SCALAR, GAUGE AND KALB–RAMOND FIELD LOCALIZATION ON A BRANE WITH GENERALIZED DYNAMICS

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In this paper, we investigate the scalar, gauge and Kalb–Ramond field localization on a one-scalar generated brane with nonstandard kinetic terms coupled with gravity. We show that the massless zero mode of spin 0 scalar field is localized on the brane with generalized dynamics, while the vector gauge field is not localized. In order to circumvent this problem, we use a functional of the scalar filed in the gauge field action to obtain vector gauge field localization in this braneworld model. We also study the localization of the Kalb–Ramond field via this procedure.

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1. Introduction

Extra dimension is an important subject in the realm of theoretical physics that provides many creative ways to solve some problems in physics such as *e.g.* hierarchy problem. Unlike the model of Arkani-Hamed *et al.* [1], Randall and Sundrum (RS) proposed an alternative scenario [2] to solve the hierarchy problem that does not require large extra dimensions. In the RS model, the size of extra dimension, r_c, is not determined by the dynamic of the model. For this scenario to be relevant, it is necessary to find a mechanism for generating a potential to stabilize the value of r_c. This mechanism which was proposed by Goldberger and Wise (GW) [3] could stabilize the size of extra dimension by a five-dimensional bulk scalar field with usual dynamics allowed to interact with gravity. Recently, Bazeia *et al.* [4] have modified the standard braneworld scenario with the inclusion of scalar fields with nonstandard dynamics. They have developed the first-order formalism for models with standard gravity but with the scalar fields having generalized

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dynamics. Other studies about braneworld models in higher dimensional space-time, which are generally based on gravity coupled to one or several scalars can be found in [5, 6].

The issue of localization of several fields and resonances in such branes is an interesting subject, as their investigation can guide us to which kind of brane structure is more acceptable phenomenologically. Generally, massless scalar fields and graviton can be localized on brane of different types. In the RS model, graviton and spin 0 field can be localized on a brane with positive tension and spin 1/2 and 3/2 fields are localized on a brane with negative tension, while the gauge field cannot be localized either on a brane with positive tension or on a brane with negative tension [7, 8]. In order to localize gauge field, the authors of [9] introduced an additional scalar field called the dilaton which can be coupled to the kinetic term of gauge fields and leads to the localization of gauge field. It was shown in Ref. [10] that gauge field localization is obtained via kinetic terms induced by localized fermions. Also in Ref. [11], the authors introduced a suitable function in the higher dimensional gauge field action to achieve gauge field zero mode localization on a thick brane generated by the coupling of a scalar field and gravity. In recent years, many authors have investigated the localization of gauge fields on branes of different types which can be found in Refs. [12–17].

The Kalb–Ramond (KR) field is an antisymmetric tensor field which was first introduced in the string theory. However, the free KR field cannot be localized on the brane in five-dimensional space-time [18]. But if the KR field couples to a dilaton field [19] or with the background kink scalar field [20], the localization of the KR field can be realized. More studies about tensor gauge field localization can be found in [11, 21].

Furthermore, it is important whether the fermions could be localized on the thick branes. By introducing the scalar–fermion coupling, fermions have normalizable zero modes in higher dimensional space-time [22–24]. In Ref. [24] the fermion localization has been considered on a thick brane constructed from one scalar field with nonstandard kinetic terms coupled to gravity. The authors of [24] used the analytical expressions for small α and investigated the contribution of this nonstandard kinetic terms to the problem of fermion localization.

The plan of this paper is as follows. In Sec. 2 we briefly summarize the braneworld model with generalized dynamics developed by Ref. [4]. In Sec. 3 we study the localization of scalar field on this brane. We give the analytical expression for the scalar zero mode in the thin brane limit. In Sec. 4 we discuss gauge field localization. The brane with generalized dynamics is not capable of supporting the existence of a localized zero mode of the gauge field on the brane. Following recent results in Ref. [11], we use a smearing out function in the gauge field action to localize gauge field on this brane. Finally, in Sec. 5 we conclude with the summary of our results.

