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A NEW CLASS OF MERONIC SOLUTIONS *

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1. The study of classical solutions to field equations has represented an interesting ground of investigation both for the physical insights that such configurations can offer and for a deeper understanding of the formal aspects of the theory (1). In particular some attention has been devoted to the symmetry properties of the various solutions (2).

Among these we quote the instanton and meron solutions. Since more complicated cases rely on the properties of the scalar equation

$$\Box h(x) + \frac{\lambda^2}{2} h^3(x) = 0 \qquad (1)$$

we can refer to it for a brief orientation. While the theory leading to Eq. (1) is invariant under the full 15 parameter conformal group, particular solutions are selected by their more restricted properties of symmetry. Thus if we require that

$$\delta h_I(x) \equiv i\varepsilon_\mu \left[\frac{1-x^2}{2}\partial_\mu + (x.\partial + 2)x_\mu\right] h(x) = 0 \qquad (2)$$

then

$$h_I(x) \propto \frac{1}{1+x^2} \qquad (3)$$

and instanton solutions are determined by the invariance subgroup O(5) (if one works Euclidean). Meron solution, with

$$h_M(x) \propto \frac{1}{\sqrt{x^2}} \qquad (4)$$

are invariant under O(4) x O(1,1) (or O(4) x O(2)in the Minkowski case),

$$\delta h_M(x) \equiv \varepsilon(x.\partial + 1)\, h(x) = 0 \ . \qquad (5)$$

For more complicated systems endowed with internal degrees of freedom the analogous invariance requirements have to be supplemented with compensating operations on these additional variables. A well known example is represented by the solutions of the O(3), (or O(4)). Yang-Mills theory

$$A_\mu = A_\mu^\alpha \frac{\tau^\alpha}{2i} = -i\,\sigma_{\mu\nu}\,\partial_\nu \ln h_\mu(x) \ , \tag{6}$$

where rotational invariance is achieved provided orbital, spin and color parts are included. Similar considerations can be done for the nonlinear σ -model (in two and in four dimensions).

2. The aim of this note is to develop similar considerations in the case of dilatations when additional degrees of freedom allow solutions which, even if not strictly dilatation invariant, are however of the power type. As a starting example we consider the case of the conformal invariant interaction of charged particles i.e. using on Euclidean formulation

$$\mathcal{L} = -\left(\partial_\mu \phi^+\right)\left(\partial_\mu \phi\right) - \frac{\lambda^2}{4}\left(\phi^+\phi\right)^2 . \tag{7}$$

The equation of motion is

$$\Box\,\phi = \frac{\lambda^2}{2}\left(\phi^+\phi\right)\phi \ . \tag{8}$$

If we look for solutions invariant under four dimensional rotations and dilatations this immediately leads to the power function (being - 1 the canonical dimension of the field). However it is easy to see that the charge degree of freedom (complex ϕ!) allows a larger set of solutions, already derived several years ago by Castell using different arguments[3]. Indeed consider the general power form

$$\phi = \frac{a}{r}\,r^{\gamma} \ , \qquad \phi^+ = \frac{a^*}{r}\,r^{-\gamma} \ , \tag{9}$$

where the general definition of hermitian conjugation has been used[4]

$$\phi^+(r) = r^{-2}\,\phi\left(\frac{1}{r}\right) . \tag{10}$$

Then using the well-known relation

$$x^2\,\Box = -L^2 + x.\partial\left(x.\partial + 2\right) \tag{11}$$

one obtains

$$|a|^2 = \frac{2}{\lambda^2}\left(\gamma^2 - 1\right) . \tag{12}$$

(If γ is real and smaller than 1, as well known, the coupling constant has the wrong sign since positivity of energy requires $\lambda^2 > 0$.)

The above form of the solutions suggests a simple interpretation in terms of dilatation. Since

$$\left(r\frac{d}{dr} + 1 - \gamma\right)\phi = 0 \tag{13}$$

this amounts to requirement of invariance under the combined action of

$$\Delta = i D - \gamma\, Q \ , \tag{14}$$

where Q generators the phase transformations of first kind, $\delta\phi = i\,\varepsilon\,\phi$. In other words, the existence of the charge Q allows a sort of "anomalous" dimension γ to appear.

