Lorentz symmetry breaking in \( \mathcal{N} = 2 \) superspace

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Abstract – In this paper, we will study the deformation of a three-dimensional theory with \( \mathcal{N} = 2 \) supersymmetry. This theory will be deformed by the presence of a constant vector field. This deformation will break the Lorentz symmetry. So, we will analyse this theory using \( \mathcal{N} = 2 \) aether superspace. The \( \mathcal{N} = 2 \) aether superspace will be obtained from a deformation of the usual \( \mathcal{N} = 2 \) superspace. This will be done by deforming the generators of the three-dimensional \( \mathcal{N} = 2 \) supersymmetry. After analysing this deformed superalgebra, we will derive an explicit expression for the superspace propagators in this deformed superspace. Finally, we will use these propagators for performing perturbative calculations.

Introduction. – Lorentz symmetry is one of the most important symmetries in nature. However, there are strong theoretical indications from various approaches to quantum gravity, that this might only be an effective symmetry in nature. Initially the study in this area was motivated by developments in string theory. This is because when perturbative string vacuum is unstable, Lorentz symmetry will be naturally broken [1,2]. This is because in this case certain tensors acquire non-zero vacuum expectation value and this introduces a preferential direction in spacetime. There is a deep relation between string theory and non-commutativity and this can also lead to breaking of Lorentz symmetry [3,4]. In fact, the breaking of Lorentz symmetry at Planck scale is generally expected to arise in most theories of quantum gravity [5]. Furthermore, even though gravity is not renormalizable, it can be made renormalizable by adding higher-order curvature invariants to the original gravitational Lagrangian [6]. However, this spoils the unitarity of the resultant theory [7]. One way out of this problem is to take a different Lifshitz scaling for space and time and thus add terms containing higher-order spatial derivatives without adding any term containing higher-order temporal derivative. This approach to quantum gravity is called Horava-Lifshitz gravity and in it Lorentz invariance is broken in the high energy limit of the theory [8,9]. Lorentz symmetry breaking has also been studied in the context of loop quantum gravity [10,11]. Lorentz symmetry breaking has also been studied in the context of modified dispersion relations and this approach has led to the development of doubly special relativity [12,13]. In this theory both the velocity of light and the Planck energy are invariant quantities. This assumption naturally incorporates the existence of a maximal momentum and modifies the first quantized field theory. This modification for General Relativity has also been studied and this has led to the development of Gravity’s Rainbow [14,15]. Furthermore, the Lorentz symmetry breaking can be used as a possible way to solve the problem of time in quantum gravity [16]. Hence, there are various motivations to study Lorentz symmetry breaking.

The Lorentz symmetry breaking has also been studied in supersymmetric theories [17,18]. In fact, it is possible to analyse Lorentz symmetry violation in which a sub-group of the Lorentz group is preserved. Thus, for example Lorentz symmetry can be violated without violating a three-dimensional rotation subgroup by choosing a background time-like vector field. Such theories have been studied in detail and are called aether theories [19,20]. These theories have also been applied in the study of Lorentz symmetry violating models of electrodynamics and in these models a Carroll-Field Jackiw term is added to the original Lagrangian [21]. This term arises as a quantum correction if a Lorentz-violating axial term is included in the fermionic sector of the original Lagrangian [22]. This term also breaks the CPT symmetry [23]. It is natural to try to study the supersymmetric theories in aether superspace. In fact, Lorentz symmetry breaking can be incorporated by deforming the structure of the generators of supersymmetry and this in turn...
modifies there superalgebra [24]. Then superderivatives can be constructed such that they anticommute with these modified generators of supersymmetry. Some attempts to implement this approach at tree level have also been made [25]. In fact, Lagrangian for supersymmetric scalar field theory with $N = 1$ supersymmetry has been constructed using eather superspace [26]. Furthermore, Lagrangian for $N = 1$ Abelian gauge theories has also been constructed using eather superspace [27]. In this paper, we will extend this work and study a supersymmetric field theory with $N = 2$ in aether superspace. We will also obtain explicit expression for propagators for this theory and use them for performing perturbative calculations.

