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Letter

Remarks on Galilean electromagnetism

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ABSTRACT

Symmetries of the nonrelativistic electrodynamics are studied. It is shown that equations describing the Galilean electromagnetism in the presence of sources hold invariant under the ℓ–conformal Galilei group for an arbitrary (half)integer parameter ℓ. The group contains a transformation which links an inertial frame of reference to that moving with a constant acceleration of order up to $2\ell - 1$, thus pointing at potential dynamical instability of the system.

1. Introduction

Extensive recent studies of the nonrelativistic AdS/CFT–correspondence extended the holographic dictionary to cover strongly coupled condensed matter systems (see [1] and references therein). As far as symmetries are concerned, the key ingredient here is a finite–dimensional conformal extension of the Galilei group [2,3].

Generators of the corresponding Lie algebra include (temporal) translation, dilatation, and special conformal transformation, which form $sl(2, R) \sim so(2, 1)$ subalgebra, as well as spatial translations, Galilei boosts and higher order constant accelerations.[1] The associated structure relations involve an arbitrary (half)integer parameter ℓ, which specifies $2\ell - 1$ acceleration generators at hand and gives the name to the algebra [2,3]. In the literature, the reciprocal of ℓ is called the rational dynamical exponent. It is also customary to refer to the instances of $\ell = \frac{1}{2}$ and $\ell = 1$ as the Schrödinger algebra and the conformal Galilei algebra, respectively (for a recent review see [4]).

The study of dynamical realizations of the ℓ–conformal Galilei group revealed an interesting peculiarity. Because the number of functionally independent integrals of motion needed to integrate a differential equation correlates with its order, such systems in general involve higher derivative terms (see e.g. [5–10] and references therein[2]), the order of the highest derivative typically being $2\ell + 1$.

A particularly interesting example in the chain of models accommodating the ℓ–conformal Galilei symmetry is the Galilean electromagnetism [12]. As was demonstrated in [13], the corresponding equations of motion in the absence of sources hold invariant under the infinite dimensional extension of the $\ell = 1$ conformal Galilei group, which was called in [13] the infinite dimensional Galilean conformal group.

The goal of this paper is twofold. Firstly, it is shown that equations describing the Galilean electromagnetism in the presence of sources actually hold invariant under the ℓ–conformal Galilei group for an arbitrary (half)integer parameter ℓ. Secondly, it is emphasized that the system may potentially be plagued with dynamical instabilities, because the group contains a transformation which links an inertial frame of reference to that moving with a constant acceleration of order up to $2\ell - 1$ (see Section 5).

The work is organized as follows. In the next section, the electric and magnetic nonrelativistic limits of the Maxwell equations are outlined. In Section 3, the structure of the ℓ–conformal Galilei group is briefly reminded. In Section 4, it is demonstrated that equations describing the Galilean electromagnetism in the presence of sources hold invariant under the ℓ–conformal Galilei group for an arbitrary (half)integer parameter ℓ. Physical implications of such symmetry are discussed in the concluding Section 5.

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[1] The algebra also involves generators of spatial rotations, which in what follows will be disregarded.

[2] The only known example in the literature of a second order system, for which all constants of motion are functionally independent, was built in [11] by making recourse to geodesics on a Ricci–flat spacetime with the ℓ–conformal Galilei isometry group.

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Throughout the paper, summation over repeated indices is understood unless otherwise stated.

2. Nonrelativistic limits of the Maxwell equations

When exploring nonrelativistic limits of the Maxwell equations[3]

$$\partial_\mu F^{\nu\mu} = -\frac{4\pi}{c} j^\nu, \qquad \partial_\mu j^\mu = 0,$$

it is customary to rescale the scalar and vector potentials $\varphi \to \alpha(c)\varphi$, $\mathbf{A} \to \beta(c)\mathbf{A}$, as well as the charge density $\rho \to \gamma(c)\rho$ and the current density $\mathbf{j} \to \sigma(c)\mathbf{j}$, and then choose the factors $\alpha(c), \beta(c), \gamma(c), \sigma(c)$ in such a way that the resulting equations are nonsingular in the limit $c \to \infty$.

