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Spin degrees of freedom incorporated in conformal group: Introduction of an intrinsic momentum operator

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Abstract

Considering spin degrees of freedom incorporated in the conformal generators, we introduce an intrinsic momentum operator π_μ, which is feasible for the Bhabha wave equation. If a physical state ψ_{ph} for spin s is annihilated by the π_μ, the degree of ψ_{ph}, $\deg\psi_{\text{ph}}$, should equal twice the spin degrees of freedom, $2(2s+1)$ for a massive particle, where the multiplicity 2 indicates the chirality. The relation $\deg\psi_{\text{ph}} = 2(2s+1)$ holds in the representation $\mathrm{R}_5(s,s)$, irreducible representation of the Lorentz group in five dimensions.

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1 Introduction

Conformal symmetry [1] has many applications in string theory and critical phenomena in condensed matter and statistical physics. For a scalar field, the conformal generators are composed of dilatation D, momentum P_μ, special conformal K_μ, and angular momentum $L_{\mu\nu}$. For a multicomponent field Φ, where spin degrees of freedom is incorporated as $L_{\mu\nu} \to L_{\mu\nu} + s_{\mu\nu}$, the D and K_μ are generalized as $D \to D + \Delta$ and $K_\mu \to K_\mu + \kappa_\mu$, while the P_μ, in an ordinary context [1], remains unchanged as $P_\mu \to P_\mu$. The unchangeability of P_μ may be because Φ transforms as a scalar under spacetime translation. If we assume that $\Phi(x) \to \Phi'(x') = \Phi(x)$ under $x \to x' = x + a$, that is, $\Phi'(x) = \Phi(x-a) = \mathrm{e}^{-a\cdot P}\Phi(x)$, we find it unnecessary to introduce an intrinsic momentum operator π_μ as $P_\mu \to P_\mu + \pi_\mu$. Even if we admit the scalar property of $\Phi(x)$ under $x \to x + a$, we can introduce π_μ in such a way that the π_μ may annihilate physical states.

This paper aims to introduce such an intrinsic momentum operator π_μ, to find that π_μ can realize for a matrix structure in parafermion-based Dirac-like equations, such as spin-1 Kemmer equation [2], and more generally, Bhabha equation [3]. In Sec. 2, we give some preliminaries concerning the conformal algebra, together with its Casimir operator. In Secs. 3-5, we deal with the π_μ in the case of spin $\frac{1}{2}, 1, \frac{3}{2}$, respectively. We devote Sec. 6 to the summary.

2 Preliminaries

We begin with the commutation relations between the intrinsic conformal generators Δ, π_μ, κ_μ, and $s_{\mu\nu}$, corresponding to D, P_μ, K_μ, and $L_{\mu\nu}$, respectively. If the intrinsic conformal generators satisfy the same commutation relations as ordinary conformal generators, we can write the non-vanishing commutation relations as

$$[\Delta, \pi_\mu] = \mathrm{i}\pi_\mu\,, \quad [\Delta, \kappa_\mu] = -\mathrm{i}\kappa_\mu\,, \quad [\kappa_\mu, \pi_\nu] = 2\mathrm{i}(g_{\mu\nu}\Delta - s_{\mu\nu})\,, \tag{1}$$

$$[\pi_\rho, s_{\mu\nu}] = \mathrm{i}(g_{\rho\mu}\pi_\nu - g_{\rho\nu}\pi_\mu)\,, \quad [\kappa_\rho, s_{\mu\nu}] = \mathrm{i}(g_{\rho\mu}\kappa_\nu - g_{\rho\nu}\kappa_\mu)\,, \tag{2}$$

$$[s_{\mu\nu}, s_{\rho\sigma}] = \mathrm{i}(g_{\nu\rho}s_{\mu\sigma} + g_{\mu\sigma}s_{\nu\rho} - g_{\mu\rho}s_{\nu\sigma} - g_{\nu\sigma}s_{\mu\rho})\,, \tag{3}$$

while the vanishing commutation relations are given by

$$[\Delta, s_{\mu\nu}] = [\pi_\mu, \pi_\nu] = [\kappa_\mu, \kappa_\nu] = 0\,. \tag{4}$$

It should be remarked that (1)-(4) are invariant under the scaling of π_μ and κ_μ, and also under the substitution between π_μ and κ_μ as

$$(\Delta, \pi_\mu, \kappa_\mu, s_{\mu\nu}) \to (\Delta, \lambda\pi_\mu, \lambda^{-1}\kappa_\mu, s_{\mu\nu})\,, \tag{5}$$

$$(\Delta, \pi_\mu, \kappa_\mu, s_{\mu\nu}) \to (-\Delta, \kappa_\mu, \pi_\mu, s_{\mu\nu})\,, \tag{6}$$

where $\lambda \in \mathbb{C} \setminus \{0\}$, and use has been made of $s_{\nu\mu} = -s_{\mu\nu}$ in (6). Note that (5) represents the "chiral" transformation $g \to g' = \mathrm{e}^{\theta\Delta} g \mathrm{e}^{-\theta\Delta}$ ($g \in \{\Delta, \pi_\mu, \kappa_\mu, s_{\mu\nu}\}$), where $\lambda = \mathrm{e}^{\mathrm{i}\theta}$.