2. The braneworld models with generalized dynamics

The action of a bulk scalar field, ϕ, on the warped braneworld model with nonstandard kinetic terms which has been described by Ref. [4] is

$$S = \int d^5x\sqrt{|g|}\left[\frac{-1}{4}R + L(\phi, X)\right], \tag{1}$$

where R is the five-dimensional scalar curvature and $g \equiv \mathrm{Det}(g_{MN})$. The Lagrangian density $L(\phi, X) = K(X) - V(\phi)$, where $K(X)$ and $V(\phi)$ are the nonstandard kinetic term and the potential, respectively. In Ref. [4], the authors have studied two specific forms for the nonstandard kinetic term, $K(X) = X + \alpha|X|X$ and $K(X) = -X^2$, where α is a real non-negative parameter which drives the model away from the standard case and $X = -\frac{\dot{\phi}^2}{2}$, where the dot is used to represent derivative with respect to y. The general form of the warped metric for a five-dimensional space-time is given by

$$ds^2 = g_{MN}dx^M dx^N = e^{2A(y)}g_{\mu\nu}dx^\mu dx^\nu - dy^2\,, \tag{2}$$

where $M, N = 0, 1, 2, 3, 4$ and e^{2A} is the warp factor. Also $g_{\mu\nu}$ stands for four-dimensional curved brane with signature $(+, -, -, -,)$ and $\mu, \nu = 0, 1, 2, 3$. We assume that the scalar field and the warp factor depend only on the extra coordinate y. In this article, we choose the first model for $K(X)$ as $K(X) = X + \alpha|X|X$ [24]. With these assumptions, the action given by Eq. (1) with the ansatz Eq. (2) leads to the following equations of motion for the scalar field

$$\begin{aligned} \ddot{\phi} + 4\dot{A}\dot{\phi} - V_\phi &= -\alpha\left(3\ddot{\phi} + 4\dot{\phi}\dot{A}\right)\dot{\phi}^2\,, & (3)\\ \ddot{A} + 2\dot{A}^2 &= \frac{-1}{3}\left(1 + \frac{\alpha}{2}\dot{\phi}^2\right)\dot{\phi}^2 - \frac{2}{3}V\,, & (4)\\ \dot{A}^2 &= \frac{1}{6}\left(1 + \frac{3}{2}\alpha\dot{\phi}^2\right)\dot{\phi}^2 - \frac{1}{3}V\,. & (5)\end{aligned}$$

Now, we apply the first order formalism to the braneworld scenario [25], hence

$$\dot{A} = -\frac{1}{3}W(\phi)\,, \tag{6}$$

where the superpotential $W(\phi)$ is, in principle, an arbitrary function of the field ϕ with the corresponding potential

$$V(\phi) = \frac{1}{8}W_\phi^2 - \frac{\alpha}{64}W_\phi^4 - \frac{1}{3}W^2\,. \tag{7}$$

We consider the superpotential $W(\phi)$ of the form [6]

$$W(\phi) = 3bc \sin\left(\sqrt{\frac{2}{3b}}\phi\right) , \tag{8}$$

where b and c are positive parameters that are related to the thickness of the brane (c) and the anti-de Sitter curvature (bc). By using the above superpotential, one can obtain [24]

$$A(y) = b\ln[\mathrm{sech}(cy)] + \frac{3}{4}\alpha b^2 c^2 \tanh^2(cy) , \tag{9}$$

$$\phi(y) = \sqrt{\frac{3b}{2}}\arcsin[\tanh(cy)] - \frac{3\sqrt{6}}{4}\alpha b^{\frac{3}{2}} c^2 \tanh(cy)\mathrm{sech}(cy) . \tag{10}$$

3. Localization of scalar fields

In this section, we first study the localization of the zero mode for spin 0 scalar field on the braneworld models with nonstandard kinetic term assumed to be $K(X) = X + \alpha|X|X$. After that, we analyze the zero mode of scalar field in the thin brane limit and we discuss the relevance of the parameter α on the localization of zero mode of the scalar field.

3.1. Localized zero mode

Let us consider the action of a massless real scalar coupled to gravity

$$S_0 = \frac{1}{2}\int d^5x\sqrt{g}g^{MN}\partial_M\Phi\partial_N\Phi , \tag{11}$$

from which the equation of motion can be derived

$$\frac{1}{\sqrt{g}}\partial_M\left(\sqrt{g}g^{MN}\partial_N\Phi\right) = 0 . \tag{12}$$

If we decompose Φ as $\Phi(x;y) = \zeta(x)\chi(y)$ and demand $\zeta(x)$ to satisfy the four-dimensional massive Klien–Gordon equation $g^{\mu\nu}\partial_\mu\partial_\nu\zeta = -m^2\zeta$, we can get the following equation for the y dependence

$$4\dot{A}\frac{\partial\chi}{\partial y} + \frac{\partial^2\chi}{\partial^2 y} + m^2 e^{-2A(y)}\chi = 0 . \tag{13}$$