In general, more than one consistent nonrelativistic limit is possible and additional insight comes from the Lorentz boost transformation

$$\varphi = \frac{\varphi' + \frac{v}{c}A'^1}{\sqrt{1-\left(\frac{v}{c}\right)^2}}, \qquad A^1 = \frac{A'^1 + \frac{v}{c}\varphi'}{\sqrt{1-\left(\frac{v}{c}\right)^2}}, \qquad A^2 = A'^2, \qquad A^3 = A'^3,$$
$$j^0 = \frac{j'^0 + \frac{v}{c}j'^1}{\sqrt{1-\left(\frac{v}{c}\right)^2}}, \qquad j^1 = \frac{j'^1 + \frac{v}{c}j'^0}{\sqrt{1-\left(\frac{v}{c}\right)^2}}, \qquad j^2 = j'^2, \qquad j^3 = j'^3. \quad (1)$$

It suggests a particular rescaling $\mathbf{A} \to \frac{1}{c}\mathbf{A}$, with other fields unchanged, which guarantees that both the resulting transformation[4]

$$\varphi = \varphi', \qquad \mathbf{A} = \mathbf{A}' + \mathbf{v}\varphi', \qquad \rho = \rho', \qquad \mathbf{j} = \mathbf{j}' + \mathbf{v}\rho', \quad (2)$$

and the equations of motion

$$\mathrm{div}\mathbf{E} = 4\pi\rho, \quad \mathrm{rot}\mathbf{E} = 0, \quad \mathrm{rot}\mathbf{H} = \frac{\partial \mathbf{E}}{\partial t} + 4\pi\mathbf{j}, \quad \mathrm{div}\mathbf{H} = 0,$$
$$\frac{\partial \rho}{\partial t} + \mathrm{div}\mathbf{j} = 0, \quad (3)$$

where $\mathbf{E} = -\mathrm{grad}\varphi$, $\mathbf{H} = \mathrm{rot}\mathbf{A}$, are nonsingular after taking the limit $c \to \infty$. In the literature, this instance is known as the electric limit of the Maxwell equations [12].

Alternatively, one can choose $\mathbf{A} \to c\mathbf{A}$, $\mathbf{j} \to c^2\mathbf{j}$ (with φ and ρ unchanged), $c \to \infty$, which result in

$$\varphi = \varphi' + (\mathbf{v}, \mathbf{A}'), \qquad \mathbf{A} = \mathbf{A}', \qquad \rho = \rho' + (\mathbf{v}, \mathbf{j}'), \qquad \mathbf{j} = \mathbf{j}', \quad (4)$$

and

$$\mathrm{div}\mathbf{E} = 4\pi\rho, \qquad \mathrm{rot}\mathbf{E} = -\frac{\partial \mathbf{H}}{\partial t}, \qquad \mathrm{rot}\mathbf{H} = 4\pi\mathbf{j},$$
$$\mathrm{div}\mathbf{H} = 0, \qquad \mathrm{div}\mathbf{j} = 0, \quad (5)$$

where $\mathbf{E} = -\mathrm{grad}\varphi - \frac{\partial \mathbf{A}}{\partial t}$, $\mathbf{H} = \mathrm{rot}\mathbf{A}$. The latter case is referred to as the magnetic limit of the Maxwell equations [12].

Note that in the original work [12] the need to introduce the nonrelativistic transformations (2) or (4) was motivated by the fact that the conventional low velocity limit of (1) does not obey the group composition law, if transformation parameters are finite.

Below we demonstrate that (3) and (5) also hold invariant under the action of the ℓ–conformal Galilei group [2,3] for an arbitrary (half)integer parameter ℓ.

3. The ℓ–conformal Galilei group

The ℓ–conformal Galilei group [2,3] includes (temporal) translation, dilatation, and special conformal transformation, which form $SL(2,R)$ subgroup, as well as spatial translations, Galilei boosts and higher order constant accelerations. Structure relations of the corresponding Lie algebra involve an arbitrary (half)integer parameter ℓ, which specifies $2\ell - 1$ acceleration generators at hand.

In a nonrelativistic spacetime parametrized by a temporal variable t and spatial coordinates x_i, $i = 1, 2, 3$, the $SL(2,R)$–transformations are realized as follows

$$t' = \frac{\alpha t + \beta}{\gamma t + \delta}, \qquad x_i' = \left(\frac{\partial t'}{\partial t}\right)^\ell x_i, \quad (6)$$

where α, β, γ, δ are real parameters obeying $\alpha\delta - \beta\gamma = 1$ and ℓ is a (half)integer number, while (no sum over n below)

$$t' = t, \qquad x_i' = x_i + a_i^{(n)} t^n, \quad (7)$$

with $n = 0, \ldots, 2\ell$, specify the spatial translation $\left(a_i^{(0)}\right)$, the Galilei boost $\left(a_i^{(1)}\right)$ and higher order constant accelerations $\left(a_i^{(2)}, \ldots, a_i^{(2\ell)}\right)$.

