Gauging of Chern-Simons $p$-Branes

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Abstract. The Chern-Simons membranes and in general the Chern-Simons $p$-branes moving in $D$-dimensional target space admit an infinite set of secondary constraints. With respect to the Poisson bracket these constraints form a closed algebra which contains classical $W_{1+\infty}$ algebra in $p$-dimensions as a subalgebra. Corresponding gauged theory in the phase-space is constructed in a Hamilton gauge as an analog of the ordinary $W$-gravity.
In the previous article [1] it was shown that in the Chern-Simons membrane case there always appears an infinite set of secondary constraints in contrast to the C-S string theory [2] in which there are two possibilities for the first class constraints: there is a finite or an infinite number of secondary constraints. There is also another, rather formal, possibility when second class constraints appear also (see Refs. 1 and 2).

When there appears an infinite set of secondary constrains for the C-S string, they satisfy an infinite algebra with respect to the Poisson bracket. This algebra contains as a subalgebra the classical (without central term) affine $SL(2, R)$ algebra, as well as the classical Virasoro algebra and their higher spin extensions which contain classical $W_{1+\infty}$ algebra [3]. Through this paper we are dealing only with classical infinite algebras.

In the case of C-S membrane the infinite set of constraints gives a linear realization of higher spin extended algebra in two-dimensions which contains affine $SL(2, R)$, Virasoro and $W_{1+\infty}$-algebras in two dimensions as subalgebras. We note, that any of these algebras can not be represented as a direct product of two infinite algebras in one dimension as in the two-dimensional conformal theory [3] and their $W_\infty$ extensions. We have to take into account also that in the C-S membrane theory we are dealing with two spatial dimensions while the time variable appears only as evolution parameter which is not the case on the ordinary 2D conformal theory. The generalization of the results for arbitrary C-S $p$-branes is straightforward. In that case we have higher spin extension of the affine $SL(2, C)$ in $p$-dimensions.

The polynomial Chern-Simons $p$-brane action was obtained in [3] from the topological $(p + 1)$-brane action (only for $D = p + 2$) [3] in the same way as in the ordinary local theory. In the paper [13] a generalization for an arbitrary space-time dimension was found.

In the present article the gauged C-S membrane theory is constructed in the Hamiltonian approach considering the Lagrange multipliers as a gauge fields with arbitrary spin. The $W_{1+\infty}$ transformation properties of the gauge fields are obtained. We note that the "No go theorem" in the case of spin $> 2$ (see Ref. [14]) does not take place because we are dealing with an infinite sequence of higher spin gauge fields. Through this paper we use

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1 In ordinary two-dimensional conformal theories we have two copies of Kac-Moody, Virasoro and $W$ algebras each of which acts only on one light-cone (holomorphic or antiholomorphic) coordinate, i.e. we have direct product of two algebras in one-dimension.

2 In the present article we use the terminology of the papers [1] and [2].
basis in which not all the constrains are independent and as a consequence there exists 
an additional symmetry of Stuckelberg type \[13\]. This symmetry allows us to exclude the 
corresponding gauge fields (with odd spin) by means of gauge fixing procedure. Integrating 
over the momentum variables we find $W_{1+\infty}$ gauge invariant action on the configuration 
space. We remaind that the Lagrangian approach to the $W$-gravity was considered in a 
lot of papers among the first of which are \[15\] and \[16\].

To proceed further we shall remind briefly some results from the papers \[8\] and \[13\], 
where in order to generalize the C-S $p$-brane to arbitrary space-time dimension the following 
notation were introduced:

$$X_{;A} = \partial AX, \quad X_{;a} = \partial aX$$

($a = 0, 1, 2, \ldots, p$). The polynomial Lagrangian for the C-S $p$-brane is given by $L = detX_{;A}^\mu$ 
which exists in $D = p + 2$ only. To extend this action for target space with arbitrary 
dimension, a generalized induced metric tensor is introduced