$N = 2$ aether superspace. – In this section we will study aether superspace formalism for three-dimensional theories. We will perform the calculations using $N = 2$ superspace formalism in three dimensions. In order to do that, we will first consider a constant vector field $v^\mu = (v^0, v^i)$, such that, $|v|^2 = v^\mu v_\mu$. Now $|v|^2 = 1$ for space-like, $|v|^2 = -1$ for time-like and $|v|^2 = 0$ for light-like cases [26]. Furthermore, this constant vector field can be used to construct a tensor field $k_{\mu\nu} = \alpha v_\mu v_\nu$, for an arbitrary parameter $\alpha$. It may be noted that in the space-like case, we have $E^2 = p^\mu p_\mu + m^2 + (2\alpha + \alpha^2)v^\mu p_\mu$, so the dynamics can be consistently defined for $\alpha > 0$ and for $\alpha < 0$, if $|\alpha| \ll 1$. However, for $\alpha < 0$ the theory turns out to be degenerate or unstable. In the light-like case, we have $E^2 = 2 = p^\mu v_\mu + m^2$, and so the dynamics is consistent for all values of $\alpha$, except $\alpha = 1$. Finally, for the light-like case, we have $E(1 - 2\alpha) = -2\alpha \sqrt{p^\mu p_\mu} \pm \sqrt{p^\mu p_\mu(1 + 2\alpha + 4\alpha^2) + m^2}$. So, the dynamics is consistent for $\alpha \ll 1$

Now the supersymmetry can be deformed using this vector field, in such a way that Lorentz symmetry is broken without affecting any supersymmetry. Thus, we can construct two supercharges in three dimensions,

\[
Q_{1a} = \partial_1a - (\gamma^\mu \partial_\mu \theta_1)_a - (\gamma^\nu k_{\mu\nu} \partial_\nu \theta_1)_a, \\
Q_{2a} = \partial_2a - (\gamma^\mu \partial_\mu \theta_2)_a - (\gamma^\nu k_{\mu\nu} \partial_\nu \theta_2)_a.
\]

Now these supercharges satisfy the following superalgebra:

\[
\begin{align*}
\{Q_{1a}, Q_{1b}\} &= 2(\gamma^\mu \partial_\mu k_{\mu\nu} \partial_\nu\theta)_{ab}, \\
\{Q_{1a}, Q_{2b}\} &= 0, \\
\{Q_{2a}, Q_{2b}\} &= 2(\gamma^\mu \partial_\mu k_{\mu\nu} \partial_\nu\theta)_{ab}.
\end{align*}
\]

We can construct superderivatives which commute with these generators of $N = 2$ supersymmetry, $\{D_{1a}, Q_{1a}\} = \{D_{2a}, Q_{1a}\} = \{D_{1a}, Q_{2a}\} = \{D_{2a}, Q_{2a}\} = 0$. These superderivatives can be written as

\[
\begin{align*}
D_{1a} &= \partial_1a + (\gamma^\mu \partial_\mu \theta_1)_a + (\gamma^\nu k_{\mu\nu} \partial_\nu \theta_1)_a, \\
D_{2a} &= \partial_2a + (\gamma^\mu \partial_\mu \theta_2)_a + (\gamma^\nu k_{\mu\nu} \partial_\nu \theta_2)_a,
\end{align*}
\]

and they satisfy

\[
\begin{align*}
\{D_{1a}, D_{1b}\} &= -2(\gamma^\mu \partial_\mu k_{\mu\nu} \partial_\nu\theta)_{ab}, \\
\{D_{1a}, D_{2b}\} &= 0, \\
\{D_{2a}, D_{2b}\} &= -2(\gamma^\mu \partial_\mu k_{\mu\nu} \partial_\nu\theta)_{ab}.
\end{align*}
\]