Generators of the corresponding infinitesimal transformations

$$H = \frac{\partial}{\partial t}, \qquad D = t\frac{\partial}{\partial t} + \ell x_i \frac{\partial}{\partial x_i},$$
$$K = t^2\frac{\partial}{\partial t} + 2\ell t x_i \frac{\partial}{\partial x_i}, \qquad C_i^{(n)} = t^n \frac{\partial}{\partial x_i}$$

obey the structure relations of the ℓ–conformal Galilei algebra [2,3] [5]

$$[H, D] = H, \qquad [H, K] = 2D, \qquad [D, K] = K,$$
$$[H, C_i^{(n)}] = nC_i^{(n-1)}, \quad [D, C_i^{(n)}] = (n-l)C_i^{(n)}, \quad [K, C_i^{(n)}] = (n-2l)C_i^{(n+1)}, \quad (8)$$

where $n = 0, \ldots, 2\ell$, $i = 1, 2, 3$, and ℓ is a (half)integer number. The fact that ℓ is (half)integer guarantees that the algebra is finite–dimensional (mind the last bracket in (8)).

If one disregards the $SL(2,R)$–transformations, the range of index n carried by the vector generators $C_i^{(n)}$ is no longer linked to ℓ. Taking n to be a natural number greater than one, one obtains a transformation which links an inertial frame of reference to that moving with a constant acceleration of order $n-1$.

4. The ℓ–conformal Galilei symmetry of the Galilean electromagnetism

Let us demonstrate that the ℓ–conformal Galilei group is a symmetry of Eq. (3).

Transformation law of the charge density $\rho(t,x)$ is obtained by fixing a value of the temporal variable t and demanding the charge contained within a three–dimensional volume element V to be invariant under (6), (7)

$$\int_{V'} d^3x' \rho'(t', x') = \int_V d^3x \rho(t, x).$$

This gives

$$\rho(t,x) = \left(\frac{\partial t'}{\partial t}\right)^{3\ell} \rho'(t', x')$$

for the $SL(2,R)$–transformation and

$$\rho(t,x) = \rho'(t', x')$$

for the transformation involving parameters $a_i^{(n)}$.

[3] In this section, we use notations similar to those in [14]. In particular, $x^\mu = (ct, x^i)$, $\partial_\mu = \frac{\partial}{\partial x^\mu}$, $\eta_{\mu\nu} = \mathrm{diag}(+,-,-,-)$, $A^\mu = (\varphi, \mathbf{A})$, $A_\mu = (\varphi, -\mathbf{A})$, $j^\mu = (c\rho, \mathbf{j})$, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, $F_{0i} = -\partial_i\varphi - \frac{1}{c}\partial_t A_i := E_i$, $F_{ij} = -\partial_i A_j + \partial_j A_i := -\epsilon_{ijk}H_k$, with $i, j, k = 1, 2, 3$, where c is the speed of light.

[4] The nonrelativistic limit of the Lorentz boost transformation acting in space–time yields the conventional Galilei boost $t = t'$, $x^i = x'^i + v^i t'$.

[5] In arbitrary spatial dimension and for half-integer ℓ the algebra admits one central extension. As far as dynamical realizations are concerned, the central charge is usually interpreted as the mass of a (higher derivative) particle. In two spatial dimensions and for integer ℓ another central extension is allowed, which is usually attributed to the noncommutativity of space or external electric/magnetic fields present. For a detailed discussion of the central charges and further references to the original literature see e.g. [15,16].

Taking into account the identities

$$\frac{\partial}{\partial t}=\left(\frac{\partial t'}{\partial t}\right)\frac{\partial}{\partial t'}+\left(\frac{\partial x'_i}{\partial t}\right)\frac{\partial}{\partial x'_i},\qquad \frac{\partial}{\partial x_i}=\left(\frac{\partial t'}{\partial x_i}\right)\frac{\partial}{\partial t'}+\left(\frac{\partial x'_j}{\partial x_i}\right)\frac{\partial}{\partial x'_j},\quad (9)$$

and demanding the invariance of the continuity equation $\frac{\partial\rho}{\partial t}+\mathrm{div}\mathbf{j}=0$, one can then establish how the current density is transformed

$$j_i(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{2\ell+1}j'_i(t',x')-\ell\left(\frac{\partial t'}{\partial t}\right)^{2\ell-1}\left(\frac{\partial^2 t'}{\partial t^2}\right)x'_i\rho'(t',x'),$$

$$j_i(t,x)=j'_i(t',x')-nt^{n-1}\rho'(t',x')a_i^{(n)}. \quad (10)$$

Given the transformation laws of the charge density above, it is straightforward to verify that $\mathrm{div}\mathbf{E}=4\pi\rho$ holds invariant provided

$$E_i(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{2\ell}E'_i(t',x'),\qquad E_i(t,x)=E'_i(t',x'), \quad (11)$$

where the former relation specifies the $SL(2,R)$–transformation of the electric field strength, while the latter links to the transformations involving $a_i^{(n)}$. Taking into account the identities (9) and applying (6), (7), (11) to $\mathrm{rot}\mathbf{E}=0$, one can establish the invariance of the latter equation as well.

