinite metric oscillators in accordance with the result above. Thus, this rule is at least true for the trivial case.

Reconsider the trivial case \( (23) \)

\[
Q = p_a \eta^a + \bar{P}^a \pi_a
\]  

(37)

Obviously one may also make the choice to interpret \( \eta^a \) as gauge generator and \( \bar{\eta}_a \) as Lagrange multiplier, and \( p_a \) and \( \pi_a \) as bosonic ghost and antighost respectively. Finiteness of the norms imply then the quantization rule: 

**Rule 2:** Bosonic ghosts and antighosts must be quantized with opposite metric states.

In addition finiteness requires sometimes that the range of the Lagrange multipliers is restricted to the group manifold. (Notice that \( \langle ph | ph \rangle = \langle \Phi | e^{[\rho, Q]} | \Phi \rangle = \langle \Phi | e^{\psi_a \psi^a} \cdots | \Phi \rangle \)

where \( e^{\psi_a \psi^a} \) is a finite group transformation.)

The restriction to inner product spaces leads obviously to severe restrictions of the allowed quantization. One may therefore question the possibility to quantize relativistic particles and strings in a manifestly Lorentz covariant way on inner product spaces. The issue here is whether or not one may represent hermitian manifestly Lorentz covariant coordinate and momentum operators \( X^\mu \) and \( P^\mu \) satisfying \( [X^\mu, P^\nu]_+ = i\eta^{\mu\nu} \) in a satisfactory way. The possible representations are given in the table below \[15\]

| Spectra       | Lorentz covariant basis | Hermitian inner products |
|---------------|-------------------------|--------------------------|
| real          | no, positive            | yes                      |
| imaginary \( x^0, p^0 \) | yes, indefinite         | yes                      |
| real          | yes, imaginary          | no                       |

The last possibility is unsatisfactory and will not be considered. Thus, we have a choice between a manifestly covariant spectrum or basis. (Off-shell states are described by Euclidean spectra and a Lorentz covariant basis as we shall see.)
Consider first the free spinless particle \([12]\). It is described by the mass shell condition. In the corresponding BRST formulation it is described by the BRST charge operator

\[
Q = \frac{1}{2}(P^2 + m^2)\eta + \pi \bar{P} = c^\dagger \phi + \phi^\dagger c
\]  

(38)

where

\[
\phi = \frac{1}{2}(P^2 + m^2) - i\pi, \quad c = \frac{1}{2}(\eta - i\bar{P})
\]  

(39)

Thus, the BRST condition \(Q|ph\rangle = 0\) leads to the possible solutions

\[
|ph\rangle_{\pm} = e^{\pm \frac{1}{2}(P^2 + m^2)v_{\pm}\bar{P}}|\Phi\rangle
\]  

(40)

where \(|\Phi\rangle\) satisfies

\[
\pi|\Phi\rangle = \eta|\Phi\rangle = \bar{\eta}|\Phi\rangle = 0
\]  

(41)

Its formal norm is

\[
\pm \langle ph|ph\rangle_{\pm} \propto \langle \phi|e^{\pm(P^2 + m^2)v}|\phi\rangle
\]  

(42)

where all ghost dependence has been eliminated. The quantization rule 1 allows for two possibilities:

Case 1: \(X^0, P^0\) has real spectra and \(\pi, v\) imaginary. This implies

\[
\pm \langle ph|ph\rangle_{\pm} \propto \int d^4p du e^{\pm i(p^2 + m^2)u} |\phi(p)|^2 = 2\pi \int d^4\delta(p^2 + m^2)|\phi(p)|^2 > 0
\]  

(43)

Case 2: \(X^0, P^0\) has imaginary spectra and \(\pi, v\) realy. This implies

\[
\pm \langle ph|ph\rangle_{\pm} = \int d^4p du e^{\pm i(p^2 - m^2)v} \phi^*(-p^0, p)\phi(p^0, p) = \\
= |_\text{finite only if } v \in (0, \infty) \text{ or } (-\infty, 0) = \\
= \int d^4p \frac{\phi^*(-p^0, p)\phi(p^0, p)}{p^2 + m^2}
\]  

(44)

The last norm is only positive if \(\phi(p)\) has even parity. A Lorentz covariant way to assure positivity is to impose invariance under strong reflection \(p^\mu \rightarrow -p^\mu\) on
the original state space $\Omega$. Notice that the measure $d^4p$ is Euclidean in (43). The
Euclidean propagator is of the form $\langle ix|e^{\rho Q}|ix'\rangle$.