$$\tilde{g}_{AB} = X_{;A}^\mu X_{;B}^\nu \eta_{\mu\nu}, \quad (1)$$

where $\eta_{\mu\nu}$ is the pseudoeuclidean metric tensor. Formally replacemint of the ordinary 
induced metric tensor in the Nambu-Goto action with the generalized metric tensor given 
by Eq. (1) gives the action for the generalized C-S $p$-branes which lives on a target space 
with arbitrary dimension

$$S = \kappa \int d\tau d^p\sigma \sqrt{-\tilde{g}}. \quad (2)$$

It is easy to check that the action (2) obeis the same $p + 1$-variable diffeomorphisms 
invariance as the ordinary $p$-branes action. As a consequence of this invariance from the 
following first class primary constraints are obtained:

$$\phi_\perp = \mathcal{P}^2 + \kappa^2 det(X_{;u}X_{;v}) \approx 0,$$

$$\phi_j = \mathcal{P}X_{;j} \approx 0,$$

$$\phi_* = \mathcal{P}X \approx 0, \quad (3)$$

where $u, v = *, 1, \ldots, p$; $j = 1, \ldots, p$. Hence, there appears one additional primary con-
straint and moreover, the degree of the first constraint is higher by two degrees than the 
degree of the corresponding constraint in the ordinary $p$-branes. The latter is a consequence 
of the fact that the Lagrangian given by Eq. (2) and the constraints (3) are obtained from
the corresponding ordinary \((p + 1)\)-brane Lagrangian and \((p + 1)\)-brane constraints substituting \(\partial_{\sigma_{p+1}} X\) by \(X\). We recall, that the ordinary bosonic string has only two (bilinear) constraints while the C-S string has three primary constraints, one of which is of degree four with respect to \(X\).

The appearance of the constraint \(\phi_\ast\) shows us that some residual symmetry from the \(p + 2\)-variable diffeomorphisms (under which the action of the \(p + 1\)-brane is invariant) survives. As a consequence of the appearance of the constraint \(\phi_\ast\) there arise some secondary constraints too. In the C-S particle case we have only one secondary constraint, while in the C-S string case there are two possibilities: four (three primary constraints and one secondary constraint) first class constraints [3] or an infinite set of first class constraints [2].

We note that in the latter case not all of the constraints are independent if we deal with finite dimensional target space. Hence we have not dynamical degree of freedom. In that case the dynamical degrees of freedom can take place only if we have infinite dimensional target space.

For any C-S \(p\)-brane the canonical Hamiltonian vanishes identically, i.e.

\[
H_0 = \mathcal{P} \dot{X} - \mathcal{L} \equiv 0, \tag{4}
\]

which is a property of the ordinary \(p\)-brane theory also.

The analyze of the constraint algebra in the case of C-S membrane shows us that there is an infinite series of secondary constraints [1]. An appropriate choice of these constraints is the following:

\[
\begin{align*}
\Psi^{m,n} &= \left( \mathcal{P} \partial_{\sigma_1}^m \partial_{\sigma_2}^n \mathcal{P} \right) \approx 0, \\
\Phi^{m,n} &= \left( X \partial_{\sigma_1}^m \partial_{\sigma_2}^n X \right) \approx 0, \\
\Gamma^{m,n} &= \left( \mathcal{P} \partial_{\sigma_1}^m \partial_{\sigma_2}^n X \right) \approx 0, \quad (m, n = 0, 1, \ldots).
\end{align*} \tag{5}
\]

We note that, as it was mentioned above, if \(D\) is finite we have only a finite number of independent constraints [3]. However, when we are dealing with infinite dimensional target space it is easy to check that all the constraints \(\Gamma\) are independent as well as those of the constraints \(\Psi\) and \(\Phi\) for which \(m + n = 2k\) \((k = 0, 1, \ldots)\). To prove the latter statement we use the following identity

\[
(\partial_{\sigma_1}^m \partial_{\sigma_2}^n XY) = \sum_{p=0}^{m} \sum_{q=0}^{n} (-)^{m+n-p-q} \binom{m}{p} \binom{n}{q} \partial_{\sigma_1}^p \partial_{\sigma_2}^q (X \partial_{\sigma_1}^m \partial_{\sigma_2}^n Y) \tag{6}
\]
which consequend from the Laibniz formula.