The solutions take a more convenient form if the singularities from O and infinity are displaced in $\pm b_\mu$, $b_\mu \equiv (0,0,0,1)$ and subsequently the rotation to Minkowski space $x_4 = it$ is performed:

$$r \to \left\{ \frac{(x-b)^2}{(x+b)^2} \right\}^{1/2} \longrightarrow \left\{ \frac{(it-1)^2 + r^2}{(it+1)^2 + r^2} \right\} \equiv e^{-i\varsigma} \quad , \qquad (15)$$

where

$$\varsigma = \operatorname{arctg}(t+r) + \operatorname{arctg}(t-r) \quad . \qquad (16)$$

Consequently, taking into account the transformation properties of the field, one gets

$$\phi \longrightarrow \frac{a}{\{(1+t_+^2)(1+t_-^2)\}^{1/2}} \; e^{-i\gamma\varsigma} \qquad (17)$$

$$\phi^* \longrightarrow \frac{a^*}{\{(1+t_+^2)(1+t_-^2)\}^{1/2}} \; e^{i\gamma\varsigma}$$

i.e. complex solutions. This casts some doubts on the effective physical significance of these results. Anyway, let us remark that observable quantities, which as a consequence of gauge invariance of first kind depend only on $(\phi^+\phi)$, turn out to be real. As well-known meron solutions have in the simple form (9) divergent action and energy; one can however improve the action and energy-momentum tensor by using the suitable conformal transformation Eq. (15). Thus the finite action is

$$I = \int d^4x \, \mathcal{L} = \frac{(\gamma^2-1)^2 \pi^3}{\lambda^2} \quad . \qquad (18)$$

The energy-momentum tensor takes on the simple form

$$\theta_{\mu\nu} = \frac{(\gamma^2-1)^2}{\lambda^2} \frac{1}{x^4} \left[-\delta_{\mu\nu} + 4 \frac{x_\mu x_\nu}{x^2} \right] \qquad (19)$$

with

$$\partial^\mu \theta_{\mu\nu} = 0 \qquad \theta^\mu_\mu = 0 \quad , \qquad (20)$$

while for the energy we obtain

$$E = \int d^3x \, \theta_{00} = \frac{4\pi^2}{\lambda^2} (\gamma^2-1)^2 \quad (E > 0, \lambda^2 > 0) \, . \qquad (21)$$

3. We want now to apply the same considerations to a model Lagrangian recently proposed by de A.F.F.[5] to describe the interaction of the gravitational field $g_{\mu\nu}$, of SU(2) gauge vector fields A^α_μ and of a neutral scalar field ϕ. The Lagrangian has the non-polynomial form

$$\mathcal{L} = -\frac{1}{4}\sqrt{g} \left\{ R + \frac{3}{2}\lambda^2\phi^2 + \frac{1}{e^2\phi^2} \sum_\alpha F^\alpha_{\mu\nu} F^\alpha_{\rho\sigma} g^{\mu\rho} g^{\nu\sigma} + \frac{2\xi}{\phi^2} g^{\mu\nu} \partial_\mu\phi \, \partial_\nu\phi \right\} \qquad (22)$$

λ^2 is the (dimensionless) cosmological constant whose presence is unavoidable to guarantee the existence of meron solutions, while ξ is a free parameter we introduce for sake of generality(*). For this Lagrangian when

$$\lambda^2 = e^2 \qquad (23)$$

the following solution exists:

$$g_{\mu\nu} = \frac{1}{C^2} \frac{\delta_{\mu\nu}}{x^2} \, , \qquad A_\mu = -i \frac{\sigma_{\mu\nu} x_\nu}{x^2} \, , \qquad \phi = c\sqrt{\frac{2}{e^2}} \qquad (24)$$

(*) Following the discussion of Ref. (5), the Newton constant does not appear in the Lagrangian density (22) and the fields $g_{\mu\nu}$, A_μ, ϕ have dimensionality -2, -1, 0 (in units of length) respectively, which ensures the dilatation invariance of the action.

