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PERTURBATION OF AN EXACT STRONG GRAVITY SOLUTION *

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ABSTRACT

Perturbations of an exact strong gravity solution are investigated. It is shown, by using the new multipole expansions previously presented, that this exact and static spherically symmetric solution is stable under odd parity perturbations.

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I. INTRODUCTION

Recently we have investigated the perturbation of strong gravity theory which was originally suggested as a gravitational analogue of vector dominance of hadron electrodynamics [1] and later on extended to include colour for the strong interacting tensor fields [2]. The difficulties associated with the non-gauge invariance of the Lagrangian have forced us to develop first a suitable form of perturbation theory for strong gravity [3]. We have also shown that the approximate strong gravity solution [4], which has been interpreted as an effective potential in order to provide an interesting mechanism for quark confinement [5], is unstable under odd parity perturbations. Later we have given new and suitable multipole expansions [6]. These have been applied to the linear wave equations and static spherically symmetric strong gravity solution. We have separated out the effective Hamiltonian of small fluctuations for the modes of a given angular momentum and parity by using these multipole expansions and the Hamiltonian formalism of general relativity. It was shown that for some values of the odd parity canonical variables the perturbed Hamiltonian is a negative definite function and thus concluded again that the approximate strong gravity solution is unstable.

However, a different choice of interaction Lagrangian leads to a different solution [3], [6]. As a matter of fact, Isham and Storey have obtained [7] a new exact solution and have also shown that the solution previously analyzed is an approximate solution.

The aim of this paper is to analyze this new exact solution by using the multipole expansion mentioned above. Here only the odd parity perturbations of exact strong gravity solution are investigated, both because of the extreme difficulty in testing the positivity of the Hamiltonian of even fluctuations [8] and because of the fact that our previous results were only valid for the odd parity perturbations. In our analysis we shall again ignore gravitation by taking $g_{\mu\nu} = \eta_{\mu\nu}$ (flat space-time). However, it turns out that this cannot be done arbitrarily and this specific value of $g_{\mu\nu}$ introduces a constraint for the existence of an exact solution. By taking into account this constraint we shall show that exact strong gravity solution is stable under odd parity perturbations only for some restricted values of parameter α appearing in the interaction Lagrangian.

The plan of the paper is as follows. In the next section the Lagrangian and exact general strong gravity solution is summarized, an exact solution which will be analyzed here is deduced and the action functional is given in ADM formalism [9]. In the last section, the perturbed action and the effective

Hamiltonian of the odd parity fluctuations are calculated. It is shown that for all of the values of the canonical variables the effective Hamiltonian is positive definite. Therefore we conclude that the strong gravity solution is stable under odd parity perturbations when the parameter α is suitably restricted.

II. SOLUTIONS

The Lagrangian from which the general solution has been obtained is [7] given by

$$\mathcal{L} = \mathcal{L}_g + \mathcal{L}_f + \mathcal{L}_{fg} ,$$

$$\mathcal{L}_g = \frac{1}{K_g^2}\sqrt{-g}\,R(g) \; ; \quad \mathcal{L}_f = \frac{1}{K_f^2}\sqrt{-f}\,R(f)$$

$$\mathcal{L}_{fg} = \frac{-M^2}{4(K_g^2+K_f^2)}\sqrt{-g}\left(\frac{f}{g}\right)^{\alpha}(f^{\mu\nu}-g^{\mu\nu})(f^{\kappa\lambda}-g^{\kappa\lambda})(g_{\kappa\mu}g_{\nu\lambda}-g_{\kappa\lambda}g_{\mu\nu}) \qquad (2.2)$$

which is just the Einstein Lagrangian for $g_{\mu\nu}$ and $f_{\mu\nu}$ fields and generally covariant mixing term which causes the interactions between the field $g_{\mu\nu}$ and $f_{\mu\nu}$. K_g and K_f are the gravitational constant and the coupling constant of the strongly interacting f-meson, respectively. M is the f-meson mass.

