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Some model-independent phenomenological consequences of flexible braneworlds

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Abstract

In this work, we will review the main properties of braneworld models with low tension. Starting from very general principles, it is possible to obtain an effective action for the relevant degrees of freedom at low energies (branons). Using the cross sections for high-energy processes involving branons, we set bounds on the different parameters appearing in these models. We also show that branons provide a WIMP candidate for dark matter in a natural way. We consider cosmological constraints on its thermal and non-thermal relic abundances. We derive direct detection limits and compare those limits with the preferred parameter region in the case in which the EGRET excess in the diffuse galactic gamma rays is due to dark matter annihilation. Finally, we will discuss the constraints coming from the precision tests of the standard model and the muon anomalous magnetic moment.

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(Some figures in this article are in colour only in the electronic version)

1. Introduction

In the braneworld scenario [1], our universe is understood as a three-brane embedded in a higher dimensional space. If the extra space is compact, the large radius R of the n extra dimensions implies that the fundamental gravity scale in $D = 4 + n$ dimensions M_D can be much lower than the Planck scale M_P. Indeed $M_P^2 = M_D^{2+n} R^n$. Thus, typically it is possible to have weak-scale gravity $M_D \simeq 1$ TeV for $n = 2$ and $R \sim 1$ mm or $n = 7$ and $R \sim 10$ Fermi. However, apart from M_D, there is another fundamental scale in these models which is given by the finite tension of the brane $\tau = f^4$. The relative size of these two scales determines two completely different regimes. In the rigid-brane limit ($f \gg M_D$), the new relevant degrees of freedom are the Kaluza–Klein (KK) modes of fields propagating in the extra dimensions. In the flexible-brane limit ($f \ll M_D$), the KK modes decouple from the standard model (SM) fields [2, 3] and the only new relevant degrees of freedom are those corresponding to the brane fluctuations (branons). In this work, we will review the main phenomenological consequences of braneworld models in this latter case.

2. Branons versus Kaluza–Klein gravitons

Let us consider our four-dimensional spacetime M_4 to be embedded in a D-dimensional bulk space whose coordinates will be denoted by (x^μ, y^m), where x^μ, with $\mu = 0, 1, 2, 3$, correspond to the ordinary four-dimensional spacetime and y^m, with $m = 4, 5, \ldots, D-1$, are coordinates of the compact extra space.

From the point of view of particle physics, the first new effects of braneworlds are related to the KK mode expansion of the bulk gravitational field:

$$g_{\mu\nu}(x, \vec{y}) = \sum_{\vec{k}} g^{\vec{k}}_{\mu\nu}(x)\, \mathrm{e}^{\mathrm{i}\vec{k}\cdot\vec{y}/R} \tag{1}$$

where a toroidal compactification has been assumed for simplicity and $\vec{k}$ is an n-dimensional vector with components $k^m = 0, 1, 2, \ldots$. Therefore, a bulk graviton can be understood as a KK tower of four-dimensional massive gravitons with masses of the order of k/R with k being any natural number (for the $n = 1$ case) so that the distance in the mass spectrum between two consecutive KK gravitons is of the order of $1/R$. This means in particular that the KK graviton spectrum can be considered as approximately continuous for large extra dimensions. In principle, we expect two kinds of effects from the KK graviton tower, namely graviton production and virtual effects on other particle production or observables (see for instance [4] and references therein). The rates for the different processes can be computed by linearizing the bulk gravitational field and by coupling the graviton field to the SM energy–momentum tensor $T^{\mu\nu}_{\mathrm{SM}}$. Then expanding the gravitational field in terms of the KK modes one finds the corresponding Feynman rules.

However, apart from the existence of KK modes, the fact that the brane is a physical object with finite tension requires the existence of new fields which parametrize the brane fluctuations [5]. For simplicity, we will assume that the bulk metric tensor takes the following form:

$$\mathrm{d}s^2 = \tilde{g}_{\mu\nu}(x) W(y)\, \mathrm{d}x^\mu\, \mathrm{d}x^\nu - g'_{mn}(y)\, \mathrm{d}y^m\, \mathrm{d}y^n \tag{2}$$

where the warp factor is normalized as $W(0) = 1$.

