Non-relativistic strings and branes
as non-linear realizations of Galilei groups

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Abstract
We construct actions for non-relativistic strings and membranes purely as Wess-Zumino terms of the underlying Galilei groups.

1 Introduction
Recently, a closed non-relativistic (NR) string with a non-trivial spectrum of excitations was constructed [1, 2]. The construction was motivated by non-commutative open string (NCOS) theories [3, 4] in 1 + 1 dimensions [5]. Here we would like to elucidate the symmetries and the geometrical structure of the NR string. On the one hand, our goal will be to generalize the study of the free non-relativistic particle action as a Wess-Zumino (WZ) term of the ordinary Galilei group [6]. On the other hand, our analysis will parallel the study of the relativistic Green-Schwarz string action which contains a WZ term [7] predicated on the non-trivial third cohomology group of (N = 2 SuperPoincaré) / SO(9, 1) [8]. After treating the string case, we will then discuss the extension to non-relativistic d-branes.

2 Non-relativistic strings
The contraction of the Poincaré group associated with the NR string in n spacetime dimensions is obtained by letting c → ∞ after rescaling the coordinates

\[ x^0 \rightarrow cx^0 \equiv ct , \quad x^1 \rightarrow cx^1 \equiv cx , \quad x^a \rightarrow X^a , \quad a = \{2, \ldots, n-1\} \]  (1)
as well as rescaling the Poincaré generators

\[ P^0 \rightarrow H/c, \quad P^1 \rightarrow P/c \]
\[ J_{0a} \rightarrow cJ_{0a} \equiv cK_a, \quad J_{1a} \rightarrow cJ_{1a} \equiv cJ_a \]
\[ P_a \rightarrow P_a, \quad J_{01} \rightarrow J_{01} \equiv K, \quad J_{ab} \rightarrow M_{ab} \]  \hspace{1cm} (2)

The contracted algebra is then

\[ [M_{ab}, M_{cd}] = i(\theta_{ac}M_{bd} + \theta_{bd}M_{ac} - \delta_{ad}M_{bc} - \delta_{bc}M_{ad}) \]
\[ [K_a, M_{cd}] = i(\theta_{ad}K_c - \theta_{ac}K_d), \quad [K_a, K] = iJ_a \]
\[ [J_a, M_{cd}] = i(\theta_{ad}J_c - \theta_{ac}J_d), \quad [J_a, K] = iK_a \]
\[ [M_{ab}, P^c] = i(\delta_{ac}P_b - \delta_{bc}P_a) \]
\[ [K, H] = iP, \quad [K, P] = iH \]
\[ [K_a, H] = iP_a, \quad [J_a, P] = iP_a \]  \hspace{1cm} (3)

Applying the general techniques \[9\] for non-linear realizations of spacetime symmetries, as have been applied to relativistic membranes and supermembranes \[10\], we consider the coset element

\[ g = \exp (-iH + ixP) \exp \left( ix^a P_a + iv^b K_b + i\theta^a J_a \right) \]  \hspace{1cm} (4)

where \( H, \ P \) are the unbroken translations, and \( X^a (t, x), \ v^a (t, x), \ \theta^a (t, x) \) are Goldstone fields associated with the broken generators \( P^a, \ K^a, \ J^a \), respectively. The stability group is generated by the transverse rotations \( J^{ab} \), and by \( K \). The transformation properties of \( X^a (t, x), \ v^a (t, x), \ \theta^a (t, x) \) are given (up to rotations) by

\[ t' = t_0 + t \cosh v_0 + x \sinh v_0, \quad x' = x_0 + t \sinh v_0 + x \cosh v_0 \]
\[ v^a' = v_0^a + v^a \cosh v_0 + \theta^a \sinh v_0, \quad \theta^a' = \theta_0^a + v^a \sinh v_0 + \theta^a \cosh v_0 \]
\[ X^a' = X_0^a + (t_0 + t \cosh v_0 + x \sinh v_0) v_0^a - (x_0 + t \sinh v_0 + x \cosh v_0) \theta_0^a \]  \hspace{1cm} (5)

