

BMS algebra and supertranslation-invariant Lorentz charges

Marc Henneaux
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The BMS_4 group

**BMS algebra and
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The group of asymptotic symmetries of gravity in the
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is infinite-dimensional

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and was discovered first by Bondi, Metzner and Sachs,

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hence the name “BMS” group

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The group of asymptotic symmetries of gravity in the
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is infinite-dimensional

and was discovered first by Bondi, Metzner and Sachs,

hence the name “BMS” group

or “ BMS_4 ” group if one wants to emphasize the spacetime
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In addition to the homogeneous Lorentz group,

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In addition to the homogeneous Lorentz group,
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The structure of this group was understood by Sachs.
In addition to the homogeneous Lorentz group,
there are the familiar translations
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which commute among themselves
and transform in an infinite-dimensional non-trivial
representation of the Lorentz group.

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The BMS_4 algebra is the semi-direct sum of the Lorentz algebra \mathcal{L} and the infinite-dimensional algebra \mathcal{A} spanned by the supertranslations,

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$$BMS_4 = \mathcal{L} \oplus_{\sigma} \mathcal{A}$$

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The supertranslations are parametrized by functions on the sphere,

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The $\ell = 0$ and $\ell = 1$ harmonics correspond to the ordinary translations ;

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The $\ell = 0$ and $\ell = 1$ harmonics correspond to the ordinary translations ;

the higher harmonics are the pure supertranslations.

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Schematically, if we denote the generators of the homogeneous Lorentz group by M_a ,

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Schematically, if we denote the generators of the homogeneous Lorentz group by M_a ,

one finds

$$\{M_a, M_b\} = f_{ab}^c M_c,$$

$$\{M_a, T_i\} = R_{ai}^j T_j,$$

$$\{M_a, S_\alpha\} = G_{a\alpha}^i T_i + G_{a\alpha}^{\beta} S_\beta,$$

$$\{T_i, T_j\} = 0 = \{T_i, S_\alpha\} = \{S_\alpha, S_\beta\}$$

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$$\{T_i, T_j\} = 0 = \{T_i, S_\alpha\} = \{S_\alpha, S_\beta\}$$

where T_i and S_α are respectively the generators of the standard translations and of the pure supertranslations.

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The structure constants R_{ai}^j , $G_{a\alpha}^i$ and $G_{a\alpha}^\beta$ are non zero.

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The structure constants R_{ai}^j , G_{aa}^i and G_{aa}^β are non zero.

The ordinary translations transform in the 4-dimensional vector representation of the Lorentz group (R_{ai}^j).

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The structure constants R_{ai}^j , G_{aa}^i and G_{aa}^β are non zero.

The ordinary translations transform in the 4-dimensional vector representation of the Lorentz group (R_{ai}^j).

Modulo the ordinary translations, the pure supertranslations transform in an infinite-dimensional representation of the Lorentz group (G_{aa}^β).

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There are intriguing physical features resulting from the structure of the BMS₄ algebra

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There are intriguing physical features resulting from the structure of the BMS₄ algebra

which are somewhat uncomfortable.

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Modulo the ordinary translations, the pure supertranslations transform in an infinite-dimensional representation of the Lorentz group ($G_{a\alpha}^\beta$).

There are intriguing physical features resulting from the structure of the BMS₄ algebra

which are somewhat uncomfortable.

These features were first pointed out at null infinity but are equally present at spatial infinity.

No invariant Poincaré subalgebra

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The fact that the Poisson brackets of the homogeneous Lorentz generators with the pure supertranslations involve both the pure supertranslations and the translations implies that the Poincaré subalgebra is not an ideal.

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The fact that the Poisson brackets of the homogeneous Lorentz generators with the pure supertranslations involve both the pure supertranslations and the translations implies that the Poincaré subalgebra is not an ideal.

At the same time, because the pure supertranslations do not form an ideal on account of the preceding point, they cannot be meaningfully quotientized out to get the Poincaré algebra as a quotient algebra.

Angular momentum ambiguity

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It follows from the non-vanishing of the bracket of the pure supertranslations with the homogeneous Lorentz transformations that the angular momentum transforms under pure supertranslations.