Working in the probe-brane approximation, our three-brane universe is moving in the background metric given by (2) which is not perturbed by its presence. The position of the brane in the bulk can be parametrized as $Y^M = (x^\mu, Y^m(x))$, and we assume for simplicity that the ground state of the brane corresponds to $Y^m(x) = 0$.

In the simplest case in which the metric is not warped along the extra dimensions, i.e. $W(y) = 1$, the transverse brane fluctuations are massless and they can be parametrized by the Goldstone bosons fields $\pi^\alpha(x)$, $\alpha = 4, 5, \ldots, D-1$ (branons). In that case we can choose the y coordinates so that the branon fields are proportional to the extra-space coordinates: $\pi^\alpha(x) = f^2 \delta^\alpha_m Y^m(x)$, where the proportionality constant is related to the brane tension $\tau = f^4$ [5].

In the general case, the curvature generated by the warp factor explicitly breaks the traslational invariance in the extra space. Therefore, branons acquire a mass matrix which is given precisely by the bulk Riemann tensor evaluated at the brane position:

$$M^2_{\alpha\beta} = \tilde{g}^{\mu\nu} R_{\mu\alpha\nu\beta}|_{y=0}. \tag{3}$$

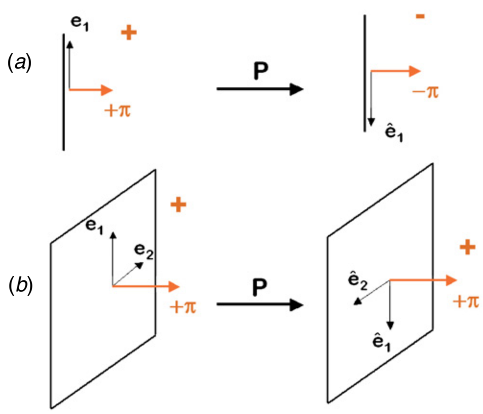

Figure 1. Parity transformation on the brane for (*a*) odd-dimensional brane. Branons fields change sign (pseudoscalars). (*b*) Even-dimensional brane (scalar branons).

The fact that the brane can fluctuate implies that the actual metric on the brane is no longer given by $\tilde{g}_{\mu\nu}$, but by the induced metric which includes the effect of warping through the mass matrix:

$$g_{\mu\nu}(x,\pi) = \tilde{g}_{\mu\nu}(x)\left(1 + \frac{M^2_{\alpha\beta}\pi^\alpha\pi^\beta}{4f^4}\right) - \frac{1}{f^4}\partial_\mu\pi^\alpha\partial_\nu\pi^\alpha + \mathcal{O}(\pi^4). \tag{4}$$

The dynamics of branons can be obtained from the Nambu-Goto action by introducing the above expansion. In addition, it is also possible to get their couplings to the ordinary particles just by replacing the spacetime by the induced metric in the standard model action. Thus we get, up to quadratic terms in the branon fields,

$$\begin{aligned} S_{\mathrm{Br}} = \int_{M_4} \mathrm{d}^4x\sqrt{\tilde{g}}\Bigg[&\frac{1}{2}\left(\tilde{g}^{\mu\nu}\partial_\mu\pi^\alpha\partial_\nu\pi^\alpha - M^2_{\alpha\beta}\pi^\alpha\pi^\beta\right) \\ &+ \frac{1}{8f^4}\left(4\partial_\mu\pi^\alpha\partial_\nu\pi^\alpha - M^2_{\alpha\beta}\pi^\alpha\pi^\beta\tilde{g}_{\mu\nu}\right)T^{\mu\nu}_{\mathrm{SM}}\Bigg]. \end{aligned} \tag{5}$$

We can see that branons interact with the SM particles through their energy–momentum tensor. The couplings are controlled by the brane tension scale f and they are universal very much like those of gravitons. For large f, branons are therefore weakly interacting particles.

The sign of the branon fields is determined by the orientation of the brane submanifold in the bulk space. Under a parity transformation on the brane ($x^i \to -x^i$), the orientation of the brane changes sign provided the ordinary space has an odd number of dimensions, whereas it remains unchanged for even spatial dimensions (see figure 1). In the case in which we are interested with three ordinary spatial dimensions, branons are therefore pseudoscalar particles. If the gravitational sector of the theory respects parity on the brane, then branons always couple to SM particles by pairs, which ensures that they are stable particles. This fact can have important consequences in cosmology as we show below.

