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# On spin-$\frac{5}{2}$ gauge fields†

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**Abstract.** A gauge theory for spin-$\frac{5}{2}$ fields is discussed.

## 1. Introduction

The beginning of a new theory for spin-$\frac{5}{2}$ fields is presented below. These results were obtained by Fang and Fronsdal (1978), and independently but later by us. The first authors actually found actions for fields of any spin (integer (Fronsdal 1978) as well as half-integer (Fang and Fronsdal 1978)). We (Berends *et al* 1979a) concentrated only on spin-$\frac{5}{2}$ and studied its quantisation (Berends *et al* 1979b), BRS rules, etc.

The Lagrangian for a free massless spin-$\frac{5}{2}$ field, coupled to an external source, is given by‡

$$\mathscr{L} = -\tfrac{1}{2}\bar{\psi}_{\mu\nu}\not{\partial}\psi_{\mu\nu} + (\bar{\psi}\cdot\gamma)_\mu[2(\partial\cdot\psi)_\mu - \not{\partial}(\gamma\cdot\psi)_\mu - \partial_\mu\psi] + \tfrac{1}{4}\bar{\psi}\not{\partial}\psi + \bar{\psi}_{\mu\nu}T_{\mu\nu} \tag{1}$$

where $\psi_{\mu\nu}$ is a symmetric tensorial Majorana spinor. This action, as one easily shows, is invariant under the following local fermionic gauge transformation:

$$\delta\psi_{\mu\nu} = \partial_\mu\epsilon_\nu + \partial_\nu\epsilon_\mu \qquad \text{with} \qquad \gamma\cdot\epsilon = 0. \tag{2}$$

The remarkable discovery is that the *restriction* $\gamma\cdot\epsilon = 0$ *still leads to only two physical modes*! One can show this either in the field equation, by choosing a suitable gauge (for example, $\gamma_\mu\psi_{\mu\nu} - \frac{1}{4}\gamma_\nu\psi = 0$), or by evaluating the residue of the propagator, taking into account that the source $T_{\mu\nu}$ satisfies some constraints due to the presence of a local gauge invariance in the action. It was widely believed that one needed $\partial_\mu T_{\mu\nu} = 0$ and thus $\delta\psi_{\mu\nu} = \partial_\mu\epsilon_\nu + \partial_\nu\epsilon_\mu$ without $\gamma\cdot\epsilon = 0$, in order to show that all unphysical modes cancel from the theory. For example, Schwinger (1970) obtained the field equation corresponding to equation (1) as well as the source constraint corresponding to equation (2) basing himself on conserved sources. Actually, as we now know, this criterion is too strong and can be weakened as follows. Namely, we can redefine the field $\psi_{\mu\nu}$ by

$$\psi'_{\mu\nu} = \psi_{\mu\nu} + a\delta_{\mu\nu}\psi + b(\gamma_\mu\gamma_\lambda\psi_{\lambda\nu} + \gamma_\nu\gamma_\lambda\psi_{\lambda\mu}) \tag{3}$$

† Based on a talk given at the Sanibel meeting on *Fundamentals of Quantum Theory and Quantum Field Theory* February 1979 in Gainesville, Florida, by P van Nieuwenhuizen.
‡ Conventions: $(\gamma\cdot\psi)_\mu = \gamma_\nu\psi_{\nu\mu}$, $\psi = \psi_{\mu\mu}$, and $\gamma_\mu\gamma_\nu + \gamma_\nu\gamma_\mu = 2\delta_{\mu\nu}$ with $\delta_{\mu\nu} = (+,+,+,+)$. Also $p^2 = \boldsymbol{p}^2 + p_4^2$ and $\not{k} = k_\mu\gamma_\mu$.

