Brane-Production and the Neutrino-Nucleon cross section at Ultra High Energies in Low Scale Gravity Models

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Abstract: The origin of the ultra high energy cosmic ray (UHECR) showers has remained as a mystery among particle physicists and astrophysicists. In low scale gravity models, where the neutrino-nucleon cross section rises to typical hadronic values at energies above $10^{20}$ eV, the neutrino becomes a candidate for the primary that initiates these showers. We calculate the neutrino-nucleon cross section at ultra high energies by assuming that it is dominated by the production of p-branes. We show, using a generalized Randall-Sundrum model, that the neutrino-nucleon cross-section at neutrino energies of $10^{11}$ GeV is of the order of 100 mb, which is required for explaining UHECR events. Similar result also follows in other models such as the Lykken-Randall model.
Low scale gravity models [1, 2] lead to the possibility that the neutrino-nucleon cross section at ultra high energies may be several orders of magnitude larger than what is predicted by the Standard Model. At center of mass energies $\sqrt{s}$ less than the scale of quantum gravity $M_*$, which we assume is of order 1 TeV, the cross section can be calculated using the perturbative Feynman rules [3]. However at higher energies, such that $\sqrt{s} > M_*$, there does not exist any reliable procedure to calculate the cross section. In applications to cosmic rays we are interested in cross section at center of mass energies of the order of $10^{12}$ GeV$^2$, where the perturbative approach is certainly not applicable. Several authors [4, 5, 6, 7, 8, 9, 10, 11] have obtained various estimates by using different models for the cross section in this energy regime. In Ref. [5] we assumed two models for the $s$ dependence of the cross section, $\sigma \sim s, s^2$, and found that neutrino-nucleon cross sections are of the order of one to several hundred mb with the precise value dependent on the model used and on the choice of $M_*$. The larger value of the cross section is obtained using the $s^2$ model. Such large values of the neutrino-nucleon cross section are very interesting since then the neutrino becomes a candidate [5, 6] for explaining the puzzling cosmic ray events which appear to violate the GZK bound [12, 13, 14, 15]. The idea that neutrino-nucleon cross section grows large at GZK energies is very old [16]. Large extra dimension, low scale gravity models [1, 2] have provided new impetus to this idea [5, 6, 7, 17, 18, 19]. Even if the cross section is not large enough to explain the GZK events, cosmic ray observations can be used to test this hypothesis and to put stringent bounds on these models.

There exist several proposals for calculating the ultra high energy scattering cross sections including the gravitational contribution. In Ref. [4, 9, 20, 21] the authors proposed an eikonal model for the cross section, which may be applicable for very large $s$ but for momentum transfer $-t < M_*^2$. For larger values of the momentum transfer, Ref. [4, 22, 23, 24, 25] have proposed that black hole production will set in. This proposal has generated great interest with several studies investigating the production rate of black holes at colliders [24, 26, 27, 28, 29, 30, 31] and cosmic ray collisions [10, 11, 32, 33, 34]. Using the eikonal model [9] and the black hole production rate [10], it is found that the neutrino-nucleon cross section is of the order of 0.1 mb at $s \approx 10^{12}$ GeV$^2$. This cross section is also large enough such that neutrinos could generate horizontal showers, which can easily be seen in future observatories.
Black hole production cross section is estimated by assuming that it is approximately equal to the geometric area of the black hole produced. A collision between partons \( i \) and \( j \) with center of mass energy \( \sqrt{s} \) is assumed to produce a black hole of mass \( M_{\text{BH}} \approx \sqrt{s} \) and hence the parton level cross section \( \hat{\sigma}(ij \rightarrow \text{BH}) \) is equal to \( \pi R_S^2 \) where \( R_S \) is the Schwarzschild radius of a black hole of mass \( M_{\text{BH}} \). This geometric cross section has been criticized by Voloshin [35] on the grounds that a high energy collision with partons \( i \) and \( j \) will have a very large amplitude to radiate, and the amplitude to produce an isolated black hole will be exponentially damped. At our current stage of understanding of quantum gravity, the various models proposed must be considered as highly speculative.