3. Phenomenology in colliders

The main processes in colliders in which branons could contribute in a relevant way are the single-photon and single-Z production in e^+e^- colliders (LEP) and the monojet and single-

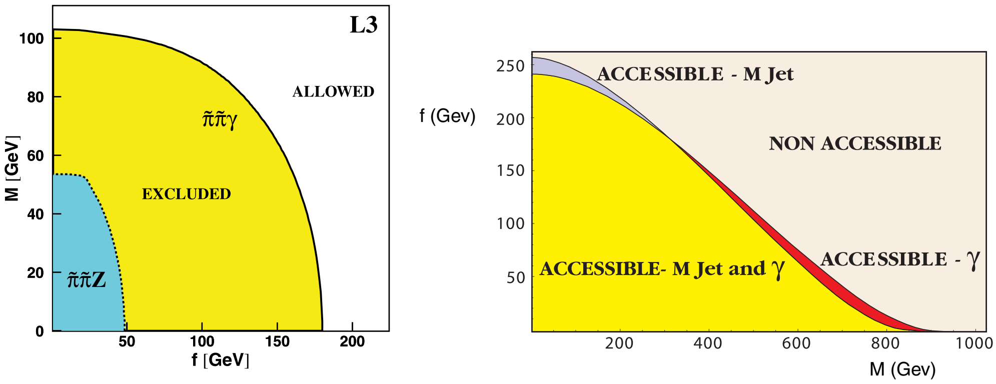

Figure 2. Collider limits on branon parameters from single-photon and single-Z processes at LEP (L3) [6] (left). Limits from monojet and single-photon processes at Tevatron-I [7] (right).

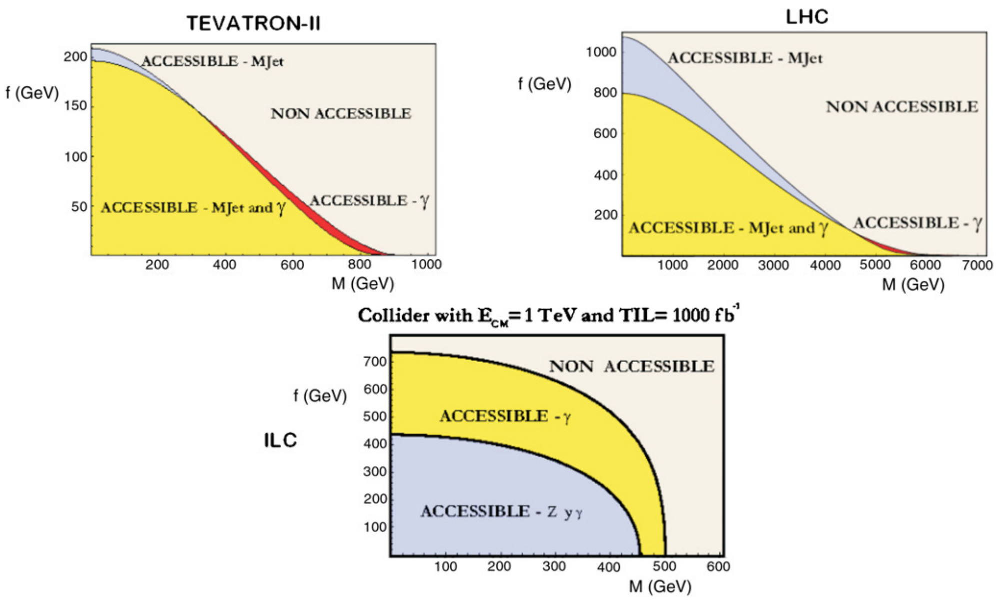

Figure 3. Collider limits on branon parameters from monojet and single-photon processes for Tevatron-II and LHC and single-photon and single-Z processes for the International Linear Collider (ILC) with typical values for the centre of mass energy of 1 TeV and total integrated luminosity of 1000 fb^{-1}.

photon production in hadron colliders such as Tevatron [6, 7]. In those processes, a pair of branons are produced in the final state which are not detected and therefore they appear as missing energy and momentum. We have computed the limits on (f, M) imposed by the absence of any deviation from the SM predictions in such processes. The results are summarized in figure 2.