Variation with respect to $g^{\mu\nu}$ and $f^{\mu\nu}$ give the following g-field and f-fields equations:

$$G_{\mu\nu}(g) = R_{\mu\nu}(g) - \frac{1}{2}g_{\mu\nu}R(g) = -\frac{1}{2}\frac{K_g^2}{\sqrt{-g}}\frac{\partial}{\partial g^{\mu\nu}}\mathcal{L}_{fg}$$

$$F_{\mu\nu}(f) = R_{\mu\nu}(f) - \frac{1}{2}f_{\mu\nu}R(f) = -\frac{1}{2}\frac{K_f^2}{\sqrt{-f}}\frac{\partial}{\partial f^{\mu\nu}}\mathcal{L}_{fg} \qquad (2.3)$$

The exact solution of the above field equations are [10] (here Schwarzschild masses are taken equal to zero)

$$g_{\mu\nu}dx^\mu dx^\nu = -\left(1+\frac{r^2}{R_g^2}\right)dt^2 + \left(1+\frac{r^2}{R_g^2}\right)^{-1}dr^2 + r^2(d\theta^2+\sin^2\theta\, d\varphi^2)$$

$$f_{\mu\nu}dx^\mu dx^\nu = -C\,dt^2 + 2D\,dt\,dr + A\,dr^2 + \frac{3}{3}r^2(d\theta^2+\sin^2\theta\, d\varphi^2) \qquad (2..)$$

with

$$C = \frac{3}{2}\Delta\left(1+\frac{r^2}{R_f^2}\right) , \quad A = \left(1+\frac{r^2}{R_g^2}\right)^{-1}\left[\frac{2}{3}+\frac{3}{2}\Delta(1-X)\right.$$

$$D^2 = \Delta\left[1-\left(1+\frac{9}{4}\Delta\right)X + \frac{9}{4}X^2\right] \qquad (2.5)$$

where

$$X = \left(1+\frac{r^2}{R_g^2}\right)^{-1}\left(1+\frac{r^2}{R_f^2}\right)$$

$$\frac{1}{R_g^2} = \frac{M^2}{12}\frac{K_g^2}{K_f^2}\left(\frac{4}{9}\Delta\right)^{\alpha}\left[\frac{3}{2}\left(\alpha-\frac{1}{2}\right)-\frac{2}{\Delta}\left(\alpha-\frac{3}{2}\right)\right]$$

$$\frac{1}{R_f^2} = \frac{M^2}{18}\left(\frac{4}{9}\Delta\right)^{\alpha-\frac{1}{2}}\left[-\frac{3}{2}\alpha+\frac{2}{\Delta}(\alpha-1)\right] , \qquad (2.6)$$

and Δ is an integration constant.

We ignore the gravitational coupling ($K_g \ll K_f$) and set $g_{\mu\nu} = \eta_{\mu\nu}$ (flat space-time). This is valid only for the special value of Δ,

$$\Delta = \frac{4}{3}\,\frac{2\alpha-3}{2\alpha-1} . \qquad (2.7)$$

The constant $1/R_g^2$ vanishes for the above value of Δ, so that the Minkowski metric $g_{\mu\nu} = \eta_{\mu\nu}$ can be an exact solution. (This shows that the solution which has been examined in Refs.3 and 6 happens to be exact for Δ = " For this value of Δ, the Lagrangian of the strong gravity theory and its solution which will be analyzed under small perturbations are given as follow:

$$\mathcal{L} = \frac{1}{K_f^2}\sqrt{-f}\,R(f) + \mathcal{L}_{mass} , \qquad (2.8)$$

$$\mathcal{L}_{mass} = -\frac{M^2\sqrt{-\eta}}{4K_f^2}\left(\frac{f}{\eta}\right)^{\alpha}(f^{\mu\nu}-\eta^{\mu\nu})(f^{\kappa\lambda}-\eta^{\kappa\lambda})(\eta_{\kappa\mu}\eta_{\lambda\nu}-\eta_{\kappa\lambda}\eta_{\mu\nu}) , \qquad (2.9$$