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It follows from the non-vanishing of the bracket of the pure supertranslations with the homogeneous Lorentz transformations that the angular momentum transforms under pure supertranslations.

This non-invariance comes on top of the familiar non-invariance of the angular momentum under ordinary translations, but there one knows how to define an intrinsic angular momentum in terms of Casimirs of the Poincaré algebra, which amounts to compute the angular momentum with respect to the center of mass worldline.

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A similar construction for supertranslations appears to be more intricate for the full BMS algebra.

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The purpose of this talk is to show that these difficulties disappear if one enlarges the BMS symmetry

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The purpose of this talk is to show that these difficulties disappear if one enlarges the BMS symmetry

by including logarithmic supertranslations, which are algebraically conjugate to the pure supertranslations.

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Logarithmic supertranslations have appeared here and then in the literature on asymptotic symmetries but have not been systematically studied.

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In particular, the algebraic structure of the logarithmic BMS algebra and the fact that one could rewrite it as a direct sum by nonlinear redefinitions were not discussed.

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Work done in collaboration with Oscar Fuentealba and Cédric Troessaert, e-Prints 2211.10941 [hep-th] and 2305.05436 [hep-th]

ADM action

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How does the BMS algebra appear at spatial infinity?

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(Dirac, Arnowitt-Deser-Misner, Regge-Teitelboim)

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We shall insist that the boundary conditions make the action :

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- and invariant under all (asymptotic) Poincaré symmetries, which are thus canonical transformations.

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This puts strong and interesting restrictions.

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The usually assumed fall-off is (in cartesian coordinates)

$$g_{ij} = \delta_{ij} + O(r^{-1}), \quad \pi^{ij} = O(r^{-2}).$$

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Without additional requirement, this fall-off is too slow because it generically leads to a logarithmic divergence in the symplectic structure (kinetic term)

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$$\int d^3x \pi^{ij} \dot{g}_{ij} \sim \ln r.$$

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One way to cure this problem would be to impose that the leading terms of the metric and its conjugate momentum have opposite parity properties under the antipodal map,

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$$h_{ij} \equiv g_{ij} - \delta_{ij} = \frac{\bar{h}_{ij}(\mathbf{n}^k)}{r} + O\left(\frac{1}{r^2}\right), \quad \bar{h}_{ij}(-\mathbf{n}^k) = \bar{h}_{ij}(\mathbf{n}^k)$$

and

$$\pi^{ij} = \frac{\bar{\pi}^{ij}(\mathbf{n}^k)}{r^2} + O\left(\frac{1}{r^3}\right), \quad \bar{\pi}^{ij}(-\mathbf{n}^k) = -\bar{\pi}^{ij}(\mathbf{n}^k).$$

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But these strict parity conditions leave no room for the BMS group.

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But these strict parity conditions leave no room for the BMS group.

The asymptotic symmetry reduces then to the Poincaré group (and not more).

Parity-twisted boundary conditions

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The strict parity conditions are too strong and kill the pure supertranslations.

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One must relax these parity conditions... but not completely if one wants to maintain finiteness of the action.

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The idea is to allow a “parity-twisted component” of a specific form in the leading orders in the asymptotic expansions of the metric and its conjugate momentum.

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The idea is to allow a “parity-twisted component” of a specific form in the leading orders in the asymptotic expansions of the metric and its conjugate momentum.

This parity-twisted component takes the form of a diffeomorphism generated by $\mathcal{O}(1)$ diffeomorphisms (rewritten in Hamiltonian form).

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Specifically, one takes for the metric

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Specifically, one takes for the metric

$$\begin{aligned}h_{ij} &\equiv g_{ij} - \delta_{ij} = U_{ij} + j_{ij}, \\U_{ij} &= \partial_i \zeta_j + \partial_j \zeta_i = O\left(\frac{1}{r}\right), \quad \zeta^i = \zeta^i(\mathbf{n}^k) = O(1), \quad \zeta^i(-\mathbf{n}^k) = \zeta^i(\mathbf{n}^k), \\j_{ij} &= \frac{(\bar{h}_{ij})^{even}(\mathbf{n}^k)}{r} + O\left(\frac{1}{r^2}\right), \quad (\bar{h}_{ij})^{even}(-\mathbf{n}^k) = (\bar{h}_{ij})^{even}(\mathbf{n}^k)\end{aligned}$$

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U_{ij} is the contribution that twists the strict parity condition on the metric by an $\mathcal{O}(1)$ -diffeomorphism (to leading order) and can be assumed to be odd (ζ_i even).