The BRST formulation of the worldline supersymmetric free massless spin-$\frac{1}{2}$
particle involves a BRST charge of the form

$$Q = P^2\eta + P \cdot \gamma c + \mathcal{P}c^2 + \pi\mathcal{P} + \kappa\bar{k}$$

(45)

where the variables satisfy the following (anti-)commutation relations (the nonzero
part):

$$[\gamma^\mu,\gamma^\nu]_+ = -2\eta^{\mu\nu}, \quad [X^\mu, P^\nu]_+ = i\eta^{\mu\nu}, \quad [\pi, v]_- = -i, \quad [k, \lambda]_+ = 1,$$

$$[\mathcal{P}, \eta]_+ = 1, \quad [\mathcal{P}, \bar{\eta}]_+ = 1, \quad [k, c]_- = -i, \quad [\bar{k}, \bar{c}]_- = -i$$

(46)

where $k, c$ are bosonic ghosts and $\bar{k}, \bar{c}$ the corresponding antighosts, $\lambda$ is a fermionic
Lagrange multiplier and $\kappa$ its conjugate momentum. $\eta^{\mu\nu}$ is a space-like Minkowski
metric. Notice that

$$[P \cdot \gamma, P \cdot \gamma]_+ = -2P^2$$

(47)

is the algebra of the world-line supersymmetry. In the matrix representation $\gamma^\mu$ is
turned into the ordinary Dirac gamma matrices. The BRST charge (45) may also
be written as $Q = \delta + \delta^\dagger$ where [12]

$$\delta = a^\dagger D + \sigma^\dagger \phi, \quad [D, D]_+ = -2\phi$$

(48)

where in turn

$$a \equiv \frac{1}{2}(c - i\bar{k}), \quad \sigma \equiv \frac{1}{2}(\eta - i\mathcal{P} - i\lambda c),$$

$$D = P \cdot \zeta + \lambda\pi - ik\omega - \omega^\dagger a, \quad \phi = P^2 - i\pi$$

(49)

A formal algebraic solution is [12]

$$|ph\rangle_\pm = e^{\pm[\rho Q]}|\Phi\rangle, \quad \rho = \mathcal{P}v + k\lambda$$

(50)

where

$$c|\Phi\rangle = \eta|\Phi\rangle = \pi|\Phi\rangle = \kappa|\Phi\rangle = 0$$

(51)
The formal norm is
\[ \pm \langle ph|ph \rangle_\pm = \langle \Phi | e^{\pm 2iP^2 v+ip\cdot \gamma \lambda+k\bar{k}+i\bar{P}\cdot \gamma \lambda P} |\Phi \rangle \] (52)

The quantization rule 2 reduces this expression to
\[ \pm \langle ph|ph \rangle_\pm \propto '(\langle \Phi | e^{\pm 2iP^2 v\cdot \gamma \lambda P} |\Phi \rangle)' \] (53)

where \(|\Phi \rangle'\) is equal to \(|\Phi \rangle\) without the bosonic ghost part. At this point there are again two cases:

**Case 1:** \(X^0, P^0\) has real spectra and \(\pi, v\) imaginary. This implies
\[ \pm \langle ph|ph \rangle_\pm \propto \int d^4 \delta(p^2) \bar{\psi}(p) p^5 \psi(p) \] (54)
which is not positive without further projections.

**Case 2:** \(X^0, P^0\) has imaginary spectra and \(\pi, v\) real. This implies
\[ \pm \langle ph|ph \rangle_\pm \propto \int d^4 p \int d^4 p' \bar{\psi}(-p^0, P) \frac{p^5 \gamma_5}{p^2 + m^2} \psi(p^0, P) \] (55)
which also is not positive without further projections.