$$\eta_{\mu\nu}dx^\mu dx^\nu = -dt^2 - \frac{2D}{C}dr\,dt + \left(1-\frac{D^2}{C^2}\right)dr^2 + r^2(d\theta^2+\sin^2\theta\, d\varphi^2)$$

$$f_{\mu\nu}dx^\mu dx^\nu = -C\,dt^2 + \frac{\Delta}{C}dr^2 + \frac{2}{3}r^2(d\theta^2+\sin^2\theta\, d\varphi^2) \qquad (2.10)$$

where

$$C = \frac{3}{2}\Delta\left(1+\frac{r^2}{R_f^2}\right), \quad A + C = \frac{2}{3} + \frac{3}{2}\Delta, \quad D^2 = \Delta - AC \qquad (2.11)$$

Here the solution is expressed most simply in spherical polar coordinates with a retarded time. This solution is similar to the previously investigated solution under small perturbations [11] for $\alpha = 0$ and $\Delta = 4$. At the end of our calculations we will therefore be able to check our previous calculations by taking $\alpha = 0$.

Since $\sqrt{-f}\; R(f)$ is identical in form with the Einstein-Lagrangian, the canonical action functional can be written in the Arnowitt, Deser and Misner form [9]. After discarding the surface terms, it reads [3], [6]

$$I = \frac{1}{K_f^2}\int d^4x\,[\pi^{ij}\dot{\gamma}_{ij} - N_t\mathcal{H}^t - N_i\mathcal{H}^i + \mathcal{L}_{mass}(N,\gamma)] \quad , \qquad (2.12)$$

where π and γ are canonical variables and N's are auxiliary. The $f^{\mu\nu}$ which denotes the inverse of $f_{\mu\nu}$ is expressed in terms of γ and N by

$$f^{\mu\nu} = \frac{1}{N_t^2}\begin{vmatrix} -1 & N^j \\ N^i & N_t^2\gamma^{ij} - N^iN^j \end{vmatrix} \quad . \qquad (2.13)$$

The indices i,j take values r, θ, φ. The other functions appearing in (2.12) are defined by

$$\mathcal{H}^t = \gamma^{-1/2}(\pi^{ij}\pi_{ij} - \tfrac{1}{2}\pi^2) - \gamma^{1/2}R(\gamma)$$

$$\mathcal{H}^i = -2\pi^{ij}{}_{|j}$$

$$\mathcal{L}_{mass}(N,\gamma) = -\frac{M^2}{4}\sqrt{-\eta}\left(\frac{f}{\eta}\right)^\alpha\Big[-6\left(2+\frac{\eta_{tt}}{N_t^2}\right) + 12\frac{\eta_{ti}}{N_t^2}N^i - \frac{6}{N_t^2}\eta_{ij}N^iN^j$$
$$+ 4\frac{N^\ell}{N_t^2}\gamma^{ij}(\eta_{tj}\eta_{\ell i} - \eta_{ij}\eta_{t\ell}) + \frac{2N^kN^\ell}{N_t^2}\gamma^{ij}(\eta_{ij}\eta_{k\ell} - \eta_{ik}\eta_{j\ell})$$
$$+ \frac{2}{N_t^2}\gamma^{ij}\{(3N_t^2+\eta_{tt})\eta_{ij} - \eta_{it}\eta_{jt}\} + \gamma^{ij}\gamma^{k\ell}(\eta_{ik}\eta_{j\ell} - \eta_{ij}\eta_{k\ell})\Big] \qquad (2.14)$$

The metric γ_{ij} is used to raise indices, vertical bar denotes the covariant differentiation with respect to γ, $\pi \equiv \pi^i{}_i$ and $R(\gamma)$ is the three-dimensional curvature scalar constructed from γ.