The Maurer-Cartan one-form \( \Omega = -ig^{-1}dg \) is

\[ \Omega = K_a dv^a + J_a d\theta^a + P_a (d v^a dt + d\theta^a dx + dX^a) - H dt + P dx \]
\[ \equiv K_a \omega_K + J_a \omega_J + P_a \omega_P - H \omega_H + P \omega_P \]  \hspace{1cm} (6)

From this, we construct the closed, invariant three-form

\[ \Omega_3 = \omega_K \wedge \omega_P \wedge \omega_P - \omega_J \wedge \omega_P \wedge \omega_H \]  \hspace{1cm} (7)

Moreover, it is easy to obtain a two-form “potential” \( \Phi_2 \) such that \( \Omega_3 = d\Phi_2 \), namely

\[ \Phi_2 = \frac{1}{2} \left( \theta^2 - v^2 \right) dt \wedge dx + v^a dX^a \wedge dx - \theta^a dX^a \wedge dt \]  \hspace{1cm} (8)

modulo addition of any closed or exact 2-form. \( \Phi_2 \) cannot be written in terms of the left-invariant one forms, so it is not invariant under the action of the group, and therefore the third cohomology group of the Galilei group as given by \(3\) is not trivial.

There are two immediate developments. (I) We can construct an extended algebra with new non-trivial commutation relations. In part, these would be given by

\[ [K_a, P_b] = i\delta_{ab} Z, \quad [K_a, J_b] = i\delta_{ab} W, \quad [P, K_a] = iN_a \]  \hspace{1cm} (9)
Note that the presence of the central charges $Z$ and $W$ break the $SO(1,1)$ invariance, since, for example, $[K_\alpha, P_\beta]=i\delta_{\alpha\beta}Z$ but $[J_\alpha, P_\beta]=0$. We will explore this further elsewhere \[1\].

(II) We can construct a WZ action given by

$$S_{WZ} = T \int \ast \Phi^2$$

where $\ast \Phi^2$ is the pullback of the two-form on the world sheet and $T = 1/(2\pi l_s^2)$ is the tension of the string. This WZ action takes the explicit form

$$S_{WZ} = T \int dt \wedge dx \left( v^a \partial_t X^a + \theta^a \partial_x X^a + \frac{1}{2} \left( \theta^2 - v^2 \right) \right)$$

(10)

If we fix the static gauge, $t=\tau$ and $x=\alpha \sigma$, $\alpha$ being a parameter, and we eliminate the nondynamical Goldstone fields, we get \[12\]

$$S_{WZ} = T \int dt dx \left( \frac{1}{2} (\partial_t X^a)^2 - \frac{1}{2} (\partial_x X^a)^2 \right)$$

(11)

The transverse coordinates are a collection of free, massless fields on the world-sheet.

If we compute the Noether charges associated with the transformations (5), we find for the central and topological elements in the gauge fixed form the formal expressions

$$Z = \int d\sigma \frac{\partial x}{\partial \sigma}, \quad W = \int d\sigma \frac{\partial x}{\partial \sigma}, \quad N_a = \int d\sigma \frac{\partial X_a}{\partial \sigma}.$$  

(12)

If $Z$ should be different from zero we need our space to be homologically non-trivial, for example $S^1 \times \mathbb{R}^{d-1}$. For a closed string the coordinate $x$ would then be wrapped \[13\] around $S^1$, which implies $x=2\pi Rk\sigma$, $\sigma \in [0,1]$, where $k$ is the winding number and $R$ is the radius of $S^1$. Similarly, if we have a homologically trivial transverse space $N_a$ is zero, but if we consider some directions of the transverse space as tori, say, we will have some of the $N_a \neq 0$.