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to bring the Lagrangian into a simpler form (see (Berends *et al* 1979a, equation (4.23)):

$$\mathscr{L} = -\tfrac{1}{2}\bar{\psi}_{\mu\nu}\not{\partial}\psi_{\mu\nu} - \tfrac{1}{4}(\bar{\psi}\cdot\gamma)_{\mu}\not{\partial}(\gamma\cdot\psi)_{\mu} + \tfrac{1}{2}\bar{\psi}\partial\cdot\gamma\cdot\psi. \tag{4}$$

In that representation the gauge invariance (2) is more complicated and given by

$$\delta\psi_{\mu\nu} = \tfrac{1}{2}\not{\partial}(\gamma_{\mu}\epsilon_{\nu} + \gamma_{\nu}\epsilon_{\mu}) + \delta_{\mu\nu}(\partial\cdot\epsilon - \tfrac{1}{2}\not{\partial}\gamma\cdot\epsilon) \tag{5}$$

which corresponds to having the source constraint

$$\partial_{\mu}T_{\mu\nu} = \not{\partial}[\tfrac{1}{2}(\gamma\cdot T)_{\nu} - \tfrac{1}{4}\gamma_{\nu}T]. \tag{6}$$

In a series of papers, Weinberg has analysed higher spin theories (Weinberg 1964). First, he shows that Lorentz invariance of the $S$ matrix requires that the source be conserved. Adapted to the spin-$\frac{5}{2}$ case, the argument runs as follows. The matrix element for emission of one spin-$\frac{5}{2}$ particle (with, say, helicity $+\frac{5}{2}$) is

$$S = \bar{u}^{+}(k)\epsilon^{+}_{\mu}(k)\epsilon^{+}_{\nu}(k)M_{\mu\nu}. \tag{7}$$

Under a boost, $\epsilon^{+}_{\mu}$ acquires a term proportional to $k_{\mu}$ and it seems that one needs $k_{\mu}M_{\mu\nu} = 0$ in order that the $S$-matrix element be Lorentz-invariant. Actually, we see that, since the spinor $u^{+}(k)$ satisfies the Dirac equation, the weak conservation in (6) is enough. The $S$ matrix is still on-shell gauge-invariant.

Subsequently, Weinberg derives restrictions on the possible couplings of massless particles with spin $J$ by considering the emission of one soft spin $J$ particle from a given process. For soft photons and gravitons, the matrix element is dominated by the poles due to emission from incoming or outgoing physical particles:

$$S\,(\text{photon}) = \sum e_i\left(\frac{P_i\cdot\epsilon}{P_i\cdot k}\right)S\,(\text{no photon})$$

$$S\,(\text{graviton}) = \sum g_i\frac{(P_i\cdot\epsilon)^2}{P_i\cdot k}S\,(\text{no graviton}).$$

Gauge invariance then requires charge conservation ($\Sigma e_i = 0$) and the equivalence principle ($g_i = \kappa$). For spin $\frac{3}{2}$, Grisaru and Pendleton (1977) applied these arguments to supergravity and found the same relations between processes with different numbers of bosons and fermions as follow from the global supersymmetry algebra. For spin $\frac{5}{2}$ it implies (Grisaru and Pendleton 1977, Grisaru *et al* 1977) that there are no long-range forces (gradient couplings).

## 2. Two modes in the field equations

To show that there are only two modes in the classical theory, we consider the field equation obtained from (1):

$$O_{\mu\nu} = \not{\partial}\psi_{\mu\nu} - \partial_{\mu}(\gamma\cdot\psi)_{\nu} - \partial_{\nu}(\gamma\cdot\psi)_{\mu} = 0. \tag{8}$$

(Note that the action is not simply $-\frac{1}{2}\bar{\psi}_{\mu\nu}O_{\mu\nu}$ as this would give a different field equation. In general, if $\bar{\psi}_{\mu\nu}F_{\mu\nu}$ is symmetric in both spinors, it is a good action. The converse is thus not true.) Choosing the relativistically invariant gauge†

$$\gamma_\mu \psi_{\mu\nu} = 0 \text{ (gauge condition)} \tag{9}$$

one finds that one can obtain this gauge, since $\delta(\gamma \cdot \psi)_\mu = \not{\partial}\epsilon_\mu$ if $\gamma \cdot \epsilon = 0$. Thus, $\epsilon_\mu = \not{\partial}^{-1}(\gamma \cdot \psi)_\mu$ and indeed

$$\Box \gamma \cdot \epsilon = -\not{\partial}\psi + 2\partial_\mu(\gamma \cdot \psi)_\mu = 0 \tag{10}$$

as one sees by tracing (8). Thus, on-shell $(\gamma \cdot \psi)_\mu = 0$, hence $\psi = 0$, and (8) yields

$$\partial_\mu \psi_{\mu\nu} = \not{\partial}\psi_{\mu\nu} = 0 \qquad \text{(on-shell in gauge } \gamma_\mu\psi_{\mu\nu} = 0\text{)}. \tag{11}$$

Hence, in this gauge, $k^2 = 0$ on-shell.