We can represent any supersymmetric theory containing two superderivatives $D_{1a}$ and $D_{2a}$ equivalent by two other derivatives which are linear combinations of these original superderivatives,

\[
\begin{pmatrix}
D_{3a} \\
D_{4a}
\end{pmatrix} =
\begin{pmatrix}
x_{11} & x_{12} \\
x_{21} & x_{22}
\end{pmatrix}
\begin{pmatrix}
D_{1a} \\
D_{2a}
\end{pmatrix},
\]

where $x_{ij}$ are c-numbers such that, $x_{11}x_{22} - x_{12}x_{21} \neq 0$, so that $D_{3a}$ and $D_{4a}$ form a valid representation of the supersymmetry. For a supersymmetric field theory, the Jacobian of this transformation can be absorbed in field redefinition. Furthermore, it may be noted that as $D_{3a}$ and $D_{4a}$ are linear combinations of $D_{1a}$ and $D_{2a}$, so they will also contain $k_{\mu\nu}$-dependent terms. It may be noted that it also is possible to analyse a non-trivially mixing of these supersymmetric derivatives [28]. Now we will use a specific form of this transformation, such that [29,30]

\[
\theta_a = \frac{1}{\sqrt{2}} [\theta_{1a} + i \theta_{2a}], \quad \bar{\theta}_a = \frac{1}{\sqrt{2}} [\theta_{1a} - i \theta_{2a}].
\]

So, the derivatives $D_a$ and $\bar{D}_a$ read

\[
\begin{align*}
D_a &= \partial_a + i(\gamma^\mu \partial_\mu \theta)_{a} + i(\gamma^\nu k_{\mu\nu} \partial_\nu\theta)_{a}, \\
\bar{D}_a &= \partial_a + i(\gamma^\mu \partial_\mu \bar{\theta})_{a} + i(\gamma^\nu k_{\mu\nu} \partial_\nu\bar{\theta})_{a}.
\end{align*}
\]

These superderivatives satisfy

\[
\begin{align*}
\{D_a, \bar{D}_b\} &= 2i(\gamma^\mu \partial_\mu)_{ab} + 2i(\gamma^\nu k_{\mu\nu} \partial_\nu\theta)_{ab}, \\
\{D_a, \bar{D}_b\} &= 0, \\
\{D_a, D_b\} &= 0.
\end{align*}
\]

We can now construct two supercharges $Q_a$ and $\bar{Q}_a$, such that these superderivatives commute with them, $\{Q_a, D_b\} = \{Q_a, \bar{D}_b\} = \{Q_a, D_a\} = \{Q_a, \bar{D}_a\} = 0$. These supercharges also can be used to represent $N = 2$ supersymmetry in three dimensions. We can represent them as follows:

\[
\begin{align*}
Q_a &= -i \partial_a - (\gamma^\mu \partial_\mu \theta)_{a} - (\gamma^\nu k_{\mu\nu} \partial_\nu\theta)_{a}, \\
\bar{Q}_a &= i \partial_a + (\gamma^\mu \partial_\mu \bar{\theta})_{a} + (\gamma^\nu k_{\mu\nu} \partial_\nu\bar{\theta})_{a}.
\end{align*}
\]

They satisfy

\[
\begin{align*}
\{Q_a, Q_b\} &= -2i(\gamma^\mu \partial_\mu k_{\mu\nu} \partial_\nu\theta)_{ab}, \\
\{Q_a, \bar{Q}_b\} &= 0, \\
\{Q_a, \bar{Q}_b\} &= 0.
\end{align*}
\]

Superfield theory in aether superspace. – In the previous section we analysed the supersymmetric algebra for three-dimensional aether superspace with $N = 2$ supersymmetry. In this section we will analyse the supersymmetric field theory for a three-dimensional theory with $N = 2$ supersymmetric in aether superspace. We can now define projections of a $N = 2$ superfield in this aether superspace as $D_a \Phi(y, \bar{\theta}) = 0$ and $D_a \phi(y, \theta) = 0$, where
\[ y^\mu = x^\mu + i\theta^\gamma \gamma^\mu. \]