To check the irreducibility of the representation for the conformal group, it may be available to obtain the Casimir operator C. Note that although the C is invariant under (5) due to the chiral transformation, the invariance of C under (6) is somewhat naive. For simplicity, we consider $(3+1)$ spacetime dimensions, where the conformal algebra is isomorphic to $\mathfrak{so}(4,2)$ [1]. In this case, the order of C is given by $2,3,4$, as in the case of $\mathfrak{so}(6)$ [4]. Explicitly, we have $C = C_2, C_3, C_4$ (the index i in C_i represents the order) as [5]

$$\begin{aligned} C_2 &= \frac{1}{2}s_{\mu\nu}s^{\mu\nu} + \frac{1}{2}\{\kappa_\mu, \pi^\mu\} - \Delta^2\,, \qquad C_3 = \epsilon^{\mu\nu\rho\sigma}\big(\Delta s_{\mu\nu} + \{\kappa_\mu, \pi_\nu\}\big)s_{\rho\sigma}\,, \\ C_4 &= \frac{1}{2}\mathcal{J}_{\mu\nu}\mathcal{J}^{\mu\nu} - \frac{1}{2}\{\mathcal{J}_{K,\mu}, \mathcal{J}_P^\mu\} - \frac{1}{16}\mathcal{J}^2\,, \end{aligned} \tag{7}$$

where $\mathcal{J}^{\mu\nu}, \mathcal{J}_K^\mu, \mathcal{J}_P^\mu$, and $\mathcal{J}$ are given by $\mathcal{J}^{\mu\nu} = \epsilon^{\mu\nu\rho\sigma}\left(\Delta s_{\rho\sigma} + \frac{1}{2}\{\kappa_\rho, \pi_\sigma\}\right)$, $\mathcal{J}_K^\mu = \epsilon^{\mu\nu\rho\sigma}\kappa_\nu s_{\rho\sigma}$, $\mathcal{J}_P^\mu = \epsilon^{\mu\nu\rho\sigma}\pi_\nu s_{\rho\sigma}$, and $\mathcal{J} = \epsilon^{\mu\nu\rho\sigma}s_{\mu\nu}s_{\rho\sigma}$, with $\epsilon^{\mu\nu\rho\sigma}$ the totally anti-symmetric Levi-Civita tensor ($\epsilon^{0123} = 1$), and $\{A, B\} = AB + BA$. It confirms that all the C's are invariant under (5). If the $\epsilon^{\mu\nu\rho\sigma}$ remains invariant under (6), the C_i's transform as $(C_2, C_3\, C_4) \to (C_2, -C_3, C_4)$. However, the invariance of $\epsilon^{\mu\nu\rho\sigma}$ under (6) is not so trivial, which will be discussed at the end of the next section and afterward.

3 Spin $\frac{1}{2}$

This section deals with the Dirac equation, which describes a spin-$\frac{1}{2}$ particle. In this case, the spin operator $s_{\mu\nu}$, which satisfies (3), can be written using the gamma matrix γ_μ as $s_{\mu\nu} = \mathrm{i}\frac{1}{4}[\gamma_\mu, \gamma_\nu]$, where $\{\gamma_\mu, \gamma_\nu\} = 2g_{\mu\nu}\mathbb{1}$. The next thing is to obtain π_μ from the first equality in (2) and $[\pi_\mu, \pi_\nu] = 0$. Considering that $[\gamma_\rho, s_{\mu\nu}] = \mathrm{i}(g_{\rho\mu}\gamma_\nu - g_{\rho\nu}\gamma_\mu)$, one may suspect that π_μ may be given by $\pi_\mu = \lambda\gamma_\mu$ ($\lambda \in \mathbb{C}$), which, however, would not be appropriate due to $[\pi_\mu, \pi_\nu] \neq 0$. This conclusion is not the end of the story. For an even spacetime dimension, there is a matrix γ_5 such that $\gamma_5^2 = \mathbb{1}$ and $\{\gamma_5, \gamma_\mu\} = 0$. Under the existence of γ_5, the choice of $\pi_\mu = \lambda(\gamma_\mu \pm \gamma_5\gamma_\mu)$ satisfies the first equality in (2) and $[\pi_\mu, \pi_\nu] = 0$. In a similar way, we obtain $\kappa_\mu = \lambda'(\gamma_\mu \pm \gamma_5\gamma_\mu)$ from the second equality in (2) and $[\kappa_\mu, \kappa_\nu] = 0$. The relation between λ and λ', along with the remaining generator Δ, can be derived from (1). To summarize, we have

$$\Delta = \pm\frac{1}{2}\mathrm{i}\gamma_5\,, \quad \pi_\mu = M\left(\frac{\mathbb{1}\pm\gamma_5}{2}\right)\gamma_\mu\,, \quad \kappa_\mu = \frac{1}{M}\left(\frac{\mathbb{1}\mp\gamma_5}{2}\right)\gamma_\mu\,, \quad s_{\mu\nu} = \frac{\mathrm{i}}{4}[\gamma_\mu, \gamma_\nu]\,, \tag{8}$$

where the multiplier $M \in \mathbb{C} \setminus \{0\}$ corresponds to λ in (5). Note that the substitution (6) can be interpreted as $\gamma_5 \to -\gamma_5$. Note also that $[\Delta, s_{\mu\nu}] = 0$.

