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ABSTRACT

This letter studies the Sp(2) covariant quantisation of gauge theories. The geometrical interpretation of gauge theories in terms of quasi principal fibre bundles $Q(M_S, G_S)$ is reviewed. It is then described the Sp(2) algebra of ordinary Yang-Mills theory. A consistent formulation of covariant lagrangian quantisation for general gauge theories based on Sp(2) BRST symmetry is established. The original $N = 1$, ten dimensional superparticle is considered as an example of infinitely reducible gauge algebras, and given explicitly its Sp(2) BRST invariant action.

Key-words: Gauge theories; BRST symmetry; Fibre bundles; Covariant quantisation

1. Gauge theories in terms of quasi-principal fibre bundles.

Gauge theories have a nice geometrical interpretation in terms of connections on a principal fibre bundle (pfb) $P(M,G)$, where M is the base space-time manifold and G is the gauge group [1,2, 3, 4]. However, quantisation of gauge theories requires the introduction of fields (c^n_m, π^n_m). It would be then desirable to have a formalism where those extra fields fit into some representation of a larger group and all the fields are components of a superfield. This is a step in the direction of recovering a geometrical interpretation of quantum gauge theories. The main ingredients in the construction of geometrical quasi-principal fibre bundles (qpfb) are a space-time base manifold M, a gauge group G, an extended superspace manifold M_S which is obtained by adding two extra Grassmann variables θ^a $(a = 1,2)$ to M, in the case of $Sp(2)$ symmetry, and a supergroup G_S. The construction is performed basically in three steps [1,2, 3]. It starts with a pfb $P(M,G)$ and extend the gauge group G to a supergroup G_S. The composition of G with a Grassmann algebra B prolongs $P(M,G)$ to a pfb $P'(M,G_S)$. The most general supergroup G_S can be represented in matrix form. In particular, $OSp(N/M)$ groups are represented by block matrices of the form

$$\begin{pmatrix} A & E \\ C & D \end{pmatrix} \tag{1.1}$$

where A, D are $(N \times N)$ and $(M \times M)$ matrices whose elements are taken from the even part of the Grassmann algebra B constructed over a complex vector space W, whilst E, C are $(N \times M)$ and $(N \times M)$ rectangular matrices whose elements belong to the odd part of B. Next, it is enlarged the base space manifold M to a superspace M_S in $P'(M,G_S)$ by adding Grassmann variables. At this stage, a pfb $P''(M_S,G_S)$ is obtained. Finally, the pfb $P''(M_S G_S)$ is transform into a quasi-principal fibre bundle $Q(M_S,G_S)$. For instance, given a one-form valued function $\alpha(x) = A_\mu dx^\mu$ on M this induces a connection ω on the pfb $P(M,G)$. The a one-form valued function α' on M_S is found by

$$\alpha'(x,\theta^a) = g^{-1}A_\mu dx^\mu g + g^{-1}dg \tag{1.2}$$

where $g = g(x^\mu, \theta^a)$, $(a = 1,2,\ldots)$ which induces a connection ω' on the qpfb $Q(M_S,G_S)$ [1,2,3,4].

2. The $Sp(2)$ BRST Algebra of Yang-Mills Theory.

It has been realized for some time [5,6,7,8] that a geometrical construction can be useful for the discussion of BRST and anti-BRST symmetry. The idea is to use a superspace with coordinates $Z^M = (x^\mu, \theta^a)$, where $(a = 1,2)$ and θ^a is an anti-commuting scalar coordinate and the BRST generators s^a are realized as differential operators on superspace, $s^a = \frac{\partial}{\partial\theta^a}$, so that $s^a s^b + s^b s^a = 0$ holds automatically [1]. For example, in Yang-Mills theory the gauge potential A^i_μ and the Faddeev-Popov ghost $(c^a)^i$ (where i is an adjoint group

[1] It is usually defined a bosonic operator $\sigma = \frac{1}{2}\epsilon_{ab}s^a s^b$ where ϵ_{ab} is the symplectic invariant form of $Sp(2)$, so that $\epsilon_{ab} = -\epsilon_{ba}$, $\epsilon^{ab}\epsilon_{bc} = \delta_{ac}$ and $\epsilon_{12} = 1$. The generator σ is invariant

index) can be combined into a super-gauge field $\mathcal{A}^i_M(Z)$ whose lowest order components are $\mathcal{A}^i_M(Z)|_{\theta^a=0} = (\mathcal{A}^i_\mu, \mathcal{A}^i_{\theta^a})|_{\theta^a=0} = (A^i_\mu, c^{ai})$. Then the standard Yang-Mills BRST transformations arise from imposing the constraints $\mathcal{F}^i_{\mu\theta^a}=0$, $\mathcal{F}^i_{\theta^a\theta^b}=0$ on the superfield strength $\mathcal{F}^i_{MN}$ [5,6]. This gives an elegant geometrical description of BRST and anti-BRST symmetry.