Expanding the tensor part of $\psi_{\mu\nu}$ into the basis with $\epsilon_\mu^+$, $\epsilon_\mu^-$, $k_\mu$ and $\bar{k}_\mu$ where $\bar{k}_\mu$ is the time-reversed $k$ (this approach was first applied to spin $\frac{3}{2}$, see Sterman *et al* (1978)), one has with $k_\mu \psi_{\mu\nu} = 0$ on-shell

$$\psi_{\mu\nu} = \sum_{\lambda,\lambda' = +,-,0} \epsilon_\mu^\lambda \epsilon_\nu^{\lambda'} G_{\lambda\lambda'} (\epsilon_\mu^0 \equiv k_\mu). \tag{12}$$

From $\not{k}\psi_{\mu\nu} = 0$ it follows that

$$\not{k} G_{\lambda\lambda'} = 0 \qquad \text{for} \qquad \lambda, \lambda' = +, -, 0. \tag{13}$$

From $\psi = 0$ and $(k_\mu)^2 = (\epsilon_\mu^+)^2 = 0$ it follows that $G_{+-} = 0$. From $\gamma_\mu \psi_{\mu\nu} = 0$ one then has

$$\not{\epsilon}^+ G^{++} = \not{\epsilon} G^{--} = \not{\epsilon}^+ G^{+0} + \not{\epsilon}^- G^{-0} = 0. \tag{14}$$

Choose now a second gauge with $\epsilon_\mu = -\epsilon_\mu^+ G^{+0} - \frac{1}{2}k_\mu G^{00} - \epsilon_\mu^- G^{-0}$. Then $\hat{\psi}_{\mu\nu} = \psi_{\mu\nu} + \partial_\mu \epsilon_\nu + \partial_\nu \epsilon_\mu$ still satisfies

$$\not{\partial}\hat{\psi}_{\mu\nu} = \gamma_\mu \hat{\psi}_{\mu\nu} = k_\mu \hat{\psi}_{\mu\nu} = 0. \tag{15}$$

Hence $\hat{\psi}_{\mu\nu} = \epsilon_\mu^+ \epsilon_\mu^+ G^{++} + \epsilon_\mu^- \epsilon_\nu^- G^{--}$. Finally, from $\not{\epsilon}^+ G^{++} = 0$, it follows that the spinor $(G^{++})$ has helicity $+\frac{1}{2}$. To see this, note that

$$\not{\epsilon}^- \not{\epsilon}^+ = 1 - \frac{1}{2} \frac{\boldsymbol{\sigma} \cdot \boldsymbol{k}}{|\boldsymbol{k}|}. \tag{16}$$

Thus,

$$\psi_{\mu\nu} = A u^+(\boldsymbol{k}) \epsilon_\mu^+(\boldsymbol{k}) \epsilon_\nu^+(\boldsymbol{k}) + B u^-(\boldsymbol{k}) \epsilon_\mu^-(\boldsymbol{k}) \epsilon_\nu^-(\boldsymbol{k}) \tag{17}$$

and there are indeed two modes in the classical theory.

One can use gauges which hold also off-shell. An example is the gauge in equation (29). One can also employ the non-relativistic gauge $\gamma_i \psi_{ij} = 0$ (summed over $i = 1, 2, 3$), which is the analogue of the Coulomb gauge (Berends *et al* 1979b). The results remain unchanged.

## 3. Two modes in the propagator

For spins below $\frac{5}{2}$, absence of ghosts in the field equations implies absence of ghosts in the propagators and vice versa. Thus it comes as no surprise that in the spin-$\frac{5}{2}$ massless propagator there are also no ghosts. In fact, the action and transformation rules in

† In the following we choose the simplest gauge, $\gamma_\mu \psi_{\mu\nu} = 0$, although this can only hold on-shell. One can repeat our analysis in the gauge $\gamma_\mu \psi_{\mu\nu} - \frac{1}{4}\gamma_\nu \psi = 0$ which can also be satisfied off-shell.

equations (1) and (2) were obtained by Berends *et al* (1979a) by the very requirement of absence of ghosts in the propagator. We now briefly sketch the procedure; for more details see Berends *et al* (1979a).