Now we expand these superfields as

\[
\Phi = (\phi(y) + \sqrt{2} \theta \psi(y)) + \phi^2 f(y) \\
= i\sqrt{2} \theta \gamma^\mu \partial_\mu \phi(x) + i\gamma^\mu \theta k_{\mu \nu} \phi(x) \\
+ \frac{i}{\sqrt{2}} \theta^2 \theta^0 \partial_\mu \phi(x) + \frac{i}{\sqrt{2}} \theta^2 \gamma^0 k_{\mu \nu} \phi(x) \\
- \frac{1}{4} \theta^2 \theta^0 \theta^0 k_{\mu \nu} \phi(x) - \frac{1}{4} \theta^2 \theta^0 \theta^0 k_{\mu \nu} \phi(x) \\
- \frac{1}{4} \theta^2 \theta^0 \theta^0 k_{\mu \nu} \phi(x) + \theta^2 f(x),
\]

\[
\Phi = \phi^*(y) + \sqrt{2} \theta \psi(y) + \phi^2 f^*(y) \\
= \phi^*(x) + \sqrt{2} \theta \psi(x) + i\theta^\gamma \gamma^\mu \phi^*(x) \\
+ i \sqrt{2} \theta \gamma^\mu k_{\mu \nu} \phi^*(x) + \frac{i}{\sqrt{2}} \theta^2 \gamma^0 k_{\mu \nu} \phi(x) \\
- \frac{1}{4} \theta^2 \theta^0 \theta^0 k_{\mu \nu} \phi^*(x) - \frac{1}{4} \theta^2 \theta^0 \theta^0 k_{\mu \nu} \phi(x) \\
+ \theta^2 f^*(x).
\]

Now we can write the action for a supersymmetric field theory in this aether superspace as

\[
S = \frac{1}{2} \int d^3x [2d^2 \theta d^2 \psi \Phi + m^2 \phi^2 + m^2 \phi^2].
\]

This action can be expanded in component form as

\[
S = \int d^3x \left[ (\phi^* f) \left( -\partial^m \partial_m - 2\theta^\nu \partial^\nu k_{\mu \nu} - k^\mu \nu \partial_\nu \partial_\mu m \right) \right] \\
\times \left[ \phi^* f + \frac{1}{2} \psi \left( -i\gamma^\mu \partial_\mu - i\gamma^\mu \gamma^0 k_{\mu \nu} \phi \right) \right] \\
\times \left[ \psi \right].
\]

Now in general for any field, the generating functional for the Green’s functions is given by

\[
Z[J, J^*] = N \exp -i \int d^3x J K^{-1} J^*,
\]

where \( N \) is normalization constant. Thus, the Green’s function for any field can be written as \( iK^{-1} \). The \( K_B \) for the bosonic part is given by

\[
K_B = \left( -\partial^m \partial_m - 2\theta^\nu \partial^\nu k_{\mu \nu} - k^\mu \nu \partial_\nu \partial_\mu m \right),
\]

which can be inverted to obtain

\[
K_B^{-1} = \frac{1}{-\partial^m \partial_m - 2\theta^\nu \partial^\nu k_{\mu \nu} - k^\mu \nu \partial_\nu \partial_\mu m - m^2} \\
\times \left( 1 + \frac{1}{-m - \partial^m \partial_m - 2\theta^\nu \partial^\nu k_{\mu \nu} - k^\mu \nu \partial_\nu \partial_\mu m} \right).
\]

Similarly, \( K_F \) for the fermionic part is given by

\[
K_F = \left( -i\gamma^\mu \partial_\mu - i\gamma^\mu k_{\mu \nu} \phi \right) \\
\times \left( -m - i\gamma^\mu \partial_\mu - i\gamma^\mu k_{\mu \nu} \phi \right),
\]

which can be inverted to obtain

\[
K_F^{-1} = \frac{1}{-i\gamma^\mu \partial_\mu - i\gamma^\mu k_{\mu \nu} \phi - m^2} \\
\times \left( 1 + \frac{1}{-m - i\gamma^\mu \partial_\mu - i\gamma^\mu k_{\mu \nu} \phi} \right).
\]