The above equation is very similar to the equation for gravity localization [2, 7]. For $m^2 = 0$, we can obtain the zero-mass solution to this equation which takes the form $\chi(y) = \chi_0$, where χ_0 is a constant. The condition for having a localized four-dimensional scalar field is that χ_0 be normalizable. We show that this constant mode is localized on the braneworld model driven by a real scalar field with nonstandard dynamic. Substituting the zero mode into the starting action Eq. (11), we get

$$\frac{1}{2}\chi_0^2 \int\limits_{-\infty}^{+\infty} dy e^{2A(y)} \int d^4x g^{\mu\nu}\partial_\mu\zeta\partial_\nu\zeta\,. \tag{14}$$

In order to localize zero mode, the y dependent part of the above expression, namely

$$I_0 = \int\limits_{-\infty}^{\infty} \mathrm{sech}^{2b}(cy) e^{\frac{3}{4}\alpha b^2 c^2 \tanh^2(cy)} dy\,, \tag{15}$$

should be finite. It is difficult to solve this integral analytically but noting that since $\tanh^2(cy) \leq 1$ for an arbitrary finite value of y, we can write

$$I_0 < \int\limits_{-\infty}^{\infty} P(y) dy\,, \tag{16}$$

where

$$P(y) = \mathrm{sech}^{2b}(cy) e^{\frac{3}{4}\alpha b^2 c^2}\,. \tag{17}$$

It is obvious that for any positive b, the function $P(y)$ is sharp on the core of brane and exhibits a narrow bell-shape profile. Therefore, it is convergent in the entire domain of the extra coordinate. Hence I_0 is finite and scalar zero mode localization can be achieved on the brane with generalized dynamics. It is interesting to see how scalar field localization occurs on the thin brane with nonstandard kinetic terms. Next, we study the zero mode localization and the Kaluza–Klein states of the scalar field in the thin brane limit.

3.2. Scalar field localization in the thin brane limit

By analyzing exact solutions (9) and (10) in the thin brane limit ($c \to \infty$ and the product bc is held fixed), one can obtain [24]

$$\phi(y) = \frac{\sqrt{6b}}{4}\pi \mathrm{sgn}(y)\,, \tag{18}$$

$$A(y) = -bc|y| + \frac{3}{4}\alpha b^2 c^2\,, \tag{19}$$

where the term $\frac{3}{4}\alpha b^2 c^2$ is the contribution of nonstandard kinetic term to the brane model geometry. When $\alpha = 0$, the function $A(y)$ given by Eq. (19)

reduces to the solution of RS model. The metric is given by Eq. (2), but it is more convenient to change it to a conformally flat metric as

$$ds^2 = e^{2A(y)}\left(g_{\mu\nu}dx^\mu dx^\nu - dz^2\right), \tag{20}$$

where the relation of the new coordinate z and y is $dz = e^{-A(y)}dy$. By considering the above conformally flat metric and using the transformation $\chi = e^{\frac{-3}{2}A(y)}\bar{\chi}$, equation (13) can be rewritten as

$$\left\{-\frac{d^2}{dz^2} + V(z)\right\}\bar{\chi} = m^2\bar{\chi}\,, \tag{21}$$

where the potential $V(z)$ is

$$V(z) = \tfrac{3}{2}\partial_z^2 A(z) + \tfrac{9}{4}(\partial_z A(z))^2\,. \tag{22}$$

The potential depends only on the warp factor exponent A and has the same form as the case of graviton. The asymptotical behavior of the potential gives us information about the presence of gaps in the continuum spectrum. In the thin brane limit, the relation between y and z is given by

$$\gamma|z| = \left(e^{bc|y|} - 1\right), \tag{23}$$

where

$$\gamma = bce^{\frac{3}{4}\alpha b^2c^2}\,. \tag{24}$$

Hence the potential of Eq. (22) can be written as

$$V(z) = \frac{15}{4}\frac{\gamma^2}{(\gamma|z|+1)^2} - \frac{3\delta(z)\gamma}{\gamma|z|+1}\,. \tag{25}$$