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It turns out that in order for the boosts to preserve the symplectic form and hence define canonical transformations,

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ζ_i must takes the specific form

$$\zeta_i = \partial_i(rU) \quad \bar{U}(\mathbf{n}^k) = O(1) = -U(-\mathbf{n}^k).$$

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$$\zeta_i = \partial_i(rU) \quad \bar{U}(\mathbf{n}^k) = O(1) = -U(-\mathbf{n}^k).$$

and so

$$U_{ij} = 2\partial_i\partial_j(rU).$$

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For the conjugate momenta, one takes

$$\pi^{ij} = V^{ij} + p^{ij}$$

$$V^{ij} = \partial^i \partial^j V - \delta^{ij} \Delta V, \quad V = V(\mathbf{n}^k) = O(1), \quad V(-\mathbf{n}^k) = V(\mathbf{n}^k)$$

$$p^{ij} = \frac{(\bar{\pi}^{ij})^{odd}(\mathbf{n}^k)}{r^2} + O\left(\frac{1}{r^3}\right), \quad (\bar{\pi}^{ij})^{odd}(-\mathbf{n}^k) = -(\bar{\pi}^{ij})^{odd}(\mathbf{n}^k)$$

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$$\pi^{ij} = V^{ij} + p^{ij}$$

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where V is a function of the angles which can be assumed to be even.

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For the conjugate momenta, one takes

$$\begin{aligned}\pi^{ij} &= V^{ij} + p^{ij} \\ V^{ij} &= \partial^i \partial^j V - \delta^{ij} \Delta V, \quad V = V(\mathbf{n}^k) = O(1), \quad V(-\mathbf{n}^k) = V(\mathbf{n}^k) \\ p^{ij} &= \frac{(\bar{\pi}^{ij})^{odd}(\mathbf{n}^k)}{r^2} + O\left(\frac{1}{r^3}\right), \quad (\bar{\pi}^{ij})^{odd}(-\mathbf{n}^k) = -(\bar{\pi}^{ij})^{odd}(\mathbf{n}^k)\end{aligned}$$

where V is a function of the angles which can be assumed to be even.

V^{ij} is the contribution that twists the strict parity of π^{ij} (at leading order) by a diffeomorphism (written in Hamiltonian form).

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For the conjugate momenta, one takes

$$\pi^{ij} = V^{ij} + p^{ij}$$

$$V^{ij} = \delta^i \partial^j V - \delta^{ij} \Delta V, \quad V = V(\mathbf{n}^k) = O(1), \quad V(-\mathbf{n}^k) = V(\mathbf{n}^k)$$

$$p^{ij} = \frac{(\bar{\pi}^{ij})^{odd}(\mathbf{n}^k)}{r^2} + O\left(\frac{1}{r^3}\right), \quad (\bar{\pi}^{ij})^{odd}(-\mathbf{n}^k) = -(\bar{\pi}^{ij})^{odd}(\mathbf{n}^k)$$

where V is a function of the angles which can be assumed to be even.

V^{ij} is the contribution that twists the strict parity of π^{ij} (at leading order) by a diffeomorphism (written in Hamiltonian form).

The twisting is thus characterized by two functions of the angles, one odd (U) and one even (V).

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These relaxed parity conditions involving a twist still lead to a consistent dynamical description
(finite action, finite symplectic form, well-defined generators).

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These relaxed parity conditions involving a twist still lead to a consistent dynamical description

(finite action, finite symplectic form, well-defined generators).

Furthermore, the group of asymptotic symmetries is the full BMS group (written in a different basis than the familiar null infinity basis).

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These relaxed parity conditions involving a twist still lead to a consistent dynamical description

(finite action, finite symplectic form, well-defined generators).

Furthermore, the group of asymptotic symmetries is the full BMS group (written in a different basis than the familiar null infinity basis).

There is complete agreement with the null infinity results.

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There is complete agreement with the null infinity results.

Can one relax even more the boundary conditions in a useful way?