III. PERTURBATIONS

Here we have an exact classical solution to a set of coupled nonlinear equations and want to test this solution for stability against small perturbations. In essence, the problem is to find the effective Hamiltonian functional of the fluctuation variables and test its positivity.

The classical solution is static so that $\pi^{ij}_{c\ell} = 0$. Since the metric is diagonal $(N_i)_{c\ell} = f_{ti} = 0$. On the other hand the fluctuation of N_t does not contribute to the odd part [6]. Therefore we define only small fluctuations q, p, and n_i by

$$\gamma_{ij} = \gamma^{c\ell}_{ij} + q_{ij}, \quad \gamma^{ij} = \gamma^{ij}_{c\ell} - q^{ij} + q^{ik}q_k{}^j$$
$$N_i = n_i, \quad N^i = n^i - q^{ij}n_j, \quad \pi^{ij} = p^{ij}$$
$$f^\alpha = (f_{c\ell})^\alpha[1 - \tfrac{1}{2}\alpha\, q^{ij}q_{ij}] \qquad (3.1)$$

Inserting these expressions in the action functional and retaining only the terms which are bilinear in the fluctuations one gets the form

$$
\begin{aligned}
[I]^{odd}_{pert} = \frac{1}{K_f^2}\int d^4x \Big\{ & p^{ij} q_{ij,t} - N_t \gamma_{cl}^{1/2} p^{ij} p_{ij} + 2 n_i p^{ij}{}_{|j} \\
& - \frac{1}{4} N_t \gamma_{cl}^{1/2} q^{ij|k} (q_{ij|k} - 2 q_{ik|j}) - \frac{M^2\sqrt{-\eta}}{4}\left(\frac{f}{\eta}\right)^{\alpha}\Big[-\frac{6}{N_t^2} n^i n^j n_{ij} \\
& - \alpha\left(\frac{3}{4} - \frac{1}{\Delta}\right) q^{ij} q_{ij} + \frac{2}{N_t^2} q^{ik} q_k{}^j \{(3N_t^2 - 1)\eta_{ij} - \eta_{it}\eta_{jt}\} - \frac{12}{N_t^2} q^{ij} n_j \eta_{ti} \\
& + (2\gamma_{cl}^{ij} q^{km} q_m{}^l - \frac{2}{N_t^2}\gamma_{cl}^{ij} n^k n^l + q^{ij} q^{kl})(\eta_{ik}\eta_{jl} - \eta_{ij}\eta_{kl}) \\
& + \frac{4}{N_t^2}(q^{ij} n^l + \gamma_{cl}^{ij} q^{lm} n_m)(\eta_{ij}\eta_{tl} - \eta_{tj}\eta_{il}) \Big] \Big\}
\end{aligned}
\tag{3.2}
$$

Here the raising and lowering of indices is with respect to the classical tensor γ^{cl}_{ij} and vertical bar "|" denotes covariant differentiations with respect to this metric.

In order to facilitate the angular decomposition, it is useful to introduce a spherical helicity basis

$$
\begin{aligned}
& \gamma^{cl}_{ij} dx^i dx^j = \omega^+\omega^- + \omega^-\omega^+ + \omega^0\omega^0 \\
& \eta_{\mu\nu} dx^\mu dx^\nu = -dt^2 - \frac{2D}{aC} dt\,\omega^0 + \frac{1}{a^2}\left(1 - \frac{D^2}{C^2}\right)\omega^0\omega^0 + \frac{3}{2}(\omega^+\omega^- + \omega^-\omega^+)
\end{aligned}
\tag{3.3}
$$

where the one-forms are defined by

$$
\begin{aligned}
& \omega^r = a\,dr, \quad \omega^\theta = b\,d\theta, \quad \omega^\varphi = b\sin\theta\, d\varphi, \quad a^2 = \frac{\Delta}{C}, \quad b^2 = \frac{2}{3}r^2 \\
& \omega^0 = \omega^r, \quad \omega^+ = \frac{1}{\sqrt{2}}(\omega^\theta + i\omega^\varphi), \quad \omega^- = \frac{1}{\sqrt{2}}(\omega^\theta - i\omega^\varphi)
\end{aligned}
\tag{3.4}
$$