3 Non-relativistic d-branes

World-volume (longitudinal) dimensions are labeled by $x^\alpha$ with $\alpha = 0, 1, \cdots, d$ and Lorentzian metric $\eta_{\alpha\beta} = diag(-1, +1, \cdots, +1)$, while the remaining (transverse) dimensions are labeled by $X^a$ with $a = 1, \cdots, D$ and Euclidean metric $\delta_{ab}$.

The group elements in the coset are

$$g = \exp(i x^\alpha p_\alpha) \exp(i X^a P_a + iv^{\alpha a} K_{\alpha a})$$

(13)

As before, the algebra, without central extensions\footnote{We will include central and topological charges \[14\] in the algebra later. See \[3\] below, and \[11\].}, is easily abstracted from the contracted \((c \rightarrow \infty)\) Galilean correspondence

$$p_\alpha \sim -i \frac{\partial}{\partial x^\alpha}, \quad P_a \sim -i \frac{\partial}{\partial X^a}, \quad K_{\alpha a} \sim -ix_\alpha \frac{\partial}{\partial X^a}$$

$$m_{\alpha\beta} \sim -i \left( x_\alpha \frac{\partial}{\partial x^\beta} - x_\beta \frac{\partial}{\partial x^\alpha} \right), \quad M_{ab} \sim -i \left( X_a \frac{\partial}{\partial X^b} - X_b \frac{\partial}{\partial X^a} \right)$$

(14)
Thus the relevant non-vanishing commutators are

\[ [K_{\alpha a}, p_{\beta}] = i\eta_{\alpha\beta} P_a \]  

(15)
as well as \[ [K_{\alpha a}, m_{\beta\gamma}] = i\eta_{\alpha\gamma} K_{\beta a} - i\eta_{\alpha\beta} K_{\gamma a} \], \[ [K_{\alpha a}, M_{bc}] = i\delta_{ac} K_{ab} - i\delta_{ab} K_{ac} \], etc. This\nleads to the one-form\n
\[ \Omega = -ig^{-1} dg \]

\[ = p_{\alpha} \, dx^\alpha + P_a \, (dX^a + v_{aa} \, dx^\alpha) + K_{aa} \, dv^{aa} \]

\[ \equiv p_{\alpha} \omega^\alpha + P_a \omega^a + K_{aa} \omega^{aa} \]  

(16)
with component one-forms defined as²

\[ \omega^\alpha = dx^\alpha, \quad \omega^a = dX^a + v_{aa} \, dx^\alpha, \quad \omega^{aa} = dv^{aa} \]  

(17)
The Maurer-Cartan equations involving the non-vanishing structure constants, \( f_{K_{\alpha a}, p_{\beta}, p_b} = \eta_{\alpha\beta} \delta_{ab} \), are given here by \( dv^\alpha = \omega^{aa} \wedge \omega_\alpha \).

We now construct from the component one-forms a closed, invariant \( d + 2 \) form

\[ \Omega_{d+2} = \frac{(-1)^d}{d!} \varepsilon_{\alpha_1 \ldots \alpha_d} \omega^{\alpha_1} \wedge \ldots \wedge \omega^{\alpha_d} \wedge \omega^{\beta b} \wedge \omega^b = d\Phi_{d+1} \]  

(18)
with \( \varepsilon_{01\ldots d} \equiv +1 = -\varepsilon^{01\ldots d} \), to produce a Wess-Zumino brane action

\[ A = T_d \int_{M_{d+2}} * \Omega_{d+2} = T_d \int_{M_{d+1}} = \partial_{M_{d+2}} * \Phi_{d+1}, \]  

(19)
where \( T_d \) is the tension of the d-brane. The boundary \( M_{d+1} = \partial M_{d+2} \) is the world-volume of the evolving brane. Straightforward calculation shows that an appropriate choice for the brane form potential is

\[ \Phi_{d+1} = \frac{1}{d!} \varepsilon_{\alpha_1 \ldots \alpha_d} \omega^{\alpha_1} \wedge \ldots \wedge dx^{\alpha_d} \wedge \left( v^{\beta b} dX^b + \frac{1}{2d+2} v_{\beta}^b v^{\gamma b} dx^\beta \right) \]  

(20)
Recall the \( a, b, \) etc. indices are contracted with Euclidean metric, while \( \alpha, \beta, \) etc. are handled with Lorentz metric.