The spin content of the tensor-spinor $\psi_{\mu\nu}$ is readily found by adding angular momenta. Since a graviton field has one spin-2, one spin-1 and two spin-0 components (of which the spin-1 and one of the spin-0 are absent in the Einstein action), it follows that

$$\{\psi_{\mu\nu}\} = \tfrac{5}{2}\oplus\tfrac{3}{2}\oplus\tfrac{3}{2}\oplus\tfrac{1}{2}\oplus\tfrac{1}{2}\oplus\tfrac{1}{2}. \tag{18}$$

For each spin one introduces a set of projection operators $P^J_{ii}$ where $ii$ distinguishes between components with the same spin (thus $ii = 11, 22$ for $J=\frac{3}{2}$ and $ii = 11, 22, 33$ for $J=\frac{1}{2}$). In addition, one introduces transition operators $P^J_{ij} (i \neq j)$ between different components of equal spin. The sum of the projection operators equals unity:

$$(P^{5/2} + P^{3/2}_{11} + P^{3/2}_{22} + P^{1/2}_{11} + P^{1/2}_{22} + P^{1/2}_{33})^{\alpha\beta}_{\mu\nu\rho\sigma} = \tfrac{1}{2}(\delta_{\mu\rho}\delta_{\nu\sigma} + \delta_{\mu\sigma}\delta_{\nu\rho})\delta^{\alpha\beta} \tag{19}$$

and the whole set of operators is orthonormal:

$$P^J_{ij}P^K_{kl} = \delta_{jk}\delta^{JK}P^J_{il}. \tag{20}$$

The most general field equation as well as the most general propagator for $\psi_{\mu\nu}$ can be expanded on this complete basis. The advantage of using this spin projection formalism is that the whole analysis is reduced to an analysis within separate and independent spin blocks. For a discussion of this spin formalism for spin-2 fields, see van Nieuwenhuizen (1973).

In Berends *et al* (1979a) the most general propagator was written down and was required to have positive residues. This put conditions on the external sources with which this propagator was sandwiched. Requiring moreover that there be only two propagating modes with helicity $\frac{5}{2}$ led to additional source constraints. These conditions were formulated in terms of spin projection operators:

$$(\alpha_j P^J_{ij})_{\mu\nu,\rho\sigma}T_{\rho\sigma} = 0 \text{ (fixed } J). \tag{21}$$

In this way we found both in the spin $J=\frac{3}{2}$ sector and in the spin $J=\frac{1}{2}$ sector one source constraint. Since, as is well known, source constraints imply gauge invariances of the field equation and vice versa, the constraints on the external spin-$\frac{5}{2}$ sources led to gauge invariances of the field equations. The establishment of the one-to-one correspondence between source constraints and gauge invariances uses the spin projection operators (van Nieuwenhuizen 1973). In this way an action was written down, again in terms of spin projection operators, which had these gauge invariances. The crucial point now was whether this action was local, i.e. whether it was linear and homogeneous in derivatives. Fortunately, it was.

## 4. Matter couplings

If one follows the successful path which led to supergravity, one replaces $\partial_\mu$ by $D_\mu$ in (1) and varies with

$$\delta\psi_{\mu\nu} = D_\mu\epsilon_\nu + D_\nu\epsilon_\mu \qquad \gamma\cdot\epsilon = 0. \tag{22}$$

If one ends up with the Einstein tensor times something:

$$\delta I(\tfrac{5}{2}) = G_{\mu\nu}\bar{\epsilon}_\rho\Sigma_{\mu\nu\rho} \tag{23}$$

then one can add the Einstein action and obtain full invariance at the local coupled level by $\delta e_{a\mu} \sim \bar{\epsilon}_\rho \Sigma_{a\mu\rho}$. However, one finds a Riemann tensor. Thus it seems one cannot couple spin-2 and -$\frac{5}{2}$. Similarly, the coupling to the electromagnetic field is inconsistent (Berends *et al* 1979b).