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(finite action, finite symplectic form, well-defined generators).

Furthermore, the group of asymptotic symmetries is the full BMS group (written in a different basis than the familiar null infinity basis).

There is complete agreement with the null infinity results.

Can one relax even more the boundary conditions in a useful way?

The answer is positive. One can further enlarge the symmetry in this manner, which then includes in addition “logarithmic supertranslations” conjugate to supertranslations.

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To exhibit the logarithmic supertranslations, one must allow terms in the canonical variables that result from diffeomorphisms that grow like $\ln r$ at infinity.

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To exhibit the logarithmic supertranslations, one must allow terms in the canonical variables that result from diffeomorphisms that grow like $\ln r$ at infinity.

The logarithmic supertranslations are diffeomorphisms that grow like $\ln r$ at infinity and which preserve by construction the new boundary conditions,

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The logarithmic supertranslations are diffeomorphisms that grow like $\ln r$ at infinity and which preserve by construction the new boundary conditions,

They define true symmetries of the action provided some conditions on the logarithmic diffeomorphisms are imposed.

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The logarithmic supertranslations are diffeomorphisms that grow like $\ln r$ at infinity and which preserve by construction the new boundary conditions,

They define true symmetries of the action provided some conditions on the logarithmic diffeomorphisms are imposed.

The (allowed) logarithmic supertranslations are then canonical transformations with a well-defined, finite charge.

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To exhibit the logarithmic supertranslations, one must allow terms in the canonical variables that result from diffeomorphisms that grow like $\ln r$ at infinity.

The logarithmic supertranslations are diffeomorphisms that grow like $\ln r$ at infinity and which preserve by construction the new boundary conditions,

They define true symmetries of the action provided some conditions on the logarithmic diffeomorphisms are imposed.

The (allowed) logarithmic supertranslations are then canonical transformations with a well-defined, finite charge.

They depend on a single function of the angles, but the $\ell = 0$ and $\ell = 1$ harmonics define proper gauge transformations : they have zero charge.

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To exhibit the logarithmic supertranslations, one must allow terms in the canonical variables that result from diffeomorphisms that grow like $\ln r$ at infinity.

The logarithmic supertranslations are diffeomorphisms that grow like $\ln r$ at infinity and which preserve by construction the new boundary conditions,

They define true symmetries of the action provided some conditions on the logarithmic diffeomorphisms are imposed.

The (allowed) logarithmic supertranslations are then canonical transformations with a well-defined, finite charge.

They depend on a single function of the angles, but the $\ell = 0$ and $\ell = 1$ harmonics define proper gauge transformations : they have zero charge.

The physical logarithmic supertranslations are thus parametrized by the same class of functions as the pure supertranslations. We denote them L_α .

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One can compute the Poisson brackets of the logarithmic supertranslation generators with themselves and with the generators of the BMS_4 algebra.

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One can compute the Poisson brackets of the logarithmic supertranslation generators with themselves and with the generators of the BMS_4 algebra.

The computation is direct and follows standard canonical methods.

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One can compute the Poisson brackets of the logarithmic supertranslation generators with themselves and with the generators of the BMS_4 algebra.

The computation is direct and follows standard canonical methods.

One finds that the generators L^α of the logarithmic supertranslations commute among themselves,

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One can compute the Poisson brackets of the logarithmic supertranslation generators with themselves and with the generators of the BMS_4 algebra.

The computation is direct and follows standard canonical methods.

One finds that the generators L^α of the logarithmic supertranslations commute among themselves,

$$\{L^\alpha, L^\beta\} = 0,$$

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One can compute the Poisson brackets of the logarithmic supertranslation generators with themselves and with the generators of the BMS_4 algebra.

The computation is direct and follows standard canonical methods.

One finds that the generators L^α of the logarithmic supertranslations commute among themselves,

$$\{L^\alpha, L^\beta\} = 0,$$

and transform in the same representation of the Lorentz group as the pure supertranslations,

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One can compute the Poisson brackets of the logarithmic supertranslation generators with themselves and with the generators of the BMS_4 algebra.

The computation is direct and follows standard canonical methods.