C is an arbitrary dimensionless constant. (The main motivation for proposing the model was to ascertain the stability properties of the meronic solution (24). Two remarks are important. It is first of all interesting that (22) is just the effective part of the $N = 4$ Lagrangian for supergravity with SU(2) x SU(2) local invariance, apart from the fact that supersymmetry fixes uniquely the ratio between λ^2 and e^2 and the value of ξ (these are $\lambda^2 = -\frac{e^2}{3}$, $\xi = 1$). Thus the solution (24) is not supersymmetric. Secondly, the Lagrangian density exhibits the simple covariance property

$$\mathcal{L} \to u^2 \mathcal{L} \tag{25}$$

when the fields undergo the transformation of multiplication by suitable powers of an arbitrary constant u :

$$g_{\mu\nu} \to u^2 g_{\mu\nu}, \quad A_\mu \to A_\mu, \quad \phi \to u^{-1}\phi \;. \tag{26}$$

The equations of motion are clearly invariant under the transformation (26) and this explains the presence of the arbitrary constant in the solutions. Furthermore the arguments of the previous section suggest that this sort of invariance may be used to derive generalized meronic solutions with "anomalous" dimensions. The following ansatz looks therefore rather plausible (keeping (26) in mind):

$$\begin{aligned} g_{\mu\nu} &= \frac{\delta_{\mu\nu}}{c^2 X^2}\, r^{2\gamma} \\ A^M_\mu &= -i\,\frac{\sigma_{\mu\nu} X_\nu}{X^2} \\ \phi &= \sqrt{\frac{2}{e^2}}\, c\, a\, r^{-\gamma} \end{aligned} \tag{27}$$

C is an arbitrary overall constant, while "a" will be determined by the theory. We then insert our ansatz in the equations of motion

which are

$$\partial_\mu\left(\sqrt{g}\, g^{\mu\mu'} F_{\mu'\nu}\, \phi^{-2}\right) = \frac{\sqrt{g}}{\phi^2}\left[F_{\mu'\nu}, A_\mu\right] g^{\mu'\mu} \tag{28a}$$

$$\xi\, \partial_\mu\left(\sqrt{g}\, g^{\mu\nu}\phi^{-2}\partial_\nu \phi\right) = \frac{\sqrt{g}}{4}\Big(3\lambda^2\phi - \frac{4\xi}{\phi^3} g^{\mu\nu}\partial_\mu\phi\,\partial_\nu\phi - \frac{2}{e^2\phi^3}\sum_\alpha F^\alpha_{\mu\nu} F^\alpha_{\lambda\rho}\, g^{\mu\lambda} g^{\nu\rho}\Big) \tag{28b}$$

$$R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} R = -2\,\theta_{\mu\nu} \tag{28c}$$

with

$$\theta_{\mu\nu} = \frac{1}{e^2\phi^2}\Big[\sum_\alpha F^\alpha_{\mu\lambda} F^\alpha_{\nu\rho}\, g^{\lambda\rho} - \frac{1}{4} g_{\mu\nu}\sum_\alpha F^\alpha_{\lambda\rho} F^\alpha_{\lambda'\rho'}\, g^{\lambda\lambda'} g^{\rho\rho'}\Big] + \frac{\xi}{\phi^2}\Big[\partial_\mu\phi\,\partial_\nu\phi - \frac{1}{2} g_{\mu\nu}\,\partial_\lambda\phi\,\partial_\rho\phi\, g^{\lambda\rho}\Big] - \frac{3\lambda^2\phi^2}{8} g_{\mu\nu} . \tag{29}$$

It is easily checked that the above choice of the powers matches the overall X dependence of the differential equations which therefore reduce to algebraic relations.

Let's examine the various equations. Since $X^\mu F^M_{\mu\nu} = 0$, Eq. (28a) is the same as in the flat case so that A^M_μ verifies it automatically as expected. Eq. (28c) leads to two constraints (corresponding to the fact that there are two independent covariants $\frac{X_\mu X_\nu}{X^2}$ and $\delta_{\mu\nu}$) which after same manipulations turn out to be

$$\frac{3\lambda^2}{e^2} = \left[\, \xi\gamma^2 + 3(\gamma^2 - 1)\right]\left[(\gamma^2 - 1) - \xi\gamma^2\right] \tag{30}$$

$$\frac{1}{a^2} = \xi\gamma^2 - (\gamma^2 - 1) = 1 + \gamma^2(\xi - 1) \tag{31}$$

Eq. (28b) gives

$$4\xi\gamma^2 = \frac{3}{a^2}\left(1 - \frac{\lambda^2}{e^2}\right) \ . \tag{32}$$

However Eq. (32) is automatically verified and it contains no new information. This is very likely due to the covariance property (25).