We have also estimated the prospects for future hadron colliders (Tevatron-II and LHC) and e^+e^- colliders (ILC). In figure 3, we show the accessible regions in the (f, M) plane.

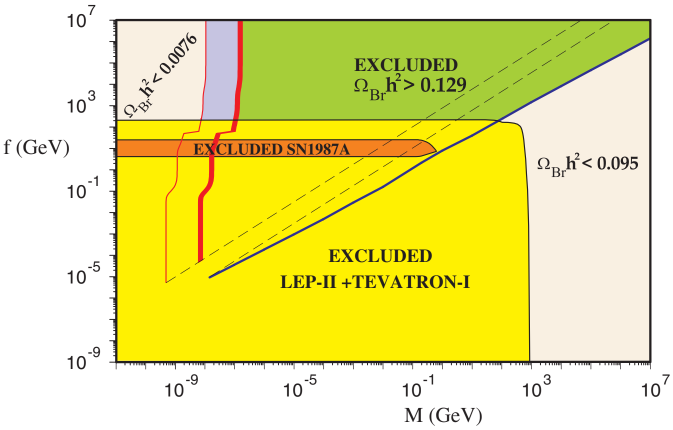

Figure 4. Thermal relic abundance in the f–M plane for a model with one branon of mass M. The two lines on the left correspond to the WMAP first-year measurements: $\Omega_{\mathrm{Br}}h^2 = 0.0076$ and $\Omega_{\mathrm{Br}}h^2 = 0.129$–$0.095$ curves for hot-warm relics, whereas the right line corresponds to the latter limits for cold relics (see [8] for details). The lower area is excluded by single-photon processes at LEP-II [7] together with monojet signal at Tevatron-I [7]. The astrophysical constraints are less restrictive and they mainly come from supernova cooling by branon emission [8].

4. Branons as dark matter

As commented before branons are stable and weakly interacting particles, this makes them natural dark matter candidates [8].

In [8], the thermal relic branon abundance has been calculated in two cases: either relativistic branons at freeze-out (hot-warm) or non-relativistic (cold), and assuming that the evolution of the universe is standard for $T < f$ (see figure 4 for comparison with the WMAP first-year limits).

However, branons can also act as non-thermal dark matter very much in the same way as axions. Indeed, if the maximum temperature reached in the universe is smaller than the branon freeze-out temperature T_f, but larger than the explicit symmetry breaking scale $(Mf^2R)^{1/2}$, then branons can be considered as essentially massless particles decoupled from the rest of matter and radiation. In such a case, there is no reason to expect that the brane initial position Y_0 should coincide with the potential minimum, but in general we would have $Y_0 \simeq R$. Thus as the universe cools down, the brane can start oscillating coherently around the equilibrium position, in a completely analogous way to the axion misalignment mechanism. The energy density of those oscillations scales as that of non-relativistic matter, and can be estimated as [9]

$$\Omega_{\mathrm{Br}}h^2 \simeq \frac{6.5 \times 10^{-20} N}{\mathrm{GeV}^{5/2}} f^4 R^2 M^{1/2}. \qquad (6)$$

Limits on the different parameters in such a case can be found in figure 5.

If branons constitute the dark matter of the universe, they will make up the galactic halo and could also be detected by direct or indirect detection experiments. In the direct case, branons interact with the nuclei in the detector and the nucleus recoil energy can be measured. Several underground experiments set limits on the branon mass and the elastic branon–nucleon cross section at zero momentum transfer σ_n [8] (see figure 6).

Branons could also be detected indirectly: their annihilations in the galactic halo can give rise to pairs of photons or e^+e^- which could be detected by γ-ray telescopes such as MAGIC