The fundamental property of π_μ (or κ_μ) is the nilpotence of order two. Let $a_\mu^\pm := (\mathbb{1}\pm\gamma_5)\gamma_\mu$. Then it follows that

$$a_\nu^+ a_\mu^+ = 0 = a_\nu^- a_\mu^-\,. \tag{9}$$

To be more exact, we can show that

$$\begin{cases} a_\mu^+ \mathrm{P}_1 = 0\,, \\ a_\mu^- \mathrm{P}_2 = 0\,, \end{cases} \qquad \begin{cases} a_\mu^+ \mathrm{P}_2 = 2\mathrm{P}_1\gamma_\mu\,, \\ a_\mu^- \mathrm{P}_1 = 2\mathrm{P}_2\gamma_\mu\,, \end{cases} \tag{10}$$

where $\mathrm{P}_1 = \frac{1}{2}(\mathbb{1} + \gamma_5)$ and $\mathrm{P}_2 = \frac{1}{2}(\mathbb{1} - \gamma_5)$ represent the projection operators such that $\mathrm{P}_1 + \mathrm{P}_2 = \mathbb{1}$ and $\mathrm{P}_i\mathrm{P}_j = \delta_{ij}\mathrm{P}_i$. In the Dirac theory, it is well known that P_1 and P_2 are employed in the chiral decomposition. In this sense, (10) can be derived without recognizing the concept of the intrinsic momentum operator π_μ; the existence of π_μ will play a substantial role in higher spin states.

Now we give some properties concerning the Casimir operators C_i's in (7). First, we discuss the invariance of C_3 under (6). Recalling that the substitution (6) corresponds to $\gamma_5 \to -\gamma_5$, and that $\gamma_5 = -\frac{1}{4!}\mathrm{i}\,\epsilon^{\mu\nu\rho\sigma}\gamma_\mu\gamma_\nu\gamma_\rho\gamma_\sigma$, we find that $\gamma_5 \to -\gamma_5$ implies that $\epsilon^{\mu\nu\rho\sigma} \to -\epsilon^{\mu\nu\rho\sigma}$. In this sense, C_3 remains invariant under (6). Next, we obtain the relation between C_2 and C_4. Note that $\mathcal{J}^{\mu\nu}$ can be rewritten as $3\Delta\epsilon^{\mu\nu\rho\sigma}s_{\rho\sigma}$, which leads to $\mathcal{J}_{\mu\nu}\mathcal{J}^{\mu\nu} = 9s_{\mu\nu}s^{\mu\nu}$. In a similar way, we have $\{\mathcal{J}_{K,\mu}, \mathcal{J}_P^\mu\} = -9\{\kappa_\mu, \pi^\mu\}$ and $\frac{1}{16}\mathcal{J}^2 = 9\Delta^2$. Thus we obtain $C_4 = 9C_2$. Anyway, there is no such operator (except a scalar multiple of identity $\mathbb{1}$) that is commutative with all the γ_μ's, so that the C_i's are given by a multiple of identity $\mathbb{1}$ as $(C_2, C_3, C_4) = \frac{15}{4}(1, 2^2, 3^2)\mathbb{1}$.

4 Spin 1

This section deals with relativistically invariant wave equations for spin $s = 1$. For the sake of simplicity, spacetime dimension d is restricted to $(3+1)$. We summarize the wave functions for a free massive particle in Table 1, to find that the π_μ is allowed for the KDP equation but not for the Proca and the WSG equations. This is because the $n \times n$ matrix π_μ such that $[\pi_\rho, s_{\mu\nu}] = \mathrm{i}(g_{\rho\mu}\pi_\nu - g_{\rho\nu}\pi_\mu)$ is allowed for $n = 10$, but not for $n = 4, 6$. In what follows, we concentrate on the KDP equation, where the β_μ's satisfy the trilinear relations

$$\beta_\mu\beta_\nu\beta_\rho + \beta_\rho\beta_\nu\beta_\mu = g_{\mu\nu}\beta_\rho + g_{\rho\nu}\beta_\mu \qquad (\mu, \nu, \rho \in \{0, 1, 2, 3\})\,. \tag{11}$$

Note that β_i ($i = 1, 2, 3$) can be identified with the non-relativistic spin-1 operator s_i in the sense that the s_i's satisfy $s_is_js_k + s_ks_js_i = \delta_{ij}s_k + \delta_{kj}s_i$.

For $n = 10$, it is known that [2] there is a matrix ω ($= \beta_5$) which is given by extending (11) to those for $\mu, \nu, \rho \in \{0, 1, 2, 3, 5\}$ with $g_{5\mu} = g_{\mu5} = \delta_{5\mu}$. Explicitly, we have

$$\omega^3 = \omega\,, \qquad \begin{cases} \{\omega^2, \beta_\mu\} = \beta_\mu\,, \\ \omega\beta_\mu\omega = 0\,, \end{cases} \qquad \begin{cases} \beta_\mu\omega\beta_\nu + \beta_\nu\omega\beta_\mu = 0\,, \\ \omega\beta_\mu\beta_\nu + \beta_\nu\beta_\mu\omega = g_{\mu\nu}\omega\,. \end{cases} \tag{12}$$

Table 1: Lorentz invariant wave equations for $s = 1$ and $d = 3+1$. For the Proca equation, the upperscript in $\psi = (A^0, A^1, A^2, A^3)$ represents the Lorentz vector component, and $\Lambda_{\mu\nu}$ represents the generator of the Lorentz transformation. For the WSG equation, s_i $(i = 1,2,3)$ is given by the (3×3) representation matrix for the non-relativistic spin-1 operator.