One finds that the generators L^α of the logarithmic supertranslations commute among themselves,

$$\{L^\alpha, L^\beta\} = 0,$$

and transform in the same representation of the Lorentz group as the pure supertranslations,

$$\{M_a, L^\alpha\} = -G_{a\beta}{}^\alpha L^\beta$$

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Putting everything together, one thus gets as non-zero brackets

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Putting everything together, one thus gets as non-zero brackets

$$\{M_a, M_b\} = f_{ab}^c M_c,$$

$$\{M_a, T_i\} = R_{ai}^j T_j,$$

$$\{M_a, S_\alpha\} = G_{a\alpha}^i T_i + G_{a\alpha}^\beta S_\beta,$$

$$\{M_a, L^\alpha\} = -G_{a\beta}^\alpha L^\beta,$$

$$\{L^\alpha, S_\beta\} = \delta_\beta^\alpha,$$

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$$\{M_a, S_\alpha\} = G_{a\alpha}^i T_i + G_{a\alpha}^\beta S_\beta,$$

$$\{M_a, L^\alpha\} = -G_{a\beta}^\alpha L^\beta,$$

$$\{L^\alpha, S_\beta\} = \delta_\beta^\alpha,$$

i.e., $\text{LogBMS}_4 = \mathcal{L} \oplus_\sigma (\mathcal{A} \oplus_c \mathcal{B})$

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$$\{M_a, M_b\} = f_{ab}^c M_c,$$

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$$\{M_a, S_\alpha\} = G_{a\alpha}^i T_i + G_{a\alpha}^\beta S_\beta,$$

$$\{M_a, L^\alpha\} = -G_{a\beta}^\alpha L^\beta,$$

$$\{L^\alpha, S_\beta\} = \delta_\beta^\alpha,$$

i.e., $\text{LogBMS}_4 = \mathcal{L} \oplus_\sigma (\mathcal{A} \oplus_c \mathcal{B})$

with \mathcal{B} being the abelian algebra of the logarithmic
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The bracket of two conserved charges is a conserved charge.

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The bracket of two conserved charges is a conserved charge.

A set $\{Q_A(q, p)\}$ of conserved charges is a complete set (of conserved charges)

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The bracket of two conserved charges is a conserved charge.

A set $\{Q_A(q, p)\}$ of conserved charges is a complete set (of conserved charges)

if any conserved charge is a function $f(Q_A)$ of these conserved charges.

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A set $\{Q_A(q, p)\}$ of conserved charges is a complete set (of conserved charges)

if any conserved charge is a function $f(Q_A)$ of these conserved charges.

The function f need not be linear.

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The bracket of two conserved charges is a conserved charge.

A set $\{Q_A(q, p)\}$ of conserved charges is a complete set (of conserved charges)

if any conserved charge is a function $f(Q_A)$ of these conserved charges.

The function f need not be linear.

[If p is a constant of the motion, p^2 or any function of p is also a constant of the motion but it is clearly not independent from p even though not a multiple of p .]

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If we denote by X_A the Hamiltonian vector field corresponding to Q_A , i.e. $X_A F = \{F, Q_A\}$,

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If we denote by X_A the Hamiltonian vector field corresponding to Q_A , i.e. $X_A F = \{F, Q_A\}$,

then the Hamiltonian vector field X associated with $f(Q)$ is given by

$$X = \frac{\partial f}{\partial Q_A} X_A$$

since $X F = \{F, f(Q)\} = \frac{\partial f}{\partial Q_A} \{F, Q_A\}$.

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If we denote by X_A the Hamiltonian vector field corresponding to Q_A , i.e. $X_A F = \{F, Q_A\}$,

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$$X = \frac{\partial f}{\partial Q^A} X_A$$

since $XF = \{F, f(Q)\} = \frac{\partial f}{\partial Q^A} \{F, Q_A\}$.

It is a linear combination $f^A X_A$ of the vector fields X_A of the complete set,

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These phase space functions are not arbitrary, however, but must depend on the phase space variables through the Q_A 's and in such a way that X is a Hamiltonian vector field, i.e., $\frac{\partial f^A}{\partial Q^B} = \frac{\partial f^B}{\partial Q^A}$.