Finally, demanding the invariance of $\mathrm{rot}\mathbf{H}=\frac{\partial\mathbf{E}}{\partial t}+4\pi\mathbf{j}$ and $\mathrm{div}\mathbf{H}=0$, one can determine how the magnetic field strength transforms under the ℓ–conformal Galilei group (no sum over n below)

$$H_i(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{\ell+1}H'_i(t',x')-\ell\left(\frac{\partial t'}{\partial t}\right)^{\ell-1}\left(\frac{\partial^2 t'}{\partial t^2}\right)\epsilon_{ijk}x'_jE'_k(t',x'),$$

$$H_i(t,x)=H'_i(t',x')-nt^{n-1}\epsilon_{ijk}a_j^{(n)}E'_k(t',x'), \quad (12)$$

where ϵ_{ijk} is the totally antisymmetric symbol with $\epsilon_{123}=1$.

Note that, when verifying the invariance of the equations of motion for the magnetic field strength, the conditions $\mathrm{rot}\mathbf{E}=0$, $\mathrm{div}\mathbf{E}=4\pi\rho$ have been used as well. To put it in other words, the full system of Eq. (3) holds invariant rather than each individual equation.

The magnetic limit of the Maxwell equations can be analyzed in a similar fashion. It is straightforward to verify that Eq. (5) hold invariant under the $SL(2,R)$–transformations

$$\rho(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{2\ell+1}\rho'(t',x')+\frac{\ell}{4\pi}\left(\frac{\partial t'}{\partial t}\right)^{2\ell-1}\left(\frac{\partial^2 t'}{\partial t^2}\right)\epsilon_{ijk}x'_j\frac{\partial H'_k(t',x')}{\partial x'_i},$$

$$j_i(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{3\ell}j'_i(t',x'),$$

$$E_i(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{\ell+1}E'_i(t',x')+\ell\left(\frac{\partial t'}{\partial t}\right)^{\ell-1}\left(\frac{\partial^2 t'}{\partial t^2}\right)\epsilon_{ijk}x'_jH'_k(t',x'),$$

$$H_i(t,x)=\left(\frac{\partial t'}{\partial t}\right)^{2\ell}H'_i(t',x'),$$

as well as under the transformations involving vector parameters $a_i^{(n)}$ (no sum over n below)

$$\rho(t,x)=\rho'(t',x')+\frac{1}{4\pi}nt^{n-1}\epsilon_{ijk}a_j^{(n)}\frac{\partial H'_k(t',x')}{\partial x'_i},\qquad j_i(t,x)=j'_i(t',x'),$$

$$E_i(t,x)=E'_i(t',x')+nt^{n-1}\epsilon_{ijk}a_j^{(n)}H'_k(t',x'),\qquad H_i(t,x)=H'_i(t',x'). \quad (13)$$

Somewhat surprisingly, the invariance of the charge contained within a three–dimensional volume element $\int_V d^3x\rho(t,x)$ under the $SL(2,R)$–transformations then requires $\ell=1$ and it also implies that both $H_i(t,x)$ and $\epsilon_{ijk}x_jH_k(t,x)$ should vanish at spatial infinity.

When verifying the symmetries of Eqs. (3) and (5) above the identities

$$\epsilon_{ijk}\epsilon_{lpk}=\delta_{il}\delta_{jp}-\delta_{ip}\delta_{jl},\qquad \epsilon_{ijk}\epsilon_{ijp}=2\delta_{kp}$$

proved useful.

It is important to stress that, discarding the $SL(2,R)$–transformations, both the electric and magnetic limit equations and the charge $\int_V d^3x\rho(t,x)$ still hold invariant under a transformation which links an inertial frame of reference to that moving with a constant acceleration of order up to $2\ell-1$. This latter fact will prove important for our discussion in the next section.

5. Discussion

Due credit given to the elegant Eqs. (3) and (5), some of their properties deserve a critical discussion. For definiteness, below we focus on the electric limit. The magnetic limit can be considered likewise.