Name	Equation	Degree of ψ	$s_{\mu\nu}$	π_μ
Proca	$(\Box + m^2)A^\mu = \partial^\mu(\partial\cdot A)$	4	$\Lambda_{\mu\nu}$	NA
WSG [6,7]	$(\Box + \gamma_{\mu\nu}\partial^\mu\partial^\nu)\psi = 2m_0^2\psi$	6	$\begin{cases} s_{0i} = \frac{1}{\mathrm{i}}\sigma_3\otimes s_i \\ s_{ij} = \mathbb{1}\otimes\epsilon_{ijk}s_k \end{cases}$	NA
KDP [2,8,9]	$(\mathrm{i}\beta_\mu\partial^\mu + m)\psi = 0$	10	$\mathrm{i}[\beta_\mu, \beta_\nu]$	✓

Then the intrinsic conformal generators are given by

$$\Delta = \pm\mathrm{i}\omega\,, \quad \pi_\mu = M\left(\beta_\mu \pm [\omega, \beta_\mu]\right), \quad \kappa_\mu = \frac{1}{M}\left(\beta_\mu \mp [\omega, \beta_\mu]\right), \quad s_{\mu\nu} = \mathrm{i}[\beta_\mu, \beta_\nu]\,. \tag{13}$$

Note that (13) reduces to (8) under $(\beta_\mu, \omega) \to \frac{1}{2}(\gamma_\mu, \gamma_5)$. It is not so difficult to obtain from (11) and (12) the nilpotence of π_μ as

$$\alpha^+_\mu\alpha^+_\nu\alpha^+_\rho = 0 = \alpha^-_\mu\alpha^-_\nu\alpha^-_\rho\,, \tag{14}$$

where $\alpha^\pm_\mu := \beta_\mu \pm [\omega, \beta_\mu]$. To be more exact, we have the following relations:

$$\begin{cases} \alpha^+_\mu \mathsf{P}_1 = 0\,, \\ \alpha^-_\mu \mathsf{P}_3 = 0\,, \end{cases} \quad \begin{cases} \alpha^+_\mu \mathsf{P}_2 = 2\mathsf{P}_1\beta_\mu\,, \\ \alpha^-_\mu \mathsf{P}_2 = 2\mathsf{P}_3\beta_\mu\,, \end{cases} \quad \begin{cases} \alpha^+_\nu\alpha^+_\mu \mathsf{P}_3 = 2\mathsf{P}_1 A_{\mu\nu}\,, \\ \alpha^-_\nu\alpha^-_\mu \mathsf{P}_1 = 2\mathsf{P}_3 A_{\mu\nu}\,, \end{cases} \tag{15}$$

where $A_{\mu\nu} = \{\beta_\mu, \beta_\nu\} - g_{\mu\nu}\mathbb{1}$, and P_i represents a projection operators as $\mathsf{P}_1 = \frac{1}{2}\omega(\omega + \mathbb{1})$, $\mathsf{P}_2 = \mathbb{1} - \omega^2$, and $\mathsf{P}_3 = \frac{1}{2}\omega(\omega - \mathbb{1})$, so that $\sum_{i=1}^3 \mathsf{P}_i = \mathbb{1}$ and $\mathsf{P}_i\mathsf{P}_j = \delta_{ij}\mathsf{P}_i$. Notice that in (15), the lower relations can derive from the corresponding upper ones through the substitution $\omega \to -\omega$. Notice further that $A_{\mu\nu}$ anticommutes with ω, that is

$$\{A_{\mu\nu}, \omega\} = 0\,. \tag{16}$$

The relation (16) leads to $[A^\mu_\mu, \omega^2] = 0$. Note that A^μ_μ and ω are Lorentz invariant in the sense that $[s_{\alpha\beta}, A^\mu_\mu] = 0 = [s_{\alpha\beta}, \omega]$. This relation implies that A^μ_μ can be written as $A^\mu_\mu = \sum_{i=0}^2 c_i\omega^i$ $(c_i \in \mathbb{C})$, where c_i $(i \geq 3)$ is not necessary due to $\omega^3 = \omega$. Here we have assumed that there is no Lorentz invariant other than $\mathbb{1}, \omega$, and ω^2. In this case, we find that $c_0 + c_2 = 0 = c_1$ from $\{A^\mu_\mu, \omega\} = 0$ by (16), and that $c_0 = 2$ from $\{\beta_\nu, \beta_\mu\beta^\mu\} = 5\beta_\nu$ by (11) and $\{\beta_\nu, \omega^2\} = \beta_\nu$ by (12). Eventually, we have

$$\beta_\mu\beta^\mu = \mathsf{P}_2 + 2\mathbb{1}\,. \tag{17}$$

Actually, the relation (17) holds in the ten-dimensional representation [2] for (11) and (12), which corresponds to the adjoint representation of the Lorentz group in five dimensions (for the adjoint representation, we have $\binom{5}{2} = 10$ Lorentz group generators). For later convenience, we rewrite $\frac{1}{2}s_{\mu\nu}s^{\mu\nu}$ using P_2 as

$$\frac{1}{2}s_{\mu\nu}s^{\mu\nu} = 4\mathbb{1} - \mathsf{P}_2\,, \tag{18}$$

where we have used (17), together with $\mathsf{P}_2^2 = \mathsf{P}_2$.