This potential is like the singular one found in the RS scenario [7]. We can see $V(z)$ tends to zero when $|z| \longrightarrow \infty$ and the shape of this potential looks like a volcano. This means that the potential provides no mass gap to separate the scalar zero mode from KK modes. In the thin brane limit, the scalar zero-mode that is the solution of the Schrödinger-like equation with $m^2 = 0$ can be given by

$$\bar{\chi}_0 = \frac{e^{\frac{9}{8}\alpha b^2c^2}}{(\gamma|z|+1)^{\frac{3}{2}}}\,, \tag{26}$$

where $\bar{\chi}_0$ is normalizable when the product bc is held fixed. The shape of the zero mode for scalar field is plotted in Fig. 1. This figure shows that as the value of α decreases, the zero-mode gets more and more localized on

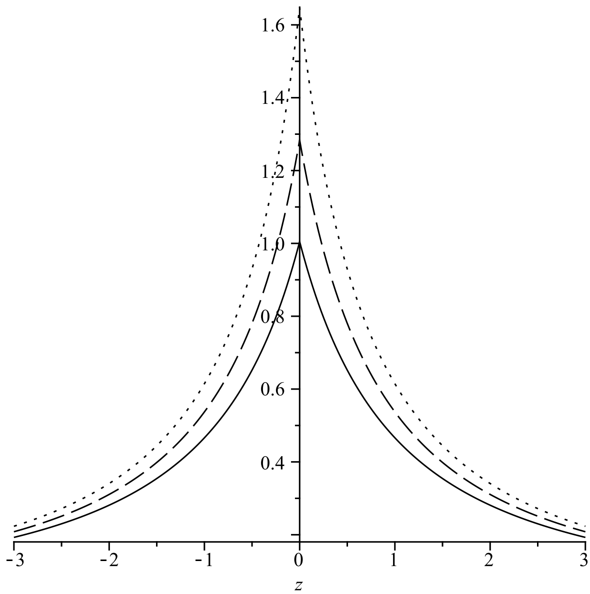

Fig. 1. The shape of the scalar field zero-mode $\bar{\chi}_0$ for $b = \frac{2}{3}$, $c = 1$ and $\alpha = 0.01$ (solid line), $\alpha = 0.5$ (dashed line), $\alpha = 1$ (dotted line).

the brane with generalized dynamic. Also the Schrödinger-like equation (21) can be rewritten as

$$\left(-\partial_z + \tfrac{3}{2}\partial_z A(z)\right)\left(\partial_z + \tfrac{3}{2}\partial_z A(z)\right)\bar{\chi} = m^2\bar{\chi}\,. \tag{27}$$

This factorization directly shows that there are no normalizable modes with negative m^2, namely, there is no tachyonic scalar mode. Thus, the scalar zero mode is the lowest mode in the spectrum. In addition to this massless mode, the potential expressed by Eq. (25) has a continuum gapless spectrum of KK modes with positive $m^2 > 0$ that can be given in terms of Bessel functions, which are similar to those obtained in [6, 7, 26].

4. Localization of gauge field

In this section, we investigate the zero mode of the vector gauge field on the brane with nonstandard kinetic term. The action of U(1) vector field is described by

$$S = -\frac{1}{4}\int d^5x\sqrt{g}g^{MN}g^{RS}F_{MR}F_{NS}\,, \tag{28}$$

where the field strength tensor is given by $F_{MN} = \partial_M A_N - \partial_N A_M$. From this action, one can determine the equations of motion

$$\frac{1}{\sqrt{g}}\partial_M\left(\sqrt{g}g^{MN}g^{RS}F_{NS}\right) = 0\,. \tag{29}$$

We assume that A_μ are Z_2-even and that A_4 is Z_2-odd with respect to the extra dimension y, which results in that A_4 has no zero mode in the effective 4-dimensional theory. Furthermore, in order to be consistent with the gauge invariant equation $\int dyA_4 = 0$, we choose our gauge condition, as $A_4 = 0$. Under these assumptions and by decomposing the vector field as $A_\mu(x,y) = a_\mu(x)\rho(y)$, the $\rho(y)$ satisfies the differential equation

$$2\dot{A}\frac{\partial\rho(y)}{\partial y} + \frac{\partial^2\rho(y)}{\partial^2 y} + m^2 e^{-2A(y)}\rho(y) = 0\,. \tag{30}$$