Let us review the construction of gauge theories in the superspace with coordinates $Z^M=(x^\mu,\theta^a)$, which gives a geometric formulation of $Sp(2)$ BRST symmetry. We consider matter fields $\Phi^i(x,\theta^a)$ and a gauge potential $\mathcal{A}^i_M(x,\theta^a) = \left(\mathcal{A}^i_\mu(x,\theta^a), \mathcal{A}^i_{\theta^a}(x,\theta^a)\right)$. These can be used to define a covariant derivative

$$\mathcal{D}_M\Phi^i = \partial_M\Phi^i - (T^k)^i{}_j\mathcal{A}^k_M\Phi^j \tag{2.1}$$

and the field strength

$$\mathcal{F}^i_{MN} = \partial_M\mathcal{A}^i_N - (-1)^{MN}\partial_N\mathcal{A}^i_M + f^i{}_{jk}\mathcal{A}^j_M\mathcal{A}^k_N \tag{2.2}$$

where $(-1)^{MN}$ is 1 unless both M and N are indices referring to anti-commuting coordinates, in which case it is -1. The gauge potential $\mathcal{A}_M$ contains more component fields than the physical gauge and ghost fields and so, as in supersymmetric theories, constraints should be imposed on the field strength $\mathcal{F}$. Appropriate constraints are [5,6]

$$\mathcal{F}_{\theta^a\theta^b}=0, \qquad \mathcal{F}_{\mu\theta^a}=0. \tag{2.3}$$

These can be written more explicitly as

$$\partial_\mu\mathcal{A}_{\theta^a} - \partial_{\theta^a}\mathcal{A}_\mu + [\mathcal{A}_\mu, \mathcal{A}_{\theta^a}] = 0 \tag{2.4}$$

$$\partial_{\theta^a}\mathcal{A}_{\theta^a} + \frac{1}{2}[\mathcal{A}_{\theta^a}, \mathcal{A}_{\theta^a}] = 0 \tag{2.5}$$

$$\epsilon^{ab}\Big(\partial_{\theta^a}\mathcal{A}_{\theta^b} + [\mathcal{A}_{\theta^a}, \mathcal{A}_{\theta^b}]\Big) = 0. \tag{2.6}$$

Defining the component expansions

$$\mathcal{A}_\mu(x,\theta^a) = A_\mu(x) + \theta^a\Lambda_{a\mu}(x) + \theta^a\theta^b\Omega_{ab\mu}(x) \tag{2.7}$$

$$\mathcal{A}^b_{\theta^a}(x,\theta^a) = c^b(x) + \theta^a\Upsilon^b_a(x) + \theta^a\theta^c\omega^b_{ac}(x), \tag{2.8}$$

however, it was found that if A_μ, c^a, π are identified with the gauge, ghost (anti-ghost) and auxiliary fields respectively then the supergauge fields have the expansions *

$$\mathcal{A}_\mu = A_\mu + \theta^a(s^aA_\mu) + \theta^a\theta^b(s^as^bA_\mu) \tag{2.9}$$

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under $Sp(2)$ and satisfies $s^a\sigma = 0$. The Sp(2) generators σ^i ($i=\pm,0$) and the fermionic charges s^a together form an algebra which is a contraction of $OSp(1,1/2)$ and denoted as $ISp(2)$ [9].