As was mentioned in Sec. 1, the π_μ should annihilate the physical state. To check the validity, we show that the rank of P_k (or equivalently, the trace of P_k) for $k=1,3$ equals the spin degrees of freedom. In the ten-dimensional representation, the eigenvalues of ω are given by $1,0,-1$ appearing $3,4,3$ times, respectively. Thus, we obtain

$$\mathrm{Rank}(\mathsf{P}_1)=\mathrm{Rank}(\mathsf{P}_3)=3\,,\qquad \mathrm{Rank}(\mathsf{P}_2)=4\,.$$

This result is quite reasonable because the number "3" equals the spin degree of freedom for a massive particle for $s=1$. To confirm the validity, we calculate the 3-dimensional spin magnitude $\langle s\rangle^2:=s_{12}{}^2+s_{23}{}^2+s_{31}{}^2$. Let $|\psi^+_{\mathrm{ph}}\rangle=\mathsf{P}_1|\psi\rangle$, $|\psi^-_{\mathrm{ph}}\rangle=\mathsf{P}_3|\psi\rangle$, and $|\psi_{\mathrm{un}}\rangle=\mathsf{P}_2|\psi\rangle$, in which we have $\alpha^\pm_\mu|\psi^\pm_{\mathrm{ph}}\rangle=0$. Recalling that $\langle s\rangle^2$ $(=\frac{1}{4}s_{\mu\nu}s^{\mu\nu})=2\mathbb{1}-\frac{1}{2}\mathsf{P}_2$ by (18), and that $\mathsf{P}_i\mathsf{P}_j=\delta_{ij}\mathsf{P}_i$, we obtain $\langle s\rangle^2|\psi^\pm_{\mathrm{ph}}\rangle=s(s+1)|\psi^\pm_{\mathrm{ph}}\rangle$ $(s=1)$ and $\langle s\rangle^2|\psi_{\mathrm{un}}\rangle=\frac{3}{2}|\psi_{\mathrm{un}}\rangle$. These relations indicate that $|\psi^\pm_{\mathrm{ph}}\rangle$ represents the spin-1 state, while $|\psi_{\mathrm{un}}\rangle$ does not. Bearing these findings in mind, we can regard $|\psi^\pm_{\mathrm{ph}}\rangle$ and $|\psi_{\mathrm{un}}\rangle$ as physical and unphysical states, respectively.

Finally, we give some properties of the Casimir operator C. As in the case of $s=\frac{1}{2}$, the invariance of C_3 under (6) is guaranteed by the statement that $(\omega\to-\omega)\Longrightarrow(\epsilon^{\mu\nu\rho\sigma}\to-\epsilon^{\mu\nu\rho\sigma})$ by $\omega=-\frac{\mathrm{i}}{4}\epsilon^{\mu\nu\rho\sigma}\beta_\mu\beta_\nu\beta_\rho\beta_\sigma$ [10, 11]. After a somewhat tedious calculation, we can write the C_i's in (7) as $(C_2,C_3,C_4)=(9,48,144)\mathbb{1}$, which confirms the irreducibility of the ten-dimensional representation.

5 Spin $\frac{3}{2}$

In this section, we consider the $(3+1)$-dimensional Minkowski space, as in the case of $s=1$. Although the Rarita-Schwinger equation is well known as a relativistic invariant wave equation for $s=\frac{3}{2}$, the intrinsic momentum operator is not allowed, as in the case of the Proca equation. Instead, we adopt a Dirac-like wave equation for parafermion of order 3, namely (massive) Bhabha wave equation [3] (see Table 2).

Extending the polynomial relations among the non-relativistic spin operators s_i's $(i=1,2,3)$ to those among s_μ's $(\mu=0,1,2,3)$ in a relativistically covariant way, we obtain

$$\begin{cases} s_\mu s_\nu s_\alpha+s_\alpha s_\nu s_\mu+g_{\mu\alpha}s_\nu=s_\mu s_\alpha s_\nu+s_\nu s_\alpha s_\mu+g_{\mu\nu}s_\alpha\,,\\ 0=\left(s_\mu s_\nu s_\alpha s_\beta-\frac{5}{4}\{s_\mu,s_\nu\}g_{\alpha\beta}+\frac{9}{16}g_{\mu\nu}g_{\alpha\beta}\right)+(\text{perm. of }\mu,\nu,\alpha,\beta)\,. \end{cases} \tag{19}$$

It may be convenient to rewrite the first relation of (19) as $[s_\mu,[s_\nu,s_\alpha]]=g_{\mu\nu}s_\alpha-g_{\mu\alpha}s_\nu$. Note that $\frac{1}{2}\gamma_\mu$ satisfies both relations in (19). This implies that there should exist a polynomial relation such that $p(s_0,s_1,s_2,s_3)=0$ with $p(s_0,s_1,s_2,s_3)|_{s_\mu\to\frac{1}{2}\gamma_\mu}\neq0$. However, we neglect, for the time being, such a polynomial relation because it is not irrelevant to the following discussion. Suppose that there exists an operator s_5 which satisfies (19) for $\mu,\nu,\alpha,\beta\in\{0,1,2,3,5\}$, with

Table 2: Lorentz invariant wave equations for $s=\frac{3}{2}$. For the Rarita equation, ψ is composed of four Dirac spinors as $\psi:=(\psi_0,\psi_1,\psi_2,\psi_3)$, where the subscript represents the Lorentz vector component, so that Λ $(=\{\Lambda_{\mu\nu}\}):\psi\mapsto\psi'$ acts as $(\psi')_\mu=\Lambda_\mu^{\ \nu}\psi_\nu$.