It is easy to find that the equation above has a constant solution $\rho(y) = \rho_0$ with $m = 0$. In order to see whether it is a normalizable one or not, we must substitute this solution in the action Eq. (28) so that

$$S = -\frac{1}{4}\rho_0^2\int\limits_{-\infty}^{+\infty}dy\int d^4x g^{\mu\nu}g^{\alpha\beta}f_{\mu\alpha}f_{\nu\beta}\,, \tag{31}$$

where we have used $f_{\mu\nu} = \partial_\mu a_\nu - \partial_\nu a_\mu$. Obviously, the suppressing warp factor is now absent of the effective action and the integral is divergent which shows that the vector field cannot be localized on the brane. Also we have checked that other normalizable nontrivial solutions of the equation of motion for the vector field do not exist. The result is the same as the RS model case, *i.e.* the zero mode of the spin 1 vector field cannot be localized either on a brane with positive tension or on a brane with negative tension so the Dvali–Shifman mechanism [27] must be considered for the vector field localization.

4.1. Using a smearing out function to localize gauge fields

Recently, the authors of [11] have proposed a model that leads to gauge field zero mode localization on thick branes by means of an effective model obtained via the introduction of a smearing out function in the higher-dimensional gauge field action. Here, we use their model to investigate gauge field zero-mode localization on the brane with generalized dynamic. Consider a suitable function in the five-dimensional gauge field Lagrangian, which leads to a normalizable zero mode after the dimensional reduction such as proposed in Ref. [11]

$$S = -\frac{1}{4}\int d^5x\sqrt{g}G(\phi)g^{MN}g^{RS}F_{MR}F_{NS}\,, \tag{32}$$

where $G(\phi)$ is a smearing out function of the minimum energy solution, which demonstrates the brane. From this action, the equation of motion is given by

$$\frac{1}{\sqrt{g}}\partial_M\left(\sqrt{g}G(\phi)g^{MN}g^{RS}F_{NS}\right)=0\,. \tag{33}$$

From this equation of motion and the conditions stated in the beginning of this section for the vector field A_M, one arrives at

$$\left(2\dot{A}(y)+\frac{\dot{G}(\phi)}{G(\phi)}\right)\frac{\partial\rho(y)}{\partial y}+\frac{\partial^2\rho(y)}{\partial^2 y}+m^2e^{-2A(y)}\rho(y)=0\,. \tag{34}$$

For the zero mode of the gauge field, $\rho(y)=\rho_0$ is a solution. Substituting the zero mode into the action (32), we get

$$S=-\frac{1}{4}\rho_0^2\int\limits_{-\infty}^{+\infty}G(\phi)dy\int d^4xg^{\mu\nu}g^{\alpha\beta}f_{\mu\alpha}f_{\nu\beta}\,. \tag{35}$$

In order to localize gauge fields on the brane with nonstandard kinetic term, one requires that $G(\phi)$ should be normalizable in the entire domain of the extra coordinate. In order to obtain the function $G(\phi)$, we use the procedure of Ref. [20]. By defining $\rho(y)=e^{-pA}\tilde{\rho}(z)$, where p is a coupling constant and using conformally flat metric, Eq. (34) can be rewritten as

$$\begin{aligned}&\partial_z^2\tilde{\rho}(z)+\left[(1-2p)\partial_zA+\frac{\partial_zG}{G}\right]\partial_z\tilde{\rho}(z)\\&+\left[-p\partial_z^2A+\left(p^2-p\right)(\partial_zA)^2-p\partial_zA\frac{\partial_zG}{G}+m^2\right]\tilde{\rho}(z)=0\,.\end{aligned} \tag{36}$$

By discarding terms of first order in derivatives, we have

$$(1-2p)\partial_zA+\frac{\partial_zG}{G}=0\,. \tag{37}$$

From the above equation, we can obtain

$$G(z)=e^{(2p-1)A(z)}\,, \tag{38}$$

and by using the expressions $\partial_zA=e^{A(y)}\partial_yA$ and $\partial_zG=e^{A(y)}\partial_yG$, we get

$$G(\phi)=\mathrm{sech}^{b(2p-1)}(cy)e^{\frac{3}{4}\alpha b^2c^2(2p-1)\tanh^2(cy)}\,. \tag{39}$$