As a special case, reconsider the NR string, with \( d = 1 \). We have

\[ \Omega_3 = \varepsilon_{\alpha\beta} \omega^\alpha \wedge \omega^b \wedge \omega^{\beta b} = \varepsilon_{\alpha\beta} \omega^\alpha \wedge (dX^b + v_{\gamma}^b dx^\gamma) \wedge dv^{\beta b} = d\Phi_2 \]

\[ \Phi_2 = \varepsilon_{\alpha\beta} dx^\alpha \wedge \left( v^{\beta b} dX^b + \frac{1}{4} v_{\beta}^b v^{\gamma b} dx^\beta \right) \]  

(21)
in agreement with (7) and (8). We note in passing that \( \Omega_3 \) is an apparent modification of the usual torsion 3-form for non-linearly realized (compact) semi-simple Lie groups \[15\] when expressed in terms of the component forms (i.e. vielbeins)³.

²The string case is obtained by writing \( \omega^b = \phi^b \) and \( \theta^b = \phi^b \).

³Naively, this other 3-form, \( \Omega_3 \), would be built from the component forms by using the structure constants \( f_{K_{\alpha a}, p_{\beta}, p_b} \) given earlier, and differs from \( \Omega_3 \) by replacement of \( \varepsilon_{\alpha\beta} \) with \( \eta_{\alpha\beta} \). For an arbitrary d-brane, it would be

\[ \tilde{\Omega}_3 = \eta_{\alpha\beta} \delta_{ab} \omega^\alpha \wedge \omega^a \wedge \omega^{\beta b} = dx_a \wedge (dX^b + v_{\gamma}^b dx^\gamma) \wedge dv^{\alpha b} \]

Remarkably, this 3-form is not closed, for arbitrary \( d \), but rather gives \( d\tilde{\Omega}_3 = (dx_a \wedge dv^{\alpha b}) \wedge (dx_\beta \wedge dv^{\beta b}) \).
If we fix the static gauge $x^\alpha = \sigma^\alpha$, the pullback to the d-brane world-volume goes as follows.

$$
\Phi_{d+1}|_{X(x)} = \frac{1}{d!} \varepsilon_{\alpha_1 \cdots \alpha_d} \left( -\varepsilon^{\alpha_1 \cdots \alpha_d \gamma} \gamma_{\beta} b \frac{\partial X^b}{\partial x^\gamma} - \frac{1}{2d + 2} v_{\gamma}^b v_{\gamma}^b \varepsilon^{\alpha_1 \cdots \alpha_d \beta} \right) dx^0 \wedge dx^1 \wedge \cdots \wedge dx^d
$$

That is to say

$$
\Phi_{d+1}|_{X(x)} = \left( v_{\gamma}^b \frac{\partial X^b}{\partial x^\gamma} + \frac{1}{2} v_{\gamma}^b v_{\gamma}^b \right) dx^0 \wedge dx^1 \wedge \cdots \wedge dx^d
$$

So the action is

$$
\mathcal{A} = T_d \int_{M_{d+1}} \left( v_{\gamma}^b \frac{\partial X^b}{\partial x^\gamma} + \frac{1}{2} v_{\gamma}^b v_{\gamma}^b \right) dx^0 \wedge dx^1 \wedge \cdots \wedge dx^d
$$

Eliminating the auxiliary Goldstone fields $v_{\gamma}^b$ using their equations of motion, $v_{\gamma}^b = -\frac{\partial X^b}{\partial x^\gamma}$, gives the gauge fixed action

$$
\mathcal{A} = T_d \int_{M_{d+1}} \left( -\frac{1}{2} \eta^{\alpha \beta} \frac{\partial X^a}{\partial x^\alpha} \frac{\partial X^b}{\partial x^\beta} \right) dx^0 \wedge dx^1 \wedge \cdots \wedge dx^d
$$

So, in a straightforward generalization of the string case given previously, the $X^a$ considered as functions of the $x^\alpha$ are free massless fields on the world-volume.