Name	Equation	Degree of ψ	$s_{\mu\nu}$	π_μ
Rarita-Schwinger	$(\epsilon^{\mu\nu\rho\sigma}\gamma_5\gamma_\nu\partial_\rho+mg^{\mu\sigma})\psi_\sigma=0$	4×4	$\Lambda_{\mu\nu}+\frac{\mathrm{i}}{4}[\gamma_\mu,\gamma_\nu]$	NA
Bhabha	$(\mathrm{i}s_\mu\partial^\mu+m)\psi=0$	20	$\mathrm{i}[s_\mu,s_\nu]$	✓

$g_{5\mu} = g_{\mu 5} = \delta_{5\mu}$. Then the intrinsic conformal generators are given, as is analogous to the case of $s = \frac{1}{2}, 1$, by

$$\Delta = \pm \mathrm{i}s_5\,, \quad \pi_\mu = M\left(s_\mu \pm [s_5, s_\mu]\right)\,, \quad \kappa_\mu = \frac{1}{M}\left(s_\mu \mp [s_5, s_\mu]\right)\,, \quad s_{\mu\nu} = \mathrm{i}[s_\mu, s_\nu]\,. \tag{20}$$

Note that the first equality in (19), together with the existence of s_5, is sufficient for (20); the second equality in (19) is not necessary for (20). Recalling that the first relation in (19) is satisfied for $s_\mu \to \frac{1}{2}\gamma_\mu$ ($s = \frac{1}{2}$) and for $s_\mu \to \beta_\mu$ ($s = 1$), we find it natural that the relation (20) is the same form as (8) and (13). For later convenience, we obtain some operators which anti-commute with s_5. Such operators are exemplified as

$$\{s_5, A_\mu\} = 0 = \{s_5, A_{\rho\nu\mu} + (\text{perm. of } \rho, \nu, \mu)\}\,, \tag{21}$$

where $A_\mu = s_5 s_\mu s_5 - \frac{3}{4}s_\mu$, and $A_{\rho\nu\mu} = s_\rho s_\nu s_\mu - \frac{7}{4}g_{\rho\nu}s_\mu$.

The projection operators P_i's ($i = 1, 2, 3, 4$) can be written using the minimum polynomial $f(x)$ with respect to s_5 as $\mathsf{P}_i = \frac{1}{f'(\lambda_i)}\frac{f(s_5)\mathbb{1}}{s_5 - \lambda_i \mathbb{1}}$, where $f(x) = \prod_{i=1}^4 (x - \lambda_i)$, with $\lambda_1 = \frac{3}{2}$, $\lambda_2 = \frac{1}{2}$, $\lambda_3 = -\frac{1}{2}$, $\lambda_4 = -\frac{3}{2}$. Let $s^\pm_\mu := s_\mu \pm [s_5, s_\mu]$. Then it follows that (see Appendix A)

$$\begin{cases} s^+_\mu \mathsf{P}_1 = 0\,, \\ s^-_\mu \mathsf{P}_4 = 0\,, \end{cases} \quad \begin{cases} s^+_\mu \mathsf{P}_2 = 2\mathsf{P}_1 X_\mu\,, \\ s^-_\mu \mathsf{P}_3 = 2\mathsf{P}_4 X_\mu\,, \end{cases} \quad \begin{cases} s^+_\nu s^+_\mu \mathsf{P}_3 = 2\mathsf{P}_1 X_{\nu\mu}\,, \\ s^-_\nu s^-_\mu \mathsf{P}_2 = 2\mathsf{P}_4 X_{\nu\mu}\,, \end{cases} \quad \begin{cases} s^+_\rho s^+_\nu s^+_\mu \mathsf{P}_4 = \frac{4}{3}\mathsf{P}_1 X_{\rho\nu\mu}\,, \\ s^-_\rho s^-_\nu s^-_\mu \mathsf{P}_1 = \frac{4}{3}\mathsf{P}_4 X_{\rho\nu\mu}\,, \end{cases} \tag{22}$$

where X_μ, $X_{\nu\mu}$ and $X_{\rho\nu\mu}$ are given by

$$X_\mu = s_\mu\,, \quad X_{\nu\mu} = \{s_\nu, s_\mu\} - s g_{\nu\mu}\mathbb{1}\,, \quad X_{\rho\nu\mu} = \left[Y_{\rho\nu\mu} + (\text{perm. of } \rho, \nu, \mu)\right]\,,$$