Now using \( K_B^{-1} \) and \( K_F^{-1} \), we can solve the two-point functions for all the component fields in this theory,

\[
\langle 0|\phi(x)\phi^*(y')|0 \rangle = \frac{i}{m^2} \frac{m^2}{m^2 - m^2} \\
\times \delta^3(x - x').
\]
where
\[ M_1(x, \tilde{\theta}, \tilde{\theta}) = \frac{\delta^2(\theta - \theta') \delta^2(\tilde{\theta} - \tilde{\theta}') \delta^3(x - x')}{-\partial^2 \partial' - 2\partial^2 \partial' k_{\mu} - k^{\tau\nu} k_{\tau\nu}' \partial_{\tau} \partial'_{\nu} - m^2}. \] (22)

Similarly, using the fact that \( \langle 0 | \Phi(y, \theta) \Phi(y', \theta') | 0 \rangle = \bar{\theta}^2(0) \phi(y) f^*(y') | 0 \rangle + \bar{\theta}^2(0) f(y) \phi^*(y') | 0 \rangle + 2\theta^a \bar{\theta}^b (\bar{v}_a(y) \bar{v}_b(y') | 0 \rangle \), we obtain
\[ \langle 0 | \Phi(x, \theta, \bar{\theta}) \Phi(x', \theta', \bar{\theta}') | 0 \rangle = \frac{im}{4} D^2 M_2(x, \theta, \bar{\theta}), \] (23)
where
\[ M_2(x, \theta, \bar{\theta}) = \frac{\delta^2(\theta - \theta') \delta^2(\tilde{\theta} - \tilde{\theta}') \delta^3(x - x')}{-\partial^2 \partial' - 2\partial^2 \partial' k_{\mu} - k^{\tau\nu} k_{\tau\nu}' \partial_{\tau} \partial'_{\nu} - m^2}. \] (24)

Finally, using the fact that \( \langle 0 | \Phi(y, \theta) \Phi(y', \theta') | 0 \rangle = \theta^a \bar{\theta}^b (\bar{v}_a(y) \bar{v}_b(y') | 0 \rangle + \theta^b \bar{\theta}^a (\bar{v}_a(y) \bar{v}_b(y') | 0 \rangle + 2\theta^a \bar{\theta}^b (\bar{v}_a(y) \bar{v}_b(y') | 0 \rangle \), we obtain
\[ \langle 0 | \Phi(x, \theta, \bar{\theta}) \Phi(x', \theta', \bar{\theta}') | 0 \rangle = \frac{\delta^3(y - y') - N(\theta, \theta') \delta^3(y - y')}{-\partial^2 \partial' - 2\partial^2 \partial' k_{\mu} - k^{\tau\nu} k_{\tau\nu}' \partial_{\tau} \partial'_{\nu} - m^2}, \] (25)
where we have defined \( N(\theta, \theta') = -2\theta^a \bar{\theta}^b \gamma_{ab}(\partial_{\mu} + k_{\mu} \partial_{\nu}) + i\partial^2 \theta^2 (\partial_{\mu} + k_{\mu} \partial_{\nu}) + 2\theta^a \bar{\theta}^b (\bar{v}_a(y) \bar{v}_b(y') | 0 \rangle. \) This can be simplified to obtain the following expression:
\[ \langle 0 | \Phi(x, \theta, \bar{\theta}) \Phi(x', \theta', \bar{\theta}') | 0 \rangle = \frac{i}{16} D^2 D^2 M_3(x, \theta, \bar{\theta}), \] (26)
where
\[ M_3(x, \theta, \bar{\theta}) = \frac{\delta^2(\theta - \theta') \delta^2(\tilde{\theta} - \tilde{\theta}') \delta^3(x - x')}{-\partial^2 \partial' - 2\partial^2 \partial' k_{\mu} - k^{\tau\nu} k_{\tau\nu}' \partial_{\tau} \partial'_{\nu} - m^2}. \] (27)