In terms of helicity components the action can be written as follows

$$
\begin{aligned}
[I]^{odd}_{pert} = \frac{1}{K_f^2}\int d^4x \Big[& p_{--} q_{++,t} + p_{++} q_{--,t} + 2(p_{+0} q_{-0,t} + p_{-0} q_{+0,t}) \\
& - N_t \gamma_{cl}^{1/2} \mathcal{H}_{odd} \Big]
\end{aligned}
\tag{3.5}
$$

where

$$
\begin{aligned}
\mathcal{H}_{odd} = 2\gamma^{-1}[|p_{++}|^2 + 2|p_{+0}|^2] & + \frac{1}{2}[-|q_{++|+}|^2 + |q_{++|0}|^2 + |q_{++|-}|^2 + 2|q_{+0|-}|^2 \\
& + |q_{00|+}|^2 - |q_{+-|0}|^2] - 2[\frac{1}{2}(q_{0+|-})^2 + q_{++|-} q_{-+|-} + q_{+|0} q_{-0|-} \\
& + q_{0+|0} q_{00|-} + C.C.] + \frac{M^2\sqrt{-\eta}}{4N_t\gamma^{1/2}}\left(\frac{f}{\eta}\right)^{\alpha}\Big[\{-2\alpha\left(\frac{3}{4} - \frac{1}{\Delta}\right) + \frac{27}{2} - \frac{6}{N_t^2} \\
& - 6\eta_{00}\}|q_{++}|^2 - \frac{1}{N_t^2}(6\eta_{00} - 9)|n_+|^2 + \{-4\alpha\left(\frac{3}{4} - \frac{1}{\Delta}\right) + 9 \\
& - \frac{6}{N_t^2} - \frac{4}{a^2 N_t^2}\}|q_{+0}|^2\Big]
\end{aligned}
\tag{3.6}
$$

In obtaining Eq.(3.5) we picked up only the terms which contributed to the abnormal parity modes and eliminated n's by setting $\partial H/\partial n = 0$. In Eq.(3.6) n_+ reads

$$
n_+ = \frac{1}{\frac{2}{3}\frac{C}{\Delta} - 1}\left[-\frac{D}{a} q_{+0} + \frac{4C^2}{3M^2\sqrt{-\eta}}\left(\frac{\eta}{f}\right)^{\alpha} p_{-}{}^{j}{}_{|j}\right] \quad . \tag{3.7}
$$

Here we use the new multipole expansions

$$
\left.
\begin{aligned}
& q_{++} = \sum (2j+1) q_2(j,m)\, D^j_{2,m}(L_r^{-1}) \\
& q_{+0} = \sum (2j+1) q_1(j,m)\, D^j_{1,m}(L_r^{-1})
\end{aligned}
\right| \tag{3.8}
$$

and likewise for the momenta P_{ab} [12]. Now we insert the above expansions in Eq.(3.5) define even and odd combinations $q_{2\pm}$ and $q_{1\pm}$ by

$$
\begin{aligned}
& q_2 = \frac{1}{\sqrt{2}}(q_{2+} - i q_{2-}), \quad q_{-2} = \frac{1}{\sqrt{2}}(q_{2+} + i q_{2-}) \\
& q_1 = \frac{1}{2}(q_{1+} - i q_{1-}), \quad q_{-1} = \frac{1}{2}(q_{1+} + i q_{1-})
\end{aligned}
\tag{3.9}
$$