If one starts at the matter end, and tries to find globally invariant action (invariant under transformations with constant $\epsilon_\mu$, still with $\gamma . \epsilon = 0$) one notices

(i) that, as always, cancellation between fields occurs pairwise. Thus one can start with the free field action of any pair of fields. Adding other fields will not help.

(ii) Explicit work shows that no couplings exist for $(\frac{3}{2}, 0)$ or $(1, \frac{1}{2})$ (which has a helicity difference of $\frac{3}{2}$) or even $(\frac{5}{2}, 1)$ (the analogue of the O(2) extended supergravity model). For example, consider the spin $(0, \frac{3}{2})$ case. Since $\gamma . \epsilon = 0$:

$$\delta A = \bar{\epsilon}_\rho[\alpha \delta_{\rho\mu}]\psi_\mu \tag{24}$$

$$\delta\psi_\mu = [\beta\gamma_\mu\delta_{\rho\sigma} + \gamma\gamma_\sigma\delta_{\mu\rho}]\epsilon_\rho\partial_\sigma A. \tag{25}$$

It is easy to verify that

$$\delta \int [-\tfrac{1}{2}(\partial_\mu A)^2 - \tfrac{1}{2}\epsilon^{\mu\nu\rho\sigma}\bar{\psi}_\mu\gamma_5\gamma_\nu\partial_\rho\psi_\sigma] = 0 \tag{26}$$

requires $\alpha = \beta = \gamma = 0$.

These negative results concerning matter couplings can be understood from an algebraic point of view. The most general anticommutator for two supersymmetry generators $Q^\alpha_\mu$ subject to the condition $\gamma^\mu Q^\alpha_\mu = 0$ is given by

$$\{Q^\alpha_\mu, \bar{Q}^\beta_\nu\} = \mathrm{i}(P_\mu\gamma_\nu + P_\nu\gamma_\mu - 5\delta_{\mu\nu}\not{P} - 3\epsilon_{\mu\nu\rho\sigma}P^\rho\gamma^\sigma\gamma^5)^{\alpha\beta}. \tag{27}$$

This result follows when one requires that the right-hand side is linear in $P_\mu$ and that the charges $Q^\alpha_\mu$ are Majorana spinors and coincides with that of Baaklini (1977) if one imposes the constraint $\gamma^\mu Q^\alpha_\mu$ on the algebra of Baaklini (1977). Consequently, the Jacobi identities are satisfied if one has the usual relations between $Q^\alpha_\mu$ and the Poincaré generators.

For massless particles we go to the frame of reference where the momentum $\boldsymbol{p}$ is along the $z$ axis. Decomposing the generators into transversal and light-cone components, one has

$$Q^i_\perp = Q_x, Q_y \qquad \text{for} \qquad i = 1, 2$$

$$Q^+ = Q_3 + Q_0 \qquad Q^- = Q_3 - Q_0.$$

Using

$$p^\pm = (p_3 \pm p_0)/\sqrt{2} \qquad \text{and} \qquad \gamma^\pm = (\gamma_3 \pm \gamma_0)/\sqrt{2}$$

with $\gamma_4 = \mathrm{i}\gamma_0$ one finds the following results:

$$\begin{aligned}
&\{Q^i_\perp, Q^{j\dagger}_\perp\} = 5\delta^{ij}p_+\gamma_-\gamma_0 - 3\mathrm{i}\epsilon^{ij}p_+\gamma_-\gamma_5\gamma_0 \\
&\{Q^i_\perp, Q^\dagger_+\} = p_+\gamma^i\sqrt{2}(\gamma_+ + 2\gamma_-) \\
&\{Q^i_\perp, Q^\dagger_-\} = 0 \\
&\{Q_+, Q^\dagger_+\} = -2p_+\gamma^+\gamma^0 \\
&\{Q_+, Q^\dagger_-\} = -4p_+\gamma_-\gamma^0 \\
&\{Q_-, Q^\dagger_-\} = 0.
\end{aligned} \tag{28}$$

In a positive definite Hilbert space, this requires $Q_-=0$ as an operator, which is in contradiction to the last but one anticommutator. Therefore the supersymmetry algebra (27) cannot be realised in such a space and the coupling of matter to spin-$\frac{5}{2}$ can only be realised if the charges vanish on the matter states. These results are in agreement with the work of Haag *et al* (1975). Hence, the coupling to matter poses serious problems and it is a challenge to find a solution.