Using the Eq. (6) we obtain that the constraints $\Psi$ and $\Phi$ with arbitrary odd spin can be represented in terms of the constrains with all underlying spins:

\[
(X X^{2k-l+1,l}) = \frac{1}{2} \sum_{p=0}^{2k-l+1} \sum_{q=0}^{l} \sum_{q+p \neq 0} (-)^{p+q+1} \binom{2k-l+1}{p} \binom{l}{q} \partial_{\sigma_1}^{p} \partial_{\sigma_2}^{q} (X X^{2k-l-p+1,l-q}),
\]

(7)

where $k, l = 0, 1, 2, \ldots$. To obtain the r.h.s. of the Eq. (7) for given $k, l$ only in terms of independent quantities we have to determine all of the even lower spin quantities from the corresponding equation and then to insert them in the r.h.s. of (7). In such a way we get

\[
(X X^{2K-l+1,l}) = \sum_{M,m} C_{M,m}^{K,l} \partial_{\sigma_1}^{2(K-M)-l+m+1} \partial_{\sigma_2}^{-m} (X X^{2M-m,m}),
\]

(8)

where the sumation over $m$ is from 0 to $\min(l,2M)$ and over $M$ is from 0 to $K+(m-l+1)/2$. The coefficient $C_{M,m}^{K,l}$ can be determined by the procedure described above.

Althout, only the constraints $\Psi$ and $\Phi$ for which $m+n=2k$ are independent for convenience we do not exclude the odd spin constraints, moreover, that all the constraints $\Gamma$ are independent (for $D = \infty$).

With respect to the Poisson bracket the constraints (5) form an infinite algebra which contains $W_{1+\infty}$ algebra in two-dimensions as a subalgebra [1]:

\[
\{\Gamma^{k,l}[f], \Gamma^{m,n}[h]\}_{PB} = \sum_{p=0}^{k} \sum_{q=0}^{l} \binom{k}{p} \binom{l}{q} \Gamma^{k+m-p,l+n-q} [f \partial_{\sigma_1}^{p} \partial_{\sigma_2}^{q} h] - \sum_{r=0}^{m} \sum_{s=0}^{n} \binom{m}{r} \binom{n}{s} \Gamma^{k+m-r,l+n-s} [h \partial_{\sigma_1}^{r} \partial_{\sigma_2}^{s} f].
\]

(9)

We note, that Eq. (9) contains as subalgebras also two copies of $W_{1+\infty}$ algebras in linear realization – one of which acts on $\sigma_1$ coordinate and the other one acts on $\sigma_2$ coordinate. The constraints $\Gamma^{k,0}$ and $\Gamma^{0,k}$ appear as generators of these transformations. In the general case $\Gamma^{k,l}$ can be considered as generators of generalized diffeomorphisms

\footnote{These $W_{1+\infty}$-algebras differ from the ordinary $W_{1+\infty}$-algebras because they are not mutually commuting.}
in two dimensional space. Indeed, the Poisson bracket of $\Gamma$ with the coordinate $X^\mu$ and with the momenta $P^\mu$ give the transformation laws for the phase space coordinates:

$$
\delta_{\Gamma}^{k,l} X^\mu = \{\Gamma^{k,l}[f], X^\mu\}_{PB} = -f \partial^k_{\sigma_1} \partial^l_{\sigma_2} X^\mu,
$$

$$
\delta_{\Gamma}^{k,l} P^\mu = \{\Gamma^{k,l}[f], P^\mu\}_{PB} = (-)^{k+l} \sum_{p=0}^{k} \sum_{q=0}^{l} \binom{k}{p} \binom{l}{q} \partial^p_{\sigma_1} \partial^q_{\sigma_2} f \partial^{k-p}_{\sigma_1} \partial^{l-q}_{\sigma_2} P^\mu. \tag{10}
$$