* In Ref [6], it was obtained explicitly a geometrical formulation of BRST and anti-BRST symmetries and given the field content of Λ, Υ, Ω and ω. The components in the expansion can also be read as conditions on the mapping of the coordinates ϕ^i of the fibres over $\{U_i\}$ (covering set of M_G) and expressed as cocycle conditions.

$$A^b_{\theta^a} = c^b + \theta^a(s^a c^b) + \theta^a\theta^c(s^a s^c c^b). \tag{2.10}$$

The BRST and anti-BRST generators s^a are then identified with the superspace differential operators ∂_{θ^a} [5,6], and the complete set of BRST and anti-BRST transformations are given by [10 ,11 ,12]

$$\begin{aligned} s^a\phi^i &= R^i_\alpha c^{\alpha a}, \qquad s^a c^{\alpha b} = \epsilon^{ab}\pi^\alpha - \frac{1}{2} f^\alpha_{\beta\gamma} c^{\beta a} c^{\gamma b}, \\ s^a \pi^\alpha &= \frac{1}{2} f^\alpha_{\beta\gamma}\pi^\beta c^{\gamma a} - \frac{1}{12}(f^\alpha_{\beta\gamma} f^\beta_{\delta\tau} + f^\alpha_{\delta\tau,i} R^i_\gamma) c^{\delta a} c^{\tau e}\epsilon_{eb} c^{\gamma b}, \end{aligned} \tag{2.11}$$

where the generators R^i_α for the gauge field A^i_μ read off from $s^a A^i_\mu = (D_\mu c^a)^i$, and π^α is an auxiliary field which connects ghost and antighost sectors.

For the matter fields $\Phi^i(x, \theta^a)$, we impose the constraint

$$\mathcal{D}_{\theta^a}\Phi^i = \partial_{\theta^a}\Phi^i - (T^k)^i{}_j A^k_{\theta^a}\Phi^j = 0 \tag{2.12}$$

which implies

$$\Phi^i = \psi^i + \theta^a(s^a\psi^i), \tag{2.13}$$

and the BRST and anti-BRST transformations again corresponds to translations in the θ^a direction with $\frac{\partial}{\partial\theta^a}$ realized as differential operators on the extended space manifold M_S.

3. Sp(2) formalism for General gauge theories.

Consider a general gauge theory with classical fields $A^i(x^\mu)$ $(i = 1, 2, \ldots, n)$ and classical action $S_0(A^i)$. The action is invariant under gauge transformations

$$\delta A^i = R^i_\alpha \xi^\alpha, \tag{3.1}$$

where ξ^α is the local gauge parameter. The Noether equations are given by

$$S_{0,i}\ R^i_\alpha = 0, \qquad \alpha = 1, 2, \ldots, m \qquad (0 \le m \le n), \tag{3.2}$$

and the generators of the gauge transformations satisfy

$$R^i_{\alpha,j} R^j_\beta - (-)^{\alpha\beta} R^i_{\beta,j} R^j_\alpha = -R^i_\gamma f^\gamma{}_{\alpha\beta} - S_{0,j} M^{ij}_{\alpha\beta}. \tag{3.3}$$

To construct a covariant lagrangian formalism for general gauge theories either with open $(M^{ij}_{\alpha\beta} \neq 0)$ or closed algebras, and based on $Sp(2)$ BRST symmetry, it is needed to enlarge the base manifold M to M_S. It is then defined a superspace M_S to include classical fields Φ^A and $Sp(2)$ doublets of anti-fields Φ^*_A, Φ^{**}_A [13,14,15]. The properties of these fields and anti-fields are $\epsilon(\Phi^A) = \epsilon_A$, $\epsilon(\Phi^*_A) = \epsilon_A + 1$, $\epsilon(\Phi^{**}_A) = \epsilon_A$, $gh(\Phi^*_A) = (-)^a - gh(\Phi^A)$ and $gh(\Phi^{**}_A) = -gh(\Phi^A)$. An extended Poisson superbracket is defined by

$$(F, G) = \frac{\delta_r}{\delta\Phi^A} F \frac{\delta_l}{\delta\Phi^*_A} G - (-)^{\epsilon_F\epsilon_G} \frac{\delta_r}{\delta\Phi^A} G \frac{\delta_l}{\delta\Phi^*_A} F, \tag{3.4}$$

where ϵ_F, ϵ_G denotes the Grassmann parity of the F, G functions on M_S, and left (right) derivatives are understood with respect to anti-fields (fields) unless otherwise stated. The extended anti-bracket (3.4) satisfies

$$\epsilon\Big((F,G)\Big) = \epsilon(F) + \epsilon(G) + 1 \tag{3.5}$$

$$gh\Big((F,G)^a\Big) = -(-)^a + gh(F) + gh(G), \qquad a = 1,2 \tag{3.6}$$

$$(F,G) = -(-)^{\epsilon_F \epsilon_G}(G,F) \tag{3.7}$$

and

$$(-)^{\epsilon_F}\epsilon_G\big((F,G),H\big) + [cycl.\ perm\ \ (F,G,H)] = 0. \tag{3.8}$$