There is an \(O(2)\) extended worldline supersymmetric model for a massless spin-1 particle [16, 1]. Its gauge generators are
\[ P^2, \ P \cdot \gamma_1, \ P \cdot \gamma_2, \ \gamma_1 \cdot \gamma_2 \] (56)
where \(\gamma_i\) satisfies \( [\gamma_1, \gamma_2]_+ = 0 \) and \( [\gamma_i, \gamma_1]_+ = -2\eta^{\mu\nu} \). This is a noncanonical theory since there exists no gauge fixing to \(\gamma_1 \cdot \gamma_2\). Within a BRST quantization on inner product spaces this implies that the physical state space will contain ghost excitations and will not be positive. A remedy of this defect is to first restrict the original state space by a condition of the form \(\gamma_1 \cdot \gamma_2 \bar{\Omega} = 0\). A BRST treatment as above [12] yields then in case 1: \(\langle ph|ph \rangle \propto \int d^4 p \delta(p^2) A_2^2(p) > 0\). Case 2 yields on the other hand Euclidean norms and propagators in agreement with the path integral treatment in [8].

Above we have demonstrated that we are able to check unitarity of models for relativistic particles with our general solutions, *i.e.* we can explicitly check whether
or not the physical norms are positive. Another issue raised in the introduction concerned the relation between path integrals and the operator quantization. This relation turns out to be nontrivial \[17\]. To see this I consider the time evolution of the states \( |\Phi\rangle \). Let \( H_0 \) be a BRST invariant Hamiltonian. I have then (\([\rho, H_0]\) is assumed to be BRST invariant)

\[
\langle p\hat{h}', t'|p\hat{h}, t \rangle = \langle p\hat{h}' | e^{-i(t'-t)H_0} | p\hat{h} \rangle = \\
= \langle \Phi' | e^{\alpha[\rho, Q]} e^{-i(t'-t)H_0} e^{\alpha[\rho, Q]} | \Phi \rangle = \langle \Phi' | e^{-i(t'-t)H_0 + 2\alpha[\rho, Q]} | \Phi \rangle
\]

from \(21\). \( \alpha \) is a real constant different from zero. Since \(57\) is independent of the value of \( \alpha \) I may set

\[
2\alpha = \pm (t' - t)
\]

for \( t' \neq t \). Eqn \(57\) may then be interpreted as if the time evolution of \( |\Phi\rangle \) is in terms of a complex Hamiltonian. The conventional path integral expressions are still obtained if \( H_0 \) is represented by a real function and \([\rho, Q]\) by an imaginary one, a condition which then governs the possible choices of the quantization of the system. The latter condition requires the Lagrange multipliers to be quantized in consistency with an imaginary spectral decomposition which is in agreement with the quantization rule 1 above. Bosonic ghosts and antighosts must be quantized with opposite metric states in agreement with rule 2. For fermionic ghosts and antighosts one must choose real and imaginary odd Grassmann spectra respectively or vice versa, a choice one always may do without affecting the metric of the states. Notice also that fermionic Lagrange multipliers must be quantized with imaginary odd spectra. If one uses real spectra for bosonic Lagrange multipliers, which was proposed as a possible choice in rule 1, then one finds a complex Hamiltonian also in the path integrals except for the trivial case. Thus, the Euclidean treatment of the relativistic particles above corresponds to complex Hamiltonians in the path integrals. However, when they are analytically continued to the Minkowski space they corresponds to real Hamiltonians apart from the \( i\epsilon \)-prescription in the propagators. It is this last case that one usually makes use of in the path integral treatments.
(\(\epsilon\) is then introduced for convergence reason) although a treatment according to case 1 seems more natural mathematically. Finally one may notice that the two cases of the general solutions (21) yield a precise connection between the choice of gauge fixing function \(\rho\) and the imposed boundary conditions in the path integral a connection which is essentially in agreement with those normally made [18].

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