In the same way we obtain also:

$$
\delta_{\Phi}^{k,l} X^\mu = \{\Phi^{k,l}[f], X^\mu\}_{PB} = 0,
$$

$$
\delta_{\Phi}^{k,l} P^\mu = \{\Phi^{k,l}[f], P^\mu\}_{PB}
$$

$$
= -f \partial^k_{\sigma_1} \partial^l_{\sigma_2} X^\mu - (-)^{k+l} \sum_{p=0}^{k} \sum_{q=0}^{l} \binom{k}{p} \binom{l}{q} \partial^p_{\sigma_1} \partial^q_{\sigma_2} f \partial^{k-p}_{\sigma_1} \partial^{l-q}_{\sigma_2} X^\mu,
$$

$$
\delta_{\psi}^{k,l} X^\mu = \{\Psi^{k,l}[f], X^\mu\}_{PB}
$$

$$
= f \partial^k_{\sigma_1} \partial^l_{\sigma_2} P^\mu (-)^{k+l} \sum_{p=0}^{k} \sum_{q=0}^{l} \binom{k}{p} \binom{l}{q} \partial^p_{\sigma_1} \partial^q_{\sigma_2} f \partial^{k-p}_{\sigma_1} \partial^{l-q}_{\sigma_2} P^\mu,
$$

$$
\delta_{\psi}^{k,l} P^\mu = \{\Psi^{k,l}[f], P^\mu\}_{PB} = 0. \tag{11}
$$

Consequently, $\delta_{\Gamma}^{1,0}$ and $\delta_{\Gamma}^{0,1}$ are ordinary diffeomorphisms in two-dimensional space. We note, that the asymmetry which appears in the transformation laws of the coordinate $X$ and momentum $P$ is a consequence of the asymmetric choice of the constraint basis (5). Taking into account the identity (6) a more symmetric basis can be obtained for the constraints (5) by a simple redefinition

$$
\Lambda^{m,n} \rightarrow \tilde{\Lambda}^{m,n} = \sum_{p=0}^{m} \sum_{q=0}^{n} b^{mn}_{pq} \partial^p_{\sigma_1} \partial^q_{\sigma_2} \Lambda^{m-p,n-q},
$$

where $b$ are constants. By a suitable choice of $b$ the classical algebra (9) can be deformed to an algebra which admits diagonal central extension also, at least for the $W_{1+\infty}$ subalgebra.

Because of the vanishing of the canonical Hamiltonian given by Eq. (4) the first order action can be written in the form:

$$
S = \int d\tau d^2\sigma \left( P \ddot{X} - \alpha_{mn} \left( X X^{(m,n)} \right) - \beta_{mn} \left( P P^{(m,n)} \right) - \gamma_{mn} \left( P X^{(m,n)} \right) \right), \tag{12}
$$

where $U^{(m,n)} = \partial^m_{\sigma_1} \partial^n_{\sigma_2} U$. In order to gauge the action given by Eq. (12) we consider the lagrange multipliers $\alpha, \beta$ and $\gamma$ as fields depending on the evolution parameter $\tau$ also.

5
Then using the transformation laws for the phase-space coordinates given by Eqs. (10) and (11) with \( \tau \) depending parameters \( f \) we obtain the transformation laws for the gauge fields:

\[
\delta \Gamma^{\alpha}{_{mn}} = \sum_{k,l \geq 0} (-)^{r+s} \binom{k}{r} \binom{l}{s} \partial_{\sigma_1}^r \partial_{\sigma_2}^s \left( f_{kl} \alpha_{m-k+r,n-l+s} \right) + \sum_{k,l \geq 0} \sum_{r,s = 0} \binom{k}{r} \binom{l}{s} \alpha_{kl} \partial_{\sigma_1}^r \partial_{\sigma_2}^s f_{m-k+r,n-l+s},
\]

(13)

\[
\delta \Gamma^{\beta}{_{mn}} = -\sum_{k,l \geq 0} (-)^{k+l} \binom{k}{r} \binom{l}{s} \beta_{kl} \partial_{\sigma_1}^r \partial_{\sigma_2}^s \left( \gamma_{m-k+r,n-l+s} \right)
\]

(14)

\[
\delta \Gamma^{\gamma}{_{mn}} = -\dot{f}_{mn} - \sum_{k,l \geq 0} \sum_{r,s = 0} \binom{k}{r} \binom{l}{s} \gamma_{kl} \partial_{\sigma_1}^r \partial_{\sigma_2}^s f_{m-k+r,n-l+s},
\]

(15)

We note that the action (12) is invariant only with respect to the gauged \( W_{1+\infty} \) algebra in two dimensions. It is not invariant with respect to the local gauge transformations (11).