A bosonic action functional $S = S(\Phi^A, \Phi^*_A, \Phi^{**}_A)$ is constructed on M_S. This action satisfy the following *generating equation*

$$\bar{\Delta}^a \exp^{\frac{i}{\hbar}S(\Phi^A,\Phi^*_A,\Phi^{**}_A)} = 0, \qquad (a = 1,2), \tag{3.9}$$

together with the boundary condition

$$S(\Phi^A, \Phi^*_A, \Phi^{**}_A)\Big|_{\Phi^*_A = \Phi^{**}_A = 0} = S(\Phi_A). \tag{3.10}$$

The operator $\bar{\Delta}^a$ is defined by

$$\bar{\Delta}^a = \Delta^a + (i/\hbar)V^a, \qquad \Delta^a = (-)^{\epsilon_A}\frac{\delta}{\delta\Phi^A}\frac{\delta_l}{\delta\Phi^*_{Aa}} \qquad V^a = \epsilon^{ab}\Phi^*_{Ab}\frac{\delta_r}{\delta\Phi^{**}_A}. \tag{3.11}$$

The algebra of operators (3.11) satisfy the important property [2]

$$\bar{\Delta}^{\{a}\bar{\Delta}^{b\}} = 0. \tag{3.12}$$

The solution to the generating equation (3.9) is given as a power series of the Planck constant

$$S(\Phi^A, \Phi^*_A, \Phi^{**}_A) = \sum_{n=0}^{+\infty}\hbar^n S_{(n)}, \tag{3.13}$$

where the classical approximation $S_{(0)}$ satisfies

$$\frac{1}{2}(S_{(0)}, S_{(0)})^a + V^a S_{(0)} = 0. \tag{3.14}$$

For a theory in which the $Sp(2)$ algebra closes off-shell the classical solution $S_{(0)}$ takes the form

$$S_{(0)} = S_0 + \Phi^*_{Aa}s^a\Phi^A + \frac{1}{2}\Phi^{**}_A\epsilon_{ab}s^a s^b\Phi^A + F_{AB}\epsilon_{ab}s^a\Phi^A s^b\Phi^B. \tag{3.15}$$

For more complicated theories like superparticles or superstrings, $S_{(0)}$ has terms of higher order in the fields Φ^*_A, Φ^{**}_A to compensate those terms which makes the $Sp(2)$ algebra to close on-shell. The classical solution to the generating equation is $Sp(2)$ BRST invariant under modified BRST generators $\bar{s}^a$ which satisfy $\bar{s}^a\bar{s}^b + \bar{s}^b\bar{s}^a = 0$.

[2] A supercommutative, associative algebra $\mathcal{A}$ equipped with an extended Poisson anti-bracket structure plus a nilpotent property it is known as a BV-algebra, or coboundary Gersterhaber algebra (CGA) [16].

4. Orthosymplectic Structure of the Original BSC Superparticle.

The original BSC superparticle S_{BSC} and further models are known to yield the same spectrum as that of $D=10$, $N=1$ super-Yang-Mills theory [17]. It is used here as an illustrative example to construct its $Sp(2)$ covariant lagrangian, since the model has an infinitely reducible algebra. The BSC superparticle action is given by [18]

$$S_0 = \int d\tau [p_\mu \dot{x}^\mu - i\theta \not{p}\dot{\theta} - \frac{1}{2}ep^2]. \tag{4.1}$$

This action describes a particle with world-line parametrized by τ moving through a ten-dimensional $N=1$ superspace with coordinates (x^μ, θ_A). The superparticle action S_{BSC} is invariant under a 10 dimensional super-Poincaré symmetry

$$\delta\theta = \epsilon, \qquad \delta x^\mu = i\epsilon\Gamma^\mu\theta, \tag{4.2}$$

together with world-line reparametrisations and a local fermionic symmetry

$$\begin{aligned} \delta\theta = \not{p}\kappa, \qquad \delta e = 4i\kappa\dot{\theta} + \dot{\xi}, \\ \delta x^\mu = i\theta\Gamma^\mu\not{p}\kappa + \xi p_\mu. \end{aligned} \tag{4.3}$$