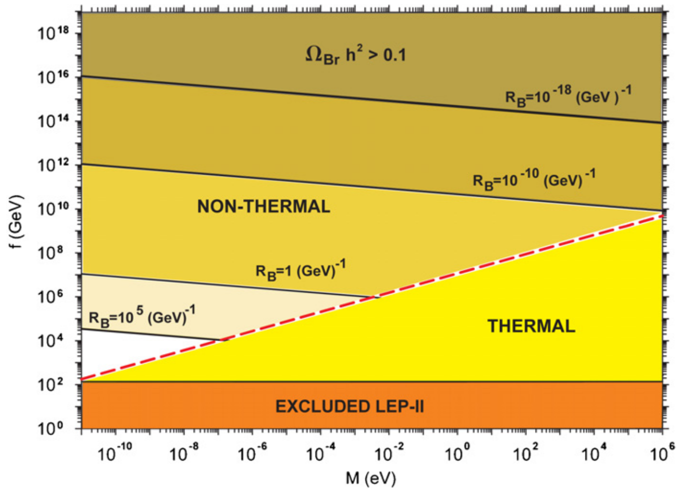

Figure 5. Thermal versus non-thermal branons regions in the f–M plane. The dashed (red) line separates the two regions and corresponds to $T_f \simeq (MM_P)^{1/2}$. The continuous (black) lines correspond to $\Omega_{\rm Br}h^2 \simeq 0.1$ for different values of the radius of the extra dimensions $R = 10^5, 1, 10^{-10}, 10^{-18}$ GeV^{-1}.

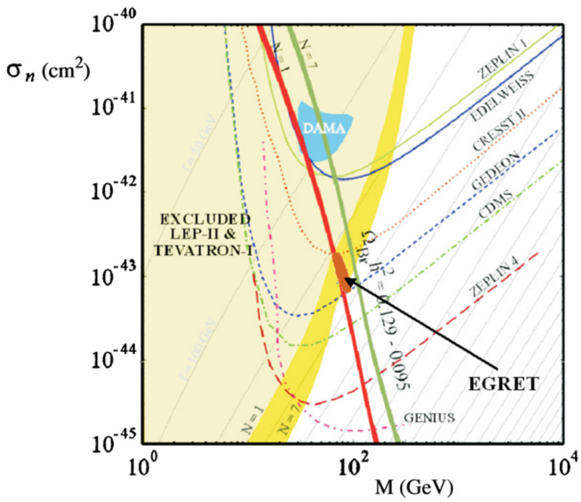

Figure 6. Elastic branon–nucleon cross section σ_n in terms of the branon mass. The thick (red and green) line corresponds to the $\Omega_{\rm Br}h^2 = 0.129$–$0.095$ curve for cold branons in figure 2 from $N = 1$ and $N = 7$ branon species. The shaded areas are the LEP-II and Tevatron-I exclusion regions. The solid lines correspond to the current limits on the spin-independent cross section from direct detection experiments. The discontinuous lines are the projected limits for future experiments. Limits obtained from [10]. The EGRET region corresponds to parameter space for which dark matter annihilation in the galactic halo could account for the gamma-ray excess observed by EGRET [11].

or GLAST or antimatter detectors (see [12] for an estimation of positron and photon fluxes from branon annihilation in AMS). Annihilation of branons trapped in the centre of the sun or the earth can give rise to high-energy neutrinos which could be detectable by high-energy neutrino telescopes.

Dark matter annihilation has been proposed as a possible explanation for the excess of flux above 1 GeV in the EGRET data of the diffuse galactic gamma rays. For that purpose, a WIMP mass around 50–70 GeV would be required. In figure 6, we include the parameter region calculated from [11]

5. Virtual branon effects

Branons can also have phenomenological consequences through virtual effects [3]. Probably the most immediate effect of virtual branons is the suppression of the coupling of SM particles and the KK modes of bulk fields like the graviton. When branon fluctuations are taken into account, this effective coupling is described by the action

$$S_h = \frac{1}{\bar{M}_P}\sum_p \int_{M_4} \mathrm{d}^4x\, h^{(p)}_{\mu\nu}(x) T^{\mu\nu}_{\mathrm{SM}}(x) f_p(\pi) \tag{7}$$

where $\bar{M}_P^2 \equiv M_P^2/4\pi$ is the squared reduced Planck mass. The $f_p(\pi)$ couplings appear due to the fact that the brane is no more sitting at $\pi = 0$, but can fluctuate around this point. Now the branons fields can be integrated out in the usual way to find

$$S_h = \frac{1}{\bar{M}_P}\sum_p \int_{M_4} \mathrm{d}^4x\, h^{(p)}_{\mu\nu}(x) T^{\mu\nu}_{\mathrm{SM}}(x) \langle f_p(\pi)\rangle \tag{8}$$

where the f_p expectation value is given by

$$\langle f_p(\pi)\rangle = \int [\mathrm{d}\pi]\, \mathrm{e}^{\mathrm{i}S^{(2)}_{\mathrm{eff}}[\pi]} f_p(\pi). \tag{9}$$