4. Before we proceed to a closer look at the properties of these solutions it is fruitful to have a simple interpretation of these results, analogous to the one of section 2. Since

$$\begin{aligned} (x.\partial + 2 - 2\gamma)\, g_{\mu\nu} &= 0 \\ (x.\partial + \gamma)\, \phi &= 0 \end{aligned} \tag{33}$$

again γ plays the role of an anomalous dimension and it indicates that the symmetry operator for the new solutions is $iD - \gamma C$, with C the "generator" of transformations (26), $\delta g_{\mu\nu} = 2\varepsilon g_{\mu\nu}$, $\delta\phi = -\varepsilon\phi$. Comparing this fact with (14), one can conclude that the generator C has a role formally similar to the charge operator Q.

We first notice that for $\gamma = 0$ the meronic solution (24) and the relation $\lambda^2 = e^2$ are reproduced. Another remarkable choice corresponds to $\gamma = 1$: in this case the emerging solution is the asymptotically flat one

$$g_{\mu\nu} = \frac{\delta_{\mu\nu}}{c^2}, \quad A_\mu = -\frac{i\sigma_{\mu\nu}x_\nu}{x^2}, \quad \phi = \sqrt{\frac{2}{e^2}}\, c\, \frac{1}{r} \ , \tag{34}$$

$$\lambda^2 = -\frac{e^2}{3}\xi^2 \ . \tag{35}$$

As noticed by de A. F. F.,[6] for $\xi = 1$ this solution is compatible with $N = 4$ supergravity and it might also provide a better understanding of the role of the cosmological term.[7]

The general set of classical configurations we have found interpolates between these two and, as it should have emerged, the new element which made it possible, is represented by the appearance of the scalar field ϕ.

The same remarks about reality of solutions after Wick rotation apply here. We have as a consequence of Eqs. (28) that

$$g_{\mu\nu} = \frac{\delta_{\mu\nu}}{c^2}\,\frac{e^{-2i\gamma\xi}}{[(1+t_+^2)(1+t_-^2)]^{1/2}}$$

$$A_\mu = -i\sigma_{\mu\nu}S^\nu$$

with

$$S_\mu = \frac{t_+}{1+t_+^2}\, y_\mu^+ + \frac{t_-}{1+t_-^2}\, y_\mu^- \tag{36}$$

$$y_\mu^\pm = \left(1, \pm\frac{\vec{x}}{|\vec{x}|}\right)$$

and

$$\phi = c\, a\, e^{i\gamma\xi} \ .$$

If we however evaluate the action and the energy-momentum tensor: these quantities turn out to be real and finite. The energy momentum tensor is

$$\begin{aligned} \Theta_{\mu\nu} = \frac{1}{x^2}\Big\{ &\left(\frac{1}{4a^2} - \frac{\xi\gamma^2}{2} - \frac{3a^2}{4e^2}\right)\delta_{\mu\nu} \\ &+ \left(\xi\gamma^2 - \frac{1}{a^2}\right)\frac{x_\mu x_\nu}{x^2}\Big\} \end{aligned} \tag{37}$$

which is conserved in the covariant sense

$$\left(\Theta_{\mu}^{\nu}\right)_{;\nu} = 0 \quad . \tag{38}$$

The lagrangian (22) can be improved by means of the conformal transformation(15) which gives

$$\mathcal{L}(x) = \frac{1}{4}\left[5e^2 + \lambda^2(4a^4-9)\right]\frac{(2b)^{4-2\gamma}}{\left[(x+b)^{2+\gamma}(x-b)^{2-\gamma}\right]^2} \tag{39}$$

and it would produce a finite action in the minkowskian metric with $b \to ib$.

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