**Interactions in aether superspace.** These superspace propagator’s can now be used for analyzing superspace perturbations. So, we can calculate the loop correction to the superspace propagator’s for different interaction terms. We can start by calculating the one-loop corrections for \( \langle 0 | \Phi(x, \theta, \bar{\theta}) \Phi(x', \theta', \bar{\theta}') | 0 \rangle \), when the interaction term is of the form \( \mathcal{L}_{int} = \lambda \bar{\theta}^2 \delta^3 \mathcal{F} \). The one-loop corrections to \( \langle 0 | \Phi(x, \theta, \bar{\theta}) \Phi(x', \theta', \bar{\theta}') | 0 \rangle \), can be written as
\[ -2\lambda^2 \left( \frac{i}{16} \right)^4 \int d^3 p_1 d^3 p_2 d^2 \theta_1 d^2 \bar{\theta}_2 D^2 \frac{\delta^2(\theta_1 - \theta_2) \delta^2(\tilde{\theta}_1 - \tilde{\theta}_2)}{\mathcal{A}(p_1 - p_2, k, m)} \times \frac{D^2 \delta^2(\theta_1 - \theta_2) \delta^2(\tilde{\theta}_1 - \tilde{\theta}_2)}{p_1^2 + 2k_{\mu} p_1 \mu p_1 - k^{\tau\nu} k_{\tau\nu}' p_1 \mu - m^2} \times \frac{D^2 \delta^2(\theta_2 - \theta_1) \delta^2(\tilde{\theta}_2 - \tilde{\theta}_1)}{p_2^2 + 2k_{\mu} p_2 \mu p_2 - k^{\tau\nu} k_{\tau\nu}' p_2 \mu - m^2} \times \frac{D^2 \delta^2(\theta_2 - \theta_1) \delta^2(\tilde{\theta}_2 - \tilde{\theta}_1)}{p_2^2 + 2k_{\mu} p_2 \mu p_2 + k^{\tau\nu} k_{\tau\nu}' p_2 \mu - m^2} \times \frac{1}{p_1^2 + 2k_{\mu} p_1 \mu p_1 + k^{\tau\nu} k_{\tau\nu}' p_1 \mu - m^2} \times \frac{1}{p_2^2 + 2k_{\mu} p_2 \mu p_2 + k^{\tau\nu} k_{\tau\nu}' p_2 \mu - m^2} = 0. \] (28)

So, even the contributions from this non-trivial diagram vanish. Thus, the vacuum energy for the aether superspace is still zero even at two-loops. It may be noted that the quantum fluctuations do not break the supersymmetry in three-dimensional \( \mathcal{N} = 2 \) aether superspace. It would be interesting to analyse general non-renormalization theorems for the aether superspace.

**Conclusion.** In this paper, we analysed a three-dimensional supersymmetric field theory with \( \mathcal{N} = 2 \) supersymmetry in aether superspace. In this superspace the Lorentz symmetry was broken without breaking any supersymmetry. We analysed this model in a representation where a mixing between the original generators of \( \mathcal{N} = 2 \) supersymmetry occurred. We then obtained an explicit expression for supercharges and superderivatives in this representation of \( \mathcal{N} = 2 \) supersymmetry. We used these superderivatives in aether superspace to...
derive explicit expressions for propagators for our model. Finally, we used these propagators for performing some perturbative calculations. We thus observed that there is no contribution to the vacuum energy from one-loop and two-loops graphs. It was argued that the supersymmetry is not broken by quantum fluctuations in aether superspace, at least till two loops. It will be interesting to perform a similar calculation for models with higher amount of supersymmetry. Thus, we could analyse a four-dimensional scalar superfield model in $N = 2$ aether superspace. It will also be interesting to study the Lorentz symmetry breaking by adding CPT odd Lorentz-breaking terms to the components of superfields and keeping the superalgebra undeformed. Furthermore, we can also add explicit Lorentz breaking terms to the superspace action. The action derived from such an approach will contain higher derivative terms.

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