with $s = \frac{3}{2}$ and $Y_{\rho\nu\mu} := s_\rho s_\nu s_\mu - g_{\rho\nu}(s s_\mu + \frac{1}{2s}s_5 s_\mu s_5) \to A_{\rho\nu\mu} - \frac{1}{3}g_{\rho\nu}A_\mu$ $(s = \frac{3}{2})$. The relations (22) lead to $s^+_\mu s^+_\nu s^+_\rho s^+_\sigma \mathsf{P}_i = 0 = s^-_\mu s^-_\nu s^-_\rho s^-_\sigma \mathsf{P}_i$ ($i = 1, 2, 3, 4$), from which, together wirh $\sum_{i=1}^4 \mathsf{P}_i = \mathbb{1}$, we obtain the nilpotence of $s^\pm_\mu$ (of order 4) as

$$s^+_\mu s^+_\nu s^+_\rho s^+_\sigma = 0 = s^-_\mu s^-_\nu s^-_\rho s^-_\sigma\,. \tag{23}$$

Note that by (21), not only have we the anti-commutativity

$$\{X_{\rho\nu\mu}, s_5\} = 0\,,$$

but also the anti-commutativities $\{\gamma_\mu, \gamma_5\} = 0$ and (16) can be rewritten using X_μ and $X_{\nu\mu}$ as

$$\{X^{(\frac{1}{2})}_\mu, \gamma_5\} = 0 = \{X^{(1)}_{\nu\mu}, \omega\}\,, \tag{24}$$

where $X^{(\frac{1}{2})}_\mu$ and $X^{(1)}_{\nu\mu}$, more generally, $X^{(s)}_{\nu\mu\ldots}$ represents the corresponding $X_{\nu\mu\ldots}$ for a given spin s. For example, we have $Y^{(\frac{1}{2})}_{\rho\nu\mu} = \frac{1}{8}\gamma_\rho\gamma_\nu\gamma_\mu - \frac{1}{8}g_{\rho\nu}\gamma_\mu$, and $Y^{(1)}_{\rho\nu\mu} = \beta_\rho\beta_\nu\beta_\mu - g_{\rho\nu}\beta_\mu$ by replacing $(s_\rho, s_\nu, s_\mu; s)$ in $Y_{\rho\nu\mu}$ with $\frac{1}{2}(\gamma_\rho, \gamma_\nu, \gamma_\mu; 1)$ and $(\beta_\rho, \beta_\nu, \beta_\mu; 1)$, respectively. Note further that we have the following vanishing relations:

$$X^{(\frac{1}{2})}_{\nu\mu} = X^{(\frac{1}{2})}_{\rho\nu\mu} = 0\,, \qquad X^{(1)}_{\rho\nu\mu} = 0\,,$$

which, in vew of (22), are due to the relations (9) and (14), respectively.

Now we discuss whether or not physical states can be given by $\mathsf{P}_k|\psi\rangle$ ($k = 1, 4$) by calculating the rank of P_k. In the Bhabha theory [3] for $s = \frac{3}{2}$, we have two irreducible representations $\mathrm{R}_5(\frac{3}{2}, \frac{3}{2})$ and $\mathrm{R}_5(\frac{3}{2}, \frac{1}{2})$, where $\mathrm{R}_5(s, \tilde{s})$ represents the spin-s Lorentz group representation in five dimensions. Let $S := \{s_1, s_2, s_3, \mathrm{i}s_0\}$. For $\mathrm{R}_5(\frac{3}{2}, \frac{3}{2})$, the eigenvalues of $x \in S$

are $\frac{3}{2}, \frac{1}{2}, -\frac{1}{2}, -\frac{3}{2}$ appearing $4, 6, 6, 4$ times, respectively; while for $\mathrm{R}_5(\frac{3}{2}, \frac{1}{2})$, the eigenvalues of $x \in S$ are $\frac{3}{2}, \frac{1}{2}, -\frac{1}{2}, -\frac{3}{2}$ appearing $2, 6, 6, 2$ times, respectively. If s_5 realizes, the eigenvalues of s_5 are identical with those of $x \in S$, so that

$$\mathrm{Rank}(\mathsf{P}_1) = \mathrm{Rank}(\mathsf{P}_4) = \begin{cases} 4 & \left(\mathrm{R}_5\left(\frac{3}{2}, \frac{3}{2}\right)\right), \\ 2 & \left(\mathrm{R}_5\left(\frac{3}{2}, \frac{1}{2}\right)\right), \end{cases} \qquad \mathrm{Rank}(\mathsf{P}_2) = \mathrm{Rank}(\mathsf{P}_3) = \begin{cases} 6 & \left(\mathrm{R}_5\left(\frac{3}{2}, \frac{3}{2}\right)\right), \\ 6 & \left(\mathrm{R}_5\left(\frac{3}{2}, \frac{1}{2}\right)\right). \end{cases}$$

Thus we obtain in the representation $\mathrm{R}_5(\frac{3}{2}, \frac{3}{2})$, the relation $\mathrm{Rank}(\mathsf{P}_1) = \mathrm{Rank}(\mathsf{P}_4) = 4$, the spin degrees of freedom for a spin-$\frac{3}{2}$ massive particle.