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($f^A dQ_A$ is exact, “integrability”)

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In general,

$$[X_A, X_B] = -\frac{\partial f_{AB}}{\partial Q_C} X_C.$$

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In general,

$$[X_A, X_B] = -\frac{\partial f_{AB}}{\partial Q_C} X_C.$$

Complete set of conserved charges can be non-linearly redefined,

$$Q_A \rightarrow \bar{Q}_A = \bar{Q}_A(Q), \quad \det\left(\frac{\partial \bar{Q}_A}{\partial Q_B}\right) \neq 0.$$

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The Q_A 's define a Poisson manifold.

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There exist theorems that generalize Darboux theorem.

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There exist theorems that generalize Darboux theorem.

[If f_{AB} in $\{Q_A, Q_B\} = f_{AB}(Q)$ is invertible, one can go to Darboux coordinates where $f_{AB} = \delta_{AB}$.]

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The above well-known properties of Hamiltonian systems,
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There are many examples where the Poisson bracket of two
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Even when the Poisson brackets are linear, one can consider
non-linear redefinitions of the charges.

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$$\{M_a, M_b\} = f_{ab}^c M_c,$$

$$\{M_a, T_i\} = R_{ai}^j T_j,$$

$$\{M_a, S_\alpha\} = G_{a\alpha}^i T_i + G_{a\alpha}^\beta S_\beta,$$

$$\{M_a, L^\alpha\} = -G_{a\beta}^\alpha L^\beta,$$

$$\{L^\alpha, S_\beta\} = \delta_\beta^\alpha,$$

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the fact that the central charge in $\{L^\alpha, S_\beta\}$ is invertible enables one to rewrite the algebra as the direct sum $\mathcal{P} \oplus (\mathcal{A}' \oplus_c \mathcal{B})$, where \mathcal{A}' is the abelian algebra of pure supertranslations.

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The logarithmic supertranslations are canonically conjugate to the pure supertranslations, and a Darboux-like procedure enables one to “decouple” them from the other generators.

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The only generators that need to be redefined are actually the Lorentz generators, as follows,

$$\tilde{M}_a = M_a - G_{a\beta}{}^i L^\beta T_i - G_{a\beta}{}^\gamma L^\beta S_\gamma \quad (5.1)$$

$$= M_a - L^\beta \{M_a, S_\beta\}. \quad (5.2)$$

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$$= M_a - L^\beta \{M_a, S_\beta\}. \quad (5.2)$$

One easily verifies

$$\{\tilde{M}_a, S_\alpha\} = \{\tilde{M}_a, L^\alpha\} = 0, \quad (5.3)$$

while the bracket $\{\tilde{M}_a, T_i\}$ does not suffer any modification.

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while the bracket $\{\tilde{M}_a, T_i\}$ does not suffer any modification.

One can also check

$$\{\tilde{M}_a, \tilde{M}_b\} = f_{ab}{}^c \tilde{M}_c, \quad \{\tilde{M}_a, T_i\} = R_{ai}{}^j T_j. \quad (5.4)$$

so that the Poincaré algebra is unchanged.

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i.e., $\text{LogBMS}_4 = \mathcal{P} \oplus (\mathcal{A}' \oplus_c \mathcal{B})$

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The redefinitions of the Lorentz generators
involve quadratic terms of the forme TL and SL .

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The redefinitions of the Lorentz generators

involve quadratic terms of the forme TL and SL .

This means that the new Lorentz transformations will differ from the old ones by field dependent supertranslations and logarithmic supertranslations.

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The redefinitions of the Lorentz generators

involve quadratic terms of the forme *TL* and *SL*.

This means that the new Lorentz transformations will differ from the old ones by field dependent supertranslations and logarithmic supertranslations.

Integrability of the charges associated with these field-dependent transformations is not an issue since we are working with the generators throughout.

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By introducing logarithmic supertranslations,
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By introducing logarithmic supertranslations,
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one can redefine the Lorentz generators so that pure
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Setting the logarithmic supertranslations to zero is illegitimate
since they are improper gauge symmetries.

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This conclusion is in line with the analysis by Yau et al, Porrati et al, Compère et al at null infinity.

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The method can be also be applied to electromagnetism in such a way that the angle-dependent $u(1)$ gauge transformations commute with the Lorentz transformations, avoiding a “angle-dependent $u(1)$ ambiguity” in the definition of the angular momentum.

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THANK YOU!