To compute the path integral above, we need to know the precise form of the $f(\pi)$ functions which depends on the geometry of the extra dimensions. For the simplest case of torus compactification, the effective action becomes

$$S_h = \frac{1}{\bar{M}_P}\sum_{\vec{k}} \int_{M_4} \mathrm{d}^4x\, g_{\vec{k}} h^{(\vec{k})}_{\mu\nu}(x) T^{\mu\nu}_{\mathrm{SM}}(x). \tag{10}$$

In other words, the effect of branon quantum fluctuations amounts to introducing the KK-mode-dependent couplings $g_{\vec{k}}$ which are given for toroidal compactification by

$$g_{\vec{k}} = \exp\left(-\frac{\Lambda^2}{32\pi^2 R^2 f^4}\sum_{m=1}^{N}(k^m)^2\left[1 - \frac{M_m^2}{\Lambda^2}\log\left(\frac{\Lambda^2}{M_m^2}+1\right)\right]\right) \tag{11}$$

where we have introduce an ultraviolet cut-off Λ in order to regularize the divergent integrals. Thus, the coupling of SM matter to higher KK modes is exponentially suppressed. This result was first obtained in [2] for massless branons by using an argument based on normal ordering. In any case, this coupling suppression has very interesting consequences from the phenomenological point of view. It improves the unitarity behaviour of the cross section for producing gravitons from SM particles and, in addition, it solves the problem of the divergences appearing even at the tree level when one considers the KK graviton tower propagators for dimension equal or larger than 2. Moreover whenever we have $v \equiv Rf^2 \ll \Lambda$, KK gravitons decouple from the SM particles, so that at low energies the only braneworld related particles that must be taken into account are branons.

In order to study the effect of virtual branons on the SM particles, it is useful to introduce the SM effective action $\Gamma^{\mathrm{eff}}_{\mathrm{SM}}[\Phi]$ obtained after integrating out the branon fields:

$$\mathrm{e}^{\mathrm{i}\Gamma^{\mathrm{eff}}_{\mathrm{SM}}[\Phi]} = \int [\mathrm{d}\pi]\, \mathrm{e}^{\mathrm{i}S_{\mathrm{SM}}[\Phi,\pi]}. \tag{12}$$

Table 1. Limits from virtual branon searches at colliders (results at the 95% c.l.) The indices a, b, c denote the two-photon, e^+e^- and e^+p (e^-p) channels, respectively. The first four analysis have been performed with real data, whereas the final four are estimations. The first two columns correspond to the centre of mass energy and total luminosity, and the third one corresponds to the lower bound on $f^2/(N^{1/4}\Lambda)$.

Experiment	$\sqrt{s}$ (TeV)	$\mathcal{L}$ (pb^{-1})	$f^2/(N^{1/4}\Lambda)$ (GeV)
HERA[c]	0.3	117	52
Tevatron-I[a,b]	1.8	127	69
LEP-II[a]	0.2	700	59
LEP-II[b]	0.2	700	75
Tevatron-II[a,b]	2.0	2×10^3	83
ILC[b]	0.5	5×10^5	261
ILC[b]	1.0	2×10^5	421
LHC[b]	14	10^5	383

This effective action is in general a non-local and divergent functional of the standard model fields and in order to obtain finite results we have used a cut-off (Λ) regularization procedure as before. The new local divergent terms are organized in powers of the fields so that, at the level of two-point functions, we have [3]

$$\Delta\mathcal{L}^{(1)} = C_1 T^{\mu}_{\mathrm{SM}\mu} \tag{13}$$

where $C_1 = -N\Lambda^4/(16(4\pi)^2 f^4)$ with N being the number of branon species. This term is proportional to the trace of the SM energy–momentum tensor and accordingly amounts to a renormalization of SM masses. To second order, we get

$$\Delta\mathcal{L}^{(2)} = W_1 T^{\mu\nu}_{\mathrm{SM}} T^{\mathrm{SM}}_{\mu\nu} + W_2 T^{\mu}_{\mathrm{SM}\mu} T^{\nu}_{\mathrm{SM}\nu} \tag{14}$$

where $W_1 = 2W_2 = N\Lambda^4/(96(4\pi)^2 f^8)$. These counterterms introduce new interaction vertices which are not present in the SM Lagrangian. The most relevant ones could be the four-fermion interactions or the fermion pair annihilation into two gauge bosons. We have used the data coming from HERA, Tevatron and LEP on this kind of processes in order to set bounds on the parameter combination $f^2/(\Lambda N^{1/4})$. The results are shown in table 1, where it is also possible to find the prospects for the future colliders mentioned above.