In terms of another world-volume parameterization, $z^\alpha, \alpha = 0, 1, \cdots, d$, and $z^\beta = \eta^{\beta \alpha} z^\alpha$, we have $dx^0 \wedge dx^1 \wedge \cdots \wedge dx^d = |M| \ dz^0 \wedge \cdots \wedge dz^d$ where

$$
M^\alpha_{\beta} \equiv \frac{\partial}{\partial z^\beta} x^\alpha, \quad M_{\beta \alpha} = \frac{\partial}{\partial x_\alpha} x^\beta, \quad \frac{\partial X^b}{\partial x^\gamma} = (M^{-1})^{\gamma \beta}_{\gamma \beta} \frac{\partial X^b}{\partial z^\beta}
$$

$$
(M^{-1})_{\alpha \beta} = \frac{1}{|M|} \mathcal{C}(M)_{\beta \alpha}, \quad |M| = \det M^\alpha_{\beta} = \varepsilon_{\alpha_\alpha_1 \cdots \alpha_d} \frac{\partial x^{\alpha_1}}{\partial z_{\beta_1}} \cdots \frac{\partial x^{\alpha_d}}{\partial z_{\beta_d}}
$$

with $\mathcal{C}(M)$ the usual matrix of cofactors, $\mathcal{C}(M)_{\beta \alpha} = (-1)^{\alpha + \beta} \det \left( M_{\text{remove } \beta \text{th row and } \alpha \text{th column}} \right)$. Explicitly

$$
\mathcal{C}(M)_{\beta \alpha} = -\frac{1}{d!} \varepsilon_{\beta_1 \cdots \beta_d} \varepsilon_{\alpha_1 \cdots \alpha_d} \frac{\partial x^{\alpha_1}}{\partial z_{\beta_1}} \cdots \frac{\partial x^{\alpha_d}}{\partial z_{\beta_d}}
$$

In this arbitrary parameterization

$$
\Phi_{d+1} = \left( v_{\gamma}^b (M^{-1})_{\gamma \beta} \frac{\partial X^b}{\partial z^\beta} + \frac{1}{2} v_{\gamma}^b v_{\gamma}^b \right) |M| \ dz^0 \wedge \cdots \wedge dz^d
$$

The action in an arbitrary gauge is therefore

$$
\mathcal{A} = T_d \int_{M_{d+1}} \left( v_{\gamma}^b \mathcal{C}(M)_{\beta \gamma} \frac{\partial X^b}{\partial z^\beta} + \frac{1}{2} v_{\gamma}^b v_{\gamma}^b \right) |M| \ dz^0 \wedge \cdots \wedge dz^d
$$

Eliminating once again the Goldstone auxiliaries using their equations of motion, $v_{\gamma}^b = -\left( M^{-1} \right)_{\gamma \alpha} \frac{\partial X^b}{\partial z^\alpha}$, gives (for $d = 2$, see [12])

$$
\mathcal{A} = T_d \int_{M_{d+1}} \left( -\frac{1}{2} |M| \ \mathcal{C}(M)_{\alpha \gamma} \eta^{\alpha \delta} \mathcal{C}(M)_{\beta \delta} \frac{\partial X^b}{\partial z^\alpha} \frac{\partial X^b}{\partial z^\beta} \right) dz^0 \wedge \cdots \wedge dz^d
$$
The quadratic in cofactors can be written more explicitly using \(27\) above.