For $p>\frac{1}{2}$, $G(\phi)=1$ on the core of the brane and $G(\phi)\to 0$ when $y\to\pm\infty$. Hence, $\int_{-\infty}^{+\infty}G(\phi)dy$ is convergent for finite product of αb^2c^2 and $p>\frac{1}{2}$

which means that gauge fields zero mode ρ_0 can be localized on the brane with generalized dynamic. To investigate the massive modes of the gauge fields, we substitute Eq. (38) into Eq. (36) and obtain a Schrödinger-like equation given by

$$\left\{-\frac{d^2}{dz^2}+V(z)\right\}\tilde{\rho}(z)=m^2\tilde{\rho}(z)\,, \tag{40}$$

where the effective potentia $V(z)$ is

$$V(z)=p\left(\partial_z^2 A+p(\partial_z A)^2\right)\,. \tag{41}$$

The equation above can be written in the form corresponding to supersymmetric quantum mechanic as

$$(\partial_z+p\partial_z A(z))(\partial_z-p\partial_z A(z))\tilde{\rho}(z)=-m^2\tilde{\rho}(z)\,. \tag{42}$$

From this equation, we can exclude the existence of tachyonic modes in the spectrum. We must consider asymptotic behavior of the potential $V(z)$ to investigate the existence of gaps in the spectrum. However, it is not possible to obtain an analytical expression for the function $A(z)$. This means that the potential must be studied numerically. In the present paper, we are

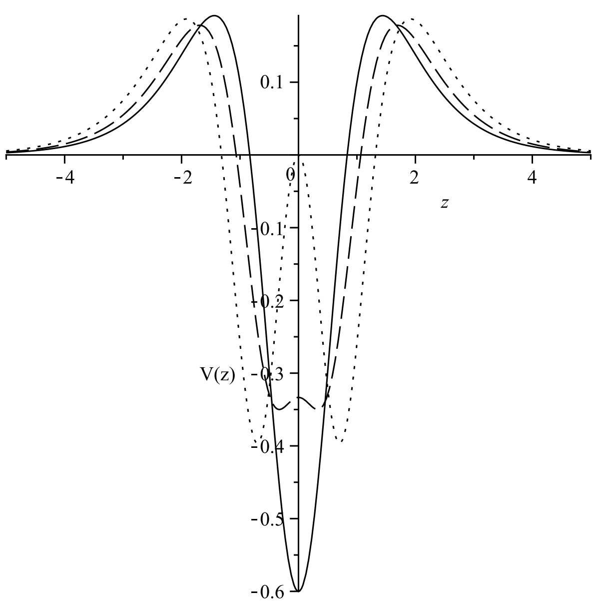

Fig. 2. The shape of potential $V(z)$ for $b=\frac{2}{3}$, $c=1$, $p=1$ and $\alpha=0.01$ (solid line), $\alpha=0.5$ (dashed line), $\alpha=1$ (dotted line).

more interested in the influence of α, so we plot the shape of the potential for various values of α in Fig. 2. From this figure we can see that $V(z)$ is a modified volcano potential with a single well for $\alpha = 0.1$ and two wells for $\alpha > 0.5$ and the potentials have the same asymptotic behavior when $z \to \infty$. Also, for $\alpha > 0.5$ the minimum at $z = 0$ splits into two minima and the distance between the two minima of $V(z)$ increases with α. The parameters b and p affect the depth of the potential and with the increasing of the p and b, the potential well gets deeper. Also, the values of the two maxima of the potential increase by increasing of p and b. Since $V(z) \to 0$ when $z \to \pm\infty$, the potential for gauge fields provides no mass gap to separate the zero mode from the excited KK modes. This means that the spectrum of the vector gauge field consists of a bound zero mode and a series of the continuous massive modes.

4.2. Localization of the Kalb–Ramond field

Here, we study the localization of the Kalb–Ramond field on the brane with generalized dynamic. In Ref. [18], it was found that there is no localized tensorial zero mode with the usual thick brane background without the coupling between the background scalar and the gauge tensor field. So we use the smearing out $G(\phi)$ function in the tensor gauge fields action to localize KR zero mode field on the brane as done in [11]. Hence the action of this field is

$$S = -\frac{1}{12}\int d^5x\sqrt{g}G(\phi)H_{MNL}H^{MNL}\,, \tag{43}$$

where $H_{MNL} = \partial_M B_{NL} + \partial_N B_{LM} + \partial_L B_{MN}$ is the field strength for the KR field and $M, N, L = 0, 1, 2, 3, 4$. From the action (43) and the metric (2), the equation of motion for the KR field can be expressed as

$$e^{2A}G(\phi)\partial_\mu H^{\mu\eta\theta} - \partial_y\left(G(\phi)H^{y\eta\theta}\right) = 0\,. \tag{44}$$