## 5. Conclusions

The O($n$) extended supergravity models have at least one- and two-loop finite quantum corrections. Only for $n>8$ does O($n$) contain SU(3)×SU(2)×U(1) as a subgroup. [O(9) ⊃ O(6)×O(3) and O(6) ≃ SU(4) ⊃ SU(3)×U(1) while O(3) ≃ SU(2)]. However, for $n>8$ these models contain spin-$\frac{5}{2}$ particles, and it was generally believed that no field theories can be constructed for particles with spin $\frac{5}{2}$ and higher. That may still be true. All we have shown is the beginning of a new theory for spin $\frac{5}{2}$ in which the free field is coupled to external sources without producing ghosts. It is not clear whether one can consistently couple also to dynamical sources. If this can be done, then this has encouraging phenomenological implications. However, our arguments indicate that the coupling to gravity and other matter configurations needs vanishing charges and absence of long-range forces.

The quantisation of the massless spin-$\frac{5}{2}$ system has been considered in Berends *et al* (1979b). Choosing the gauge fixing term

$$F_\mu=\gamma_\nu\psi_{\mu\nu}-\tfrac{1}{4}\gamma_\mu\psi \qquad (\gamma\,.\,F=0) \tag{29}$$

the effective action becomes

$$\mathscr{L}=\mathscr{L}(\text{class})-\tfrac{1}{2}\xi\bar{F}_\mu\not{\partial}F_\mu+\bar{C}^*_\mu\not{\partial}C_\mu \tag{30}$$

where from $\gamma\,.\,F=0$ and $\gamma\,.\,\epsilon=0$ it follows that

$$\bar{C}^*\,.\,\gamma=\gamma\,.\,C=0 \tag{31}$$

respectively. The BRS rules which leave this quantum action invariant are

$$\begin{aligned}&\delta\psi_{\mu\nu}=(\partial_\mu C_\nu+\partial_\nu C_\mu)\Lambda \qquad \delta C_\mu=0\\ &\delta\bar{C}^*_\mu=\Lambda\xi\bar{F}_\nu\overleftarrow{\not{\partial}}(\delta_{\mu\nu}-\tfrac{1}{4}\gamma_\nu\gamma_\mu).\end{aligned} \tag{32}$$

Of course, in a coupled theory with field-dependent structure constants, $\delta C_\mu$ will no longer vanish in general.

One can also construct a massive spin-$\frac{5}{2}$ theory without ghosts, but in this case one needs an extra spin-$\frac{1}{2}$ field $\lambda$. One can again use the spin projection formalism, but now in the $(\psi_{\mu\nu},\lambda)$ space, to invert the field equations and then find that there are no ghosts. The $(\psi_{\mu\nu},\lambda)$ action is obtained by taking a fourth root of the spin-$\frac{5}{2}$ projection operator, and choosing this fourth root such that it does not contain non-local $\square^{-1}$ terms (the root formalism is due to Ogievetski and Sokatchev (1977)). Then, after shifting fields to make this root symmetric, the action is obtained by sandwiching the shifted root between $(\psi_{\mu\nu},\lambda)$.

In the limit of vanishing mass the spin $\frac{1}{2}$ decouples, and one can discard it in the massless theory. In this sense there is a van Dam–Veltman (1970) mass discontinuity, this time already at the level of the action.

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*Note added in proof.* Since this paper was submitted several related publications on higher-spin theories have appeared (Aragone C and Deser S 1979 *Phys. Lett.* **86B** 161, *Phys. Rev.* D to be published, Berends F A and van Reisen J C J M *Nucl. Phys.* B to be published, de Wit B and Freedman D Z *Phys. Rev.* D to be published, Fronsdal C and Hata H *Nucl. Phys.* B to be published, Fransdal C *Preprint* UCLA/79/TEP/19, Curtright T 1979 *Phys. Lett.* **85B** 219). Unfortunately, the negative results reported here on the possibility of having interacting theories, with or without gravitation, have been confirmed. Up till now, no fully consistent interacting theory has been found.

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