The analogous relation holds for a general spin s. Note that by a fundamental property of the projector, we have $\mathrm{Rank}(\mathsf{P}_i) = N_i$, where N_i represents the number of the eigenvalue $(s+1-i)$ of s_5. Note also that in the representaion $\mathrm{R}_5(s, \tilde{s})$ ($\tilde{s} = s, s-1, \ldots$), the maximum and minimum eigenvalues of s_5 [that is, s and $(-s)$, respectively] occur $(2\tilde{s}+1)$ times [3]. Considering these two remarks, we obtain in the representation $\mathrm{R}_5(s, s)$, the relation $\mathrm{Rank}(\mathsf{P}_1) = \mathrm{Rank}(\mathsf{P}_{2s+1}) = 2s+1$, the spin degrees of freedom. To confirm that $|\psi^+_{\mathrm{ph}}\rangle = \mathsf{P}_1|\psi\rangle$ and $|\psi^-_{\mathrm{ph}}\rangle = \mathsf{P}_{2s+1}|\psi\rangle$, in which we have $s^{\pm}_{\mu}|\psi^{\pm}_{\mathrm{ph}}\rangle = 0$, can be regarded as physical states, we should further show $\langle s\rangle^2|\psi^{\pm}_{\mathrm{ph}}\rangle = s(s+1)|\psi^{\pm}_{\mathrm{ph}}\rangle$, which, however, will be discussed elsewhere.

6 Conclusion

We have found that the intrinsic momentum operator $\pi_\mu = s^+_\mu, s^-_\mu$, which we do not introduce in the ordinary conformal group, is feasible for the Bhabha wave equation, provided that s_5, corresponding to $\frac{1}{2}\gamma_5$ ($s = \frac{1}{2}$) and ω ($s = 1$), exists. For a general spin s, we can write the intrinsic conformal generators as the same relations as (20) and those where $s_5 \to (-s_5)$, satisfying the invariance under (5) and (6). The fundamental property of π_μ is the nilpotence of order $(2s+1)$. To be more exact, let P_i's ($i = 1, 2, \ldots, 2s+1$) be the projection operators concerning the s_5 as $\mathsf{P}_i = \frac{1}{f'(\lambda_i)} \frac{f(s_5)\mathbb{1}}{s_5 - \lambda_i \mathbb{1}}$, where $f(x) = \prod_{i=1}^{2s+1}(x - \lambda_i)$, $\lambda_i = s + 1 - i$. Then we have the same hierarchical relation as (22), where $X^{(\frac{1}{2})}_{\mu}$, $X^{(1)}_{\mu\nu}, \ldots$ anti-commute with γ_5, $\omega, \ldots$, respectively. As long as the wave function transforms as a scalar under the spacetime translation, either s^+_μ or s^-_μ should annihilate a physical state, so that the relation $\mathrm{Rank}(\mathsf{P}_k) = 2s+1$ ($k = 1, 2s+1$) is required for a massive particle. Fortunately, this relation holds in the representation $\mathrm{R}_5(s, s)$, irreducible representation of the Lorentz group in five dimensions.

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A Derivation of (22)

It is not so difficult to obtain X_μ and $X_{\nu\mu}$ by rewriting $s^+_\mu \mathsf{P}_2$ and $s^+_\nu s^+_\mu \mathsf{P}_3$ in such a way that s_5 is located as leftward as possible. However, this procedure is not practical for the calculation of $X_{\rho\nu\mu}$ because $X_{\rho\nu\mu}$ hinges on s_5 so that we may not represent $X_{\rho\nu\mu}$ uniquely due to some relations between s_5 and s_μ's. In this sense, it would be better to adopt another approach. We start with the following relation:

$$s_\mu^+ \mathsf{P}_4 = 2X_\mu \mathsf{P}_4 \qquad (X_\mu = s_\mu). \tag{A.1}$$

Keeping the form of (A.1) without rearranging s_5 leftward, and applying s_ν^+ to both sides of (A.1) from the left, then we find it rather simple to obtain

$$s_\nu^+ s_\mu^+ \mathsf{P}_4 = 2X_{\nu\mu} \mathsf{P}_4 \quad \left(X_{\nu\mu} = \{s_\nu, s_\mu\} - s\mathbb{1}, \quad s = \frac{3}{2} \right),$$

where we have used $[s_\nu^+, s_\mu] = [s_\nu, s_\mu] + g_{\nu\mu} s_5$, together with the relation $s_5 \mathsf{P}_4 = -s\mathsf{P}_4$. Further application of s_ρ^+ leads to the relation

$$s_\rho^+ s_\nu^+ s_\mu^+ \mathsf{P}_4 = \tfrac{4}{3} X_{\rho\nu\mu} \mathsf{P}_4 \quad \left(X_{\rho\nu\mu} = Y_{\rho\nu\mu} + (\text{perm. of } \rho, \nu, \mu) \right),$$

where $Y_{\rho\nu\mu} = s_\rho s_\nu s_\mu - g_{\rho\nu}(s s_\mu + \frac{1}{2s} s_5 s_\mu s_5)$. A similar calculation yields $s_\rho^- s_\nu^- s_\mu^- \mathsf{P}_1 = \frac{4}{3} X_{\rho\nu\mu} \mathsf{P}_1$. Recalling that $\{s_5, X_{\rho\nu\mu}\} = 0$ by (21) and noticing that $\mathsf{P}_1 \leftrightarrow \mathsf{P}_4$ under the substitution $s_5 \to -s_5$, we finally get the last relation in (22).

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