Firstly, because $\frac{\partial\mathbf{H}}{\partial t}$ does not contribute to (3), it appears problematic to regard (3) as a fully consistent set of *evolutionary equations* to determine $\mathbf{E}$ and $\mathbf{H}$.

Secondly, for a similar reason the magnetic field strength $\mathbf{H}$ does not contribute to the conserved energy

$$\frac{1}{8\pi}\int d^3x(\mathbf{E},\mathbf{E}), \quad (14)$$

which characterizes the system in the absence of sources.[6]

Thirdly, as was shown in the previous section, Eq. (3) hold invariant under a transformation which links an inertial frame of reference to that moving with a constant acceleration of order up to $2\ell-1$. In particular, given a solution to (3) describing constant homogeneous electric and magnetic fields in a vacuum

$$E_i(t,x)=\alpha_i,\qquad H_i(t,x)=\beta_i,\qquad \rho(t,x)=0,\qquad j_i(t,x)=0, \quad (15)$$

where α_i, β_i are arbitrary vectors, the transformations yield another solutions to the same equations

$$E_i(t,x)=\alpha_i,\qquad H_i(t,x)=\beta_i+t^{n-1}\gamma_i,\qquad \rho(t,x)=0,\qquad j_i(t,x)=0, \quad (16)$$

where γ_i is an arbitrary vector, t is the temporal variable, and $n=1,\ldots,2\ell$. Furthermore, t^{n-1} in (16) can actually be replaced with an arbitrary function $f(t)$, for example, e^t (see below). An exponential growth over time of a magnetic field strength in the direction specified by a vector γ_i in an empty space is certainly an unphysical phenomenon pointing at potential dynamical instability of the system. The fact that the magnetic field strength can grow unbounded and yet it does not contribute to the conserved energy (14) calls for a better understanding of count of physical degrees of freedom in phase space. A natural way to proceed is to implement the constrained dynamics analysis for the Lagrangian formulations in [13,17–21]. The key point is to understand whether there is a propagating degree of freedom supporting the magnetic field strength.[7]

Fourthly, in the literature on the Galilean electromagnetism, it is widely underestimated that the Galilei boost (2) is actually a particular instance of a more general symmetry

$$t'=t,\qquad x'_i=x_i+\alpha_i(t),$$

$$\rho(t,x)=\rho'(t',x'),\qquad j_i(t,x)=j'_i(t',x')-\dot\alpha_i(t)\rho'(t',x'),$$

$$E_i(t,x)=E'_i(t',x'),\qquad H_i(t,x)=H'_i(t',x')-\epsilon_{ijk}\dot\alpha_j(t)E'_k(t',x'), \quad (17)$$

in which $\alpha_i(t)$ is an arbitrary differentiable function of time and the dot designates the derivative with respect to t. The conventional Galilei boost emerges at $\alpha_i(t)=-tv_i$. In the absence of sources, the transformation (17) was first revealed in [17], where it was also argued that the vacuum equations of motion are time reparametrization invariant and one can perform a time–dependent spatial dilatations. In the presence of sources, one can establish the following symmetry transformation rules

$$t'=f(t),\qquad x'_i=g(t)x_i,$$

$$\rho(t,x)=g^3(t)\rho'(t',x'),\qquad j_i(t,x)=g^2(t)\dot f(t)j'_i(t',x')-g(t)\dot g(t)x'_i\rho'(t',x'),$$

[6] It is assumed here that $\epsilon_{ijk}E_jH_k$ vanishes at spatial infinity.

[7] The author thanks an anonymous reviewer for calling his attention to this fact.

$$E_i(t,x) = g^2(t)E_i'(t',x'), \quad H_i(t,x) = g(t)\dot{f}(t)H_i'(t',x') - \dot{g}(t)\epsilon_{ijk}x_j'E_k'(t',x'),$$

where $f(t)$ and $g(t)$ are arbitrary differentiable functions of time. The ℓ–conformal Galilei symmetry described above thus appears to be a finite dimensional subgroup of the larger infinite dimensional symmetry group.

Finally, in order to resolve some physical issues, like the problem with the definition of the conserved energy above, in the original work [12] it was suggested to double the number of fields and consider Eqs. (3) and (5) simultaneously by assigning the subscript "e" to the fields entering (3) and "m" to those contributing to (5). The Lorentz force was modified accordingly by allowing mixed contributions. Yet, the problem with potentially unstable solutions persists to the extended framework.

Turning to possible further developments, it would be interesting to explore whether the ℓ-conformal Galilei symmetry discussed in this work can be realized at the level of action functionals studied in [13, 17–21].

Data availability

No data was used for the research described in the article.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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