As we have commented above, Λ is the cut-off which limits the validity of the effective description of branon and SM dynamics. This new parameter appears when dealing with branon radiative corrections since the Lagrangian in (5) is not renormalizable. A one-loop calculation with the new effective four-fermion vertices coming from (14) is equivalent to a two-loop computation with the Lagrangian in (5), and allows us to obtain the contribution of branons to the muon anomalous magnetic moment [13]:

$$\delta a_\mu \simeq \frac{5m_\mu^2}{114(4\pi)^4}\frac{N\Lambda^6}{f^8}. \tag{15}$$

We can observe that the correction has the right sign and that it is thus possible to improve the agreement with the experimental value. In fact, the difference between the experimental and the SM prediction [14, 15] measured at the 2.7σ level by the Brookhaven g-2 Collaboration is $\delta a_\mu \equiv a_\mu(\exp) - a_\mu(\mathrm{SM}) = (23.4 \pm 9.1) \times 10^{-10}$. Thus, we can estimate the preferred parameter region for branon physics as

$$6.0\ \mathrm{GeV} \gtrsim \frac{f^4}{N^{1/2}\Lambda^3} \gtrsim 2.2\ \mathrm{GeV}(95\%\ \mathrm{c.l.}). \tag{16}$$

Another limitation to the branon parameters could be obtained from electroweak precision measurements, which use to be very useful to constrain models of new physics. The so-called oblique corrections (the ones corresponding to the W, Z and γ two-point functions) used to be described in terms of the S, T, U or the ϵ_1, ϵ_2 and ϵ_3 parameters. The experimental values obtained by LEP are consistent with the SM prediction for a light Higgs $m_H \leqslant 237$ GeV at 95% c.l. In principle, it is necessary to know this parameter in order to put constraints on new physics, but one can calculate the disfavoured regions of parameters in order to avoid fine tunings. We can estimate this area by performing a computation of the parameter $\bar{\epsilon} \equiv \delta M_W^2 / M_W^2 - \delta M_Z^2 / M_Z^2$, in a similar way as it was done for the first-order correction coming from the Kaluza–Klein gravitons in the ADD models for rigid branes [16]. The experimental value of $\bar{\epsilon}$ obtained from LEP is $\bar{\epsilon} = (1.27 \pm 0.16) \times 10^{-2}$. The theoretical uncertainties are one order of magnitude smaller and therefore we can estimate the constraints for the branon contribution at 95% c.l. as $|\delta\bar{\epsilon}| \gtrsim 3.2 \times 10^{-3}$ with the result [3]

$$\frac{f^4}{N^{1/2}\Lambda^3} \gtrsim 3.1 \text{ GeV } (95\% \text{ c.l.}). \tag{17}$$

6. Conclusions

In this work, we have reviewed the main phenomenological consequences of flexible braneworlds. The new effects can be described in a model-independent way by means of the low-energy effective action for branons which depends essentially on only two parameters: the brane tension scale f and the branon mass M. We have computed the limits on those parameters from LEP and Tevatron-I data, and the prospects for detection in future colliders.

We have shown that branons are natural candidates for dark matter both as thermal and non-thermal relics in different regions of the parameter space. We have obtained the cosmological limits from WMAP observations and also studied the prospects for detection of branons in the galactic halo from direct or indirect searches experiments.

Finally, we have considered additional constraints on branon physics from virtual effects and those coming from electroweak precision observables. The new two-loop diagrams involving branons are shown to give a non-vanishing contribution to the muon anomalous magnetic moment, which could explain the observed difference between the SM prediction and the measurements of the Brookhaven g-2 Collaboration.

Acknowledgments

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