\[
\mathcal{C}(M)_{\alpha\gamma} \eta^{\delta} \mathcal{C}(M)_{\beta\delta} = - \frac{1}{d!} \varepsilon_{\alpha_1 \cdots \alpha_d} \varepsilon_{\beta_1 \cdots \beta_d} \frac{\partial x_{\gamma_1}}{\partial z_{\alpha_1}} \frac{\partial x_{\gamma_2}}{\partial z_{\alpha_2}} \cdots \frac{\partial x_{\gamma_d}}{\partial z_{\alpha_d}} \frac{\partial x_{\gamma_d}}{\partial z_{\beta_d}}
\]

(31)

Rather deceptively, \(30\) looks like the action for a nonlinear, interacting theory coupling the transverse dimensions to the world-volume variables. Of course, in the gauge \(x^\alpha = z^\alpha\), we have \(M_{\alpha\beta} = \eta_{\alpha\beta}\), and we recover a free-field action as in the string case—a simplification typical of re-parameterization invariance.

The Noether charges for the d-brane, in the gauge fixed form \(25\), are given at any world-volume time as

\[
P_a = \int d^d x \; \Pi_a \; , \; \quad K_{\alpha a} = \int d^d x \; (x_\alpha \Pi_a - \eta_{\alpha 0} X_a)
\]

\[
p_0 = - \int d^d x \; (\frac{1}{2} \Pi_a \Pi_a + \frac{1}{2} \partial_i X_a \partial_i X_a) \equiv \int d^d x \; \mathcal{P}_0 \; , \; \quad p_i = - \int d^d x \; \Pi_a \partial_i X_a \equiv \int d^d x \; \mathcal{P}_i
\]

\[
m_{\alpha\beta} = \int d^d x \; (x_\alpha \mathcal{P}_\beta - x_\beta \mathcal{P}_\alpha) \; , \; \quad M_{ab} = \int d^d x \; (X_a \Pi_b - X_b \Pi_a)
\]

(32)

where \(a, b = 1, \cdots , D\), and \(\alpha, \beta = 0, 1, \cdots , d\), but \(i = 1, \cdots , d\). Here, \(\Pi_a\) is the canonically conjugate momentum associated with \(X_a\), so \([X_a(x), \Pi_b(x')]_{x=a=x'} = i \delta_{ab} \delta^{(d)}(x-x')\).

For the general d-brane, the commutation relations are just a straightforward generalization of the non-relativistic string case. However, a greater variety of central and topological charges may appear in the d-brane algebra, for example as

\[
[P_a, K_{ab}] = i \eta_{0a} \delta_{ab} Z \; , \; \quad [K_{0a}, K_{ib}] = i \delta_{ab} W_i \; , \; \quad [p_i, K_{0a}] = i N_{ia}.
\]

(33)

These central and topological charges are given formally for the d-brane by\(^4\)

\[
Z = \int d^d x \; , \; \quad W_i = \int x_i \; d^d x \; , \; \quad N_{ia} = \int \frac{\partial X_a}{\partial x^i} \; d^d x.
\]

(34)

To avoid ambiguities in these formal expressions, we have checked the results against various Jacobi identities. A full discussion of all Jacobi identities is deferred \[11\].

4 Conclusion

We have shown how to construct non-relativistic string and brane actions from the structure of the underlying Galilei group. We will investigate quantum properties of these models and discuss the roles of their topological charges in a subsequent paper \[11\].

In certain cases, perhaps all, passive advections of non-relativistic strings and membranes also result from a Wess-Zumino action when expressed in the formalism of Nambu mechanics \[10\]. It would be interesting to combine fully that other formalism with the group theoretic approach of the present paper.

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\(^4\)In the last relation in \(33\) and/or \(34\), we have assumed there are no conjugate momentum winding numbers, so we have set \(\int d^d x \; \partial_i \Pi_a(x) = 0\). If otherwise, we would write \(N_{ia} = \int (\partial_i X_a + x_0 \partial_i \Pi_a) \; d^d x\).
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