and similarly for p_2 and p_1. Then the action functional separates into a sum of contributions according to angular momentum and parity. If we ignore the normal parity modes and perform the angular integrations we obtain (at $m = 0$ where all variables are real),

$$[I]^{odd}_{pert} = \frac{8\pi}{3K_f^2} \sum (2j+1) \int dt\, dr\, r^2 \left(\frac{\Delta}{C}\right)^{1/2} \left[p_{1-} \partial_t q_{1-} + p_{2-} \partial_t q_{2-} - C^{1/2} \mathcal{H}^*_{odd} \right] \tag{3.10}$$

where the Hamiltonian $\mathcal{H}^*_{odd}$ is given by

$$\begin{aligned} \mathcal{H}^*_{odd} = {} & p_{1-}^2 + p_{2-}^2 + \frac{3}{8}\left[\sqrt{(j-1)(j+2)}\, \frac{q_{1-}}{r} + \sqrt{\frac{2C}{3\Delta}}\, \partial_r q_{2-}\right]^2 \\ & + \frac{3M^2}{8\sqrt{\Delta}} \left(\frac{4\Delta}{9}\right)^{\alpha} \left[\frac{1}{\Delta}\left(\frac{9}{4}\Delta - 1\right) q_{2-}^2 + \frac{3}{2C}\left(\frac{3}{2}C - 1\right) q_{1-}^2 \right] \\ & + \frac{27}{4r^2C^2} \frac{1}{\left[\frac{3}{2}\alpha - \frac{3}{\Delta}(\alpha-1)\right]} \left[\frac{D}{a} q_{1-} + \frac{4C^2}{3M^2}\left(\frac{\eta}{\gamma}\right)^{\alpha - 1/2} \left\{ \frac{\sqrt{(j-1)(j+2)}}{b} p_{2-} \right.\right. \\ & \left.\left. - \frac{1}{a}\left(\frac{3}{r} p_{1-} + \partial_r p_{1-}\right) \right]^2 \end{aligned} \tag{3.11}$$

In order to show that the solution given by Eq.(.10) is stable, we have to prove that $\mathcal{H}^*_{odd}$ is a positive definite function for all of the arbitrary values of the canonical variables. However, taking the value Δ ($\Delta > 0$ for $\alpha < 0$) given by Eq.(2.7) for which our solution is exact strong gravity solution, it is seen that all of the coefficients

$$\frac{9}{4}\Delta - 1 = \frac{4\alpha - 8}{2\alpha - 1}$$

$$\frac{3}{2}\alpha - \frac{2}{\Delta}(\alpha - 1) = \frac{1}{2(3-2\alpha)}$$

$$\frac{3}{2C}\left(\frac{3}{2}C - 1\right) = \frac{9\Delta}{4C^2}\left[\frac{3}{2(3-2\alpha)} + \frac{r^2}{R_f^2}\,\frac{10\alpha - 17}{2\alpha - 1} + \frac{9\Delta r^4}{4R_f^4}\right],$$

are positive for $\alpha < 0$ and consequently for all the arbitrary values of canonical variables,

$$\mathcal{H}^*_{odd} > 0 \quad \text{for} \quad \alpha < 0.$$

We have investigated the perturbations of the static spherically symmetric solution of the strong gravity theory and have shown that the effective perturbed Hamiltonian is a positive definite function of small perturbations. Therefore we conclude that the general solution of the strong gravity is stable under odd parity perturbations for $\alpha < 0$. This result agrees with the result of Refs.[3,6] where it has been shown that for $\alpha = 0$ and $\Delta > \frac{9}{4}$ the strong gravity solution is unstable under odd parity perturbations.

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11) See the footnote 7 of Ref.3.

12) In Ref.6, the reality conditions, the expansions of the covariant derivatives as well as other properties of the new multipole expansions have been given in detail. Unfortunately, there are some misprints in Eqs.(A.11) and (A.12) of Ref.6. The q_1 in Eqs.(A.11) and (A.12) should read $-q_1$.

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