The Grassmann spinor κ_A parametrizes the local symmetry while ξ parametrizes a linear combination of world-line diffeomorphisms and a local *trivial* local symmetry. To construct a covariant $Sp(2)$ orthosymplectic structure for this model, it is required the formalism of the previuos section since the classical infinitely reducible gauge algebra $\mathcal{A}$ closes on-shell. It is then defined a superspace M_S to include the classical fields $\Phi^A = (x^\mu, p_\mu, e, \theta_A)$ and $Sp(2)$ doublets of anti-fields Φ^*_A, Φ^{**}_A. The classical approximation $S_{(0)}$ which satisfies (3.14) is given by

$$S_{(0)} = S_{BSC} + S_1 + S_2 + S_3, \tag{4.4}$$

where S_{BSC} is the classical action of the original superparticle and S_1, S_2, and S_3 are

$$\begin{aligned} S_1 = \int d\tau [&\theta^*_a \not{p}\kappa^a_1 + e^*_a(4i\kappa^a_1\dot{\theta} + \dot{c}^a) + \kappa^*_{nab}((-)^n\not{p})(f^{ab}_c\kappa^c_{n+1} + \epsilon^{ab}\pi_n) \\ &+ x^*_{\mu a}(i\theta\gamma^\mu\not{p}\kappa^a_1 + p_\mu c) + c^*_{da}(-2if^{ad}_{rs}\kappa^r_1\not{p}\kappa^s_1 + \epsilon^{ad}\pi)], \end{aligned} \tag{4.5}$$

$$\begin{aligned} S_2 = \int d\tau [&\theta^{**}(-p^2\epsilon_{ab}f^{ab}_c\kappa^c_2) + e^{**}(-4i\epsilon_{ab}f^{ab}_c\kappa^c_2\not{p}\dot{\theta} + 2i\epsilon_{ab}\kappa^a_1\dot{\not{p}}\kappa^b_1) \\ &+ \kappa^{**}_{nr}(-p^2)(\epsilon_{ab}f^{br}_c f^{ac}_s\kappa^s_{n+2} + f^{br}_b\pi_{n+1}) \\ &+ x^{**}_\mu(ip^2\epsilon_{ab})(f^{ab}_c\theta\gamma^\mu\kappa^c_2 - \kappa^a_1\gamma^\mu\kappa^b_1) \\ &+ c^{**}_e(-4ip^2)(\epsilon_{ab}f^{be}_{rs}f^{as}_c\kappa^r_1\kappa^c_2 + f^{be}_{rb}\kappa^r_1\pi_1)], \end{aligned} \tag{4.6}$$

and

$$\begin{aligned} S_3 = \int d\tau \frac{1}{2}\epsilon^*_a[&\theta^*_b(-\kappa^c_2)(2f^{ab}_c + \epsilon^{ab}\epsilon_{rs}f^{rs}_c) \\ &+ x^*_{\mu b}(2i\kappa^a_1\gamma^\mu\kappa^b_1 + i\epsilon^{ab}\epsilon_{rs}\kappa^r_1\gamma^\mu\kappa^s_1 - i\theta\gamma^\mu\kappa^c_2(2f^{ab}_c + \epsilon^{ab}\epsilon_{rs}f^{rs}_c)) \\ &+ \kappa^*_{nAb}(-\kappa^c_{n+2}(\epsilon^{ab}\epsilon_{ps}f^{sA}_q f^{pq}_c + 2f^{Ab}_s f^{as}_c) - \pi_{n+1}(2\epsilon^{as}f^{Ab}_s + \epsilon^{ab}f^{pA}_p)) \\ &- 4ic^*_{Ab}(\kappa^r_1\kappa^c_2(\epsilon^{ab}\epsilon_{ps}f^{sA}_{rq}f^{pq}_c + 2f^{bA}_{rs}f^{as}_c) + \kappa^r_1\pi_1(\epsilon^{ab}f^{bA}_{rs} + \epsilon^{ab}f^{pA}_{qp}))]. \end{aligned} \tag{4.7}$$

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