In order to write down the action (12) on the configuration space we exclude the momentum variables by means of the equation:

\[
\frac{\delta L}{\delta P} = 0.
\]

(16)

Incerting the Lagrangian from (12) into Eq. (16) we find

\[
\dot{X}_\mu - \alpha X_{\mu}^{(m,n)} - Q P_\mu = 0,
\]

(17)

where

\[
Q = \frac{1}{2} \sum_{m,n \geq 0} \left( \beta_{mn} \partial_{\sigma_1}^m \partial_{\sigma_2}^n + \sum_{p,q = 0} \partial_{\sigma_1}^p \partial_{\sigma_2}^q \beta_{mn} \partial_{\sigma_1}^{m-p} \partial_{\sigma_2}^{n-q} \right),
\]

(18)

and the derivatives act on the right. From Eq. (17) we obtain

\[
P_\mu = Q^{-1} \left( \dot{X}_\mu - \sum_{m,n \geq 0} \alpha_{mn} X_{\mu}^{m,n} \right).
\]

(19)
Incerting the momentum from (19) into (12) we find

\[
\mathcal{L} = \frac{\lambda}{2} \left( \dot{X}^2 - X_{\sigma_1}^2 - X_{\sigma_2}^2 \right) - \sum_{m,n \geq 0} \left( \bar{\alpha}_{mn} \dot{X}^{(m,n)} - \bar{\beta}_{mn} X^{(m,n)} - \bar{\gamma}_{mn} \dot{X}^{(m,n)} \right),
\]

(20)

where the kinetic term is separated formally. New gauge fields \(\bar{\alpha}, \ldots\) are introduced instead of the infinite series of the Lagrange multipliers \(\alpha, \ldots\) and their derivatives and \(\lambda\) is a gauge invariant field with spin 1. The explicit form of these functions in terms of \(\alpha, \ldots\) can be found by power decomposition of the operator \(Q^{-1}\).

In order to gauge the configuration space Lagrangian (20) we suppose that the multipliers \(\bar{\alpha}, \ldots\) are functions of the evolution parameter \(\tau\) and the spatial world-sheet coordinates \(\sigma\). From the invariance of the action (20) with respect to the local \(W_{1+\infty}\) transformations we obtain the transformation laws for the gauge fields:

\[
\begin{align*}
\delta \bar{\alpha}_{mn} &= -\lambda f_{mn} - \sum_{k,l \geq 0} \sum_{p,q = 0} \left( \begin{array}{c} k \\ p \end{array} \right) \left( \begin{array}{c} l \\ q \end{array} \right) \left( (-)^{k+l} \partial_{\sigma_1}^p \partial_{\sigma_2}^q \left( f_{kl} \bar{\alpha}_{m-k+p,n-l+q} \right) 
+ \bar{\alpha}_{k,l} \partial_{\sigma_1}^p \partial_{\sigma_2}^q f_{m-k+p,n-l+q} \right), \\
\delta \bar{\beta}_{mn} &= -\lambda \dot{f}_{mn} + \sum_{k,l \geq 0} \sum_{p,q = 0} \left( \begin{array}{c} k \\ p \end{array} \right) \left( \begin{array}{c} l \\ q \end{array} \right) \left( (-)^{k+l} \partial_{\sigma_1}^p \partial_{\sigma_2}^q \left( f_{kl} \bar{\beta}_{m-k+p,n-l+q} \right) 
+ \dot{f}_{kl} \bar{\alpha}_{m-k+p,n-l+q} \right) + \bar{\alpha}_{k,l} \partial_{\sigma_1}^p \partial_{\sigma_2}^q \dot{f}_{m-k+p,n-l+q} + \bar{\beta}_{kl} \partial_{\sigma_1}^p \partial_{\sigma_2}^q f_{m-k+p,n-l+q}, \\
\delta \bar{\gamma}_{mn} &= -\lambda \left( \partial_{\sigma_1}^2 + \partial_{\sigma_2}^2 \right) f_{mn} + 2 \left( \partial_{\sigma_1} \dot{f}_{k-1,l} + \partial_{\sigma_2} \dot{f}_{k,l-1} \right) + f_{m-2,n} - f_{m,n-2} \\
&+ \sum_{k,l \geq 0} \sum_{p,q = 0} \left( \begin{array}{c} k \\ p \end{array} \right) \left( \begin{array}{c} l \\ q \end{array} \right) \left( (-)^{k+l} \partial_{\sigma_1}^p \partial_{\sigma_2}^q \left( f_{kl} \bar{\gamma}_{m-k+p,n-l+q} \right) 
+ \dot{f}_{kl} \bar{\beta}_{m-k+p,n-l+q} \right) + \bar{\gamma}_{k,l} \partial_{\sigma_1}^p \partial_{\sigma_2}^q f_{m-k+p,n-l+q} \right). 
\end{align*}
\]

(21)

According to Eq. (8) all the quantities \(\dot{X} X^{m,n}\) (if \(D = \infty\)) as well as the quantities \((\dot{X} \dot{X}^{m,n})\) and \((XX^{m,n})\) for which \(m + n = 2k\) are independent. So the formula (8) shows that there exists an symmetry of the action (20) with respect of the transformations.
of Stukelberg type \[15\]

\[
\tilde{\alpha}_{2M-m+1,m} = u_{2M-m+1,m},
\]

\[
\tilde{\alpha}_{2M-m,m} = -\sum_{K,l} C^{K,l}_{M,m} \partial_{\sigma_1} 2(K-M)-l+m \partial_{\sigma_2}^{-m} u_{2K+1,l},
\]

\[
\tilde{\gamma}_{2M-m+1,m} = v_{2M-m+1,m},
\]

\[
\tilde{\gamma}_{2M-m,m} = -\sum_{K,l} C^{K,l}_{M,m} \partial_{\sigma_1} 2(K-M)-l+m \partial_{\sigma_2}^{-m} v_{2K+1,l},
\]

(22)

where \(u_{m,n}\) and \(v_{m,n}\) are arbitrary functions. This invariance allows us to choose the following gauge fixing

\[
\tilde{\alpha}_{2K-l+1,l} = 0,
\]

\[
\tilde{\gamma}_{2K-l+1,l} = 0.
\]

(23)

In this gauge the even spin quantities \((\dot{X} \dot{X} 2^{2K-l+1,l})\) and \((XX 2^{2K-l+1,l})\) are canceled in the action (20). Then the Lagrangian became

\[
\mathcal{L} = \frac{\lambda}{2} \left( \dot{X}^2 - X_{\sigma_1}^2 - X_{\sigma_2}^2 \right)
\]

\[
- \sum_{m,n \geq 0} \tilde{\beta}_{mn} (\dot{X} X^{m,n}) + \sum_{2K \geq l} \left( \tilde{\alpha}_{2K-l,l} (\dot{X} \dot{X} 2^{2K-l,l}) - \tilde{\gamma}_{2K-l,l} X 2^{2K-l,l} \right),
\]

(24)

which in the case \(D = \infty\) contains only independent quantities. In any other case only finite number of quantities survives.

At the end we note that, the Hamiltonian approach applied here loses the manifest Lorentz covariance and leads to Hamilton gauge. For instance, here appear three infinite sequences of gauge fields instead of one vector gauge field sequence that appears in the manifestly Lorentz covariant approach. The generalization for the case of arbitrary C-S \(p\)-branes is straightforward. In that case we have \(p+1\) infinite sequences of gauge fields i.e. one sequence of \(p+1\) vector potentials.

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