With the gauge choice $B^{\mu y} = 0$, $\partial_\mu B^{\mu\nu} = 0$ and decomposing the field as $B^{\eta\theta} = \Sigma_n h^{\eta\theta}U_n(y)$, we have

$$\frac{\dot{G}(\phi)}{G(\phi)}\frac{\partial U_n(y)}{\partial y} + \frac{\partial^2 U_n(y)}{\partial^2 y} + m^2 e^{-2A(y)}U_n(y) = 0\,. \tag{45}$$

It is obvious that $U_0 =$ constant is a zero mode solution for the KR field. Now, let us substitute this constant solution into the effective action for the tensor field (43) which leads directly to

$$S = -\frac{1}{12}U_0^2\int_{-\infty}^{+\infty} e^{-2A}G(\phi)dy\int d^4x h_{\mu\eta\theta}h^{\mu\eta\vartheta}\,. \tag{46}$$

The condition of having a localized KR field on the brane requires that the part dependent on the extra dimension in the action above should be finite. By considering the smearing out function $G(\phi)$ given by Eq. (39), we obtain

$$\int e^{-2A}G(\phi) = \int \mathrm{sech}^{(2p-3)b}(cy)e^{\frac{3}{4}\alpha b^2c^2(2p-3)\tanh^2(cy)}dy\,, \tag{47}$$

which is convergent for finite product of $\alpha b^2 c^2$ and $p > \frac{3}{2}$. This means that the localization of the zero mode for the KR field can be achieved. We also find that for a specific coupling given by Eq. (39), the normalizable zero modes of the vector and tensor gauge field can be localized on the brane with nonstandard kinetic term under the condition $p > \frac{3}{2}$. In order to analyze the Kaluza–Klein massive spectrum, we return to Eq. (45) with $m \neq 0$ and transform it into a Schrödinger-like equation as

$$\left\{-\frac{d^2}{dz^2} + V(z)\right\}\tilde{U}(z) = m^2\tilde{U}(z)\,, \tag{48}$$

where we have used $U(y) = e^{-\acute{p}}\tilde{U}(z)$ with $\acute{p} = p - 1$. Also the effective potential $V(z)$ is given by

$$V(z) = \acute{p}\left(\partial_z^2 A + \acute{p}(\partial_z A)^2\right)\,. \tag{49}$$

This potential is very similar to the one given in Eq. (41). Hence we can conclude the $V(z)$ is a volcano-like potential which provides no mass gap to separate the KR zero mode from the excited KK modes. Also the equation above can be written in the form corresponding to supersymmetric quantum mechanic as

$$\left(\partial_z + \acute{p}\partial_z A(z)\right)\left(\partial_z - \acute{p}\partial_z A(z)\right)\tilde{U}(z) = -m^2\tilde{U}(z)\,. \tag{50}$$

From this equation, we can exclude the existence of tachyonic modes in the spectrum.

5. Conclusions

In this paper, we investigated the localization of scalar and gauge field on a thick brane with nonstandard kinetic terms $L = K(X) - V(\phi)$, where $K = X + \alpha|X|X$ (type-I model in [4]). First, we showed that massless spin 0 scalar fields can be localized on the brane. Next, we considered thin brane limit and we found that as the value of α decreases the zero-mode of the scalar field gets more and more localized on the brane with generalized dynamic. Furthermore, by considering thin brane limit, we find that the potential of KK modes in the corresponding Schrödinger equation is a volcano-like potential which means the potential provides no mass gap to separate the scalar zero mode from KK modes. We analyzed several aspects of the localization properties of the vector and tensor gauge field on a brane generated by one scalar field with generalized dynamic. In order to localize the gauge fields, we applied a mechanism which is proposed in Ref. [11]. We used a smearing out function as $e^{(2p-1)A(z)}$ in the gauge field Lagrangian which simultaneously leads to vector and tensor gauge field zero mode localization on the brane with nonstandard kinetic term under condition $p > \frac{3}{2}$. Also, we find the massive spectrum of the KR and vector gauge field satisfies a Schrödinger-like equation. For both tensor and vector gauge field, the effective potential is volcano-like which means the potential for the gauge field provides no mass gap to separate the zero mode from the excited KK modes. In this work, our choice for the braneworld model was type-I model of Ref. [4]. It would be interesting to study the matter field localization on the brane configuration which is type-II model of Ref. [4].

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