Recently Ref. [36] has proposed that besides black holes, high energy collisions will also produce branes. They argue that in certain cases the production cross section of p-branes of mass \( M_p \approx \sqrt{s} \) is much higher than that of black holes of the same mass. The precise value of the cross section depends on the dimensionality of the brane and the size of the extra dimensions. The largest cross section for a p-brane is obtained when the brane is completely wrapped on the small-size extra dimensions where it is assumed that there exist \( m \) extra dimensions compactified on length scale of order \( L_s < M_s^{-1} \) and the remaining \( n - m \) extra dimensions are compactified on length scale of order \( L' >> M_s^{-1} \). The ratio \( \Sigma(s; n, m, p \leq m) \) of the cross section of the p-brane to black hole production in this case is given by [36]

\[
\Sigma(s; n, m, p \leq m) \approx \left( \frac{L}{L_s} \right)^{-\frac{2p}{n-p+1}} \frac{\gamma(n, p)^2}{\gamma(n, 0)^2} \left( \frac{s}{s_*} \right)^{-\frac{m-1}{n+1}} \tag{1}
\]

where \( n \) is the total number of extra dimensions, \( L_s = M_s^{-1} \),

\[
\gamma(n, p) = \left[ \frac{8\Gamma \left( \frac{n+3-p}{2} \right)}{(2+n)\sqrt{(1-\frac{p}{n+2})/(p+1)}} \right]^{\frac{1}{n+1}} \tag{2}
\]

and \( w = 1/[1-p/(n+1)] \). The branes, once produced, will decay producing gravitons and matter fields.

We calculate the neutrino-nucleon total cross section \( \sigma_{\nu p} \) at ultra high energies assuming that the cross section is dominated by brane production. We perform this calculation within the framework of a generalized Randall-Sundrum (RS) model [2] and the model proposed by Lykken and Randall in
Ref. [37] such that these models have an arbitrary number of compact extra dimensions instead of the single compact dimension assumed in Ref. [2, 37]. Besides the single compact dimension which involves the warp factor, the remaining \( m \) compact dimensions are assumed to be factorizable. These \( m \) extra dimensions are compactified on a \( m \)-torus of radius \( R \). The generalized RS model, therefore, consists of two 3-branes with opposite tension situated at orbifold points \( y = 0, y_c \) in a \((4+m+1)\)-dimensional space-time. The metric can be written as

\[
ds_{D+1}^2 = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + R^2 \delta_{ab} d\theta_a d\theta_b + dy^2
\]

where \( \mu, \nu = 0, ..., 3 \), \( a, b = 1, ..., m \), \( \theta_a \) are the compacted dimension with range \([0,2\pi]\) and \( y \) is the warped dimension with range \([-y_c, y_c]\). The Einstein equations can be written as

\[
G_{AB} = \frac{1}{4M^{m+3}} \left( T^{(B)}_{AB} + T^{(br)}_{AB} \right)
\]

The components of the bulk energy-momentum tensor are assumed to be

\[
T^{(B)}_{\nu\mu} = \delta^\nu_\mu \Lambda_0, \quad T^{(B)}_{y\mu} = \Lambda_0, \quad T^{(B)}_{a\mu} = \Lambda_0 \delta_a^\mu
\]

The components of the brane energy-momentum tensor are assumed to be

\[
T^{(br)}_{\nu,i} = \delta(y - y_i) v_0^i \delta^\nu_\nu, \quad T^{(br)}_{y,i} = 0, \quad T^{(br)}_{n,i} = \delta(y - y_i) v_0^i \delta_n^n,
\]

where \( i = 1, 2 \), and all other non-diagonal components are taken to be zero. Inhomogeneities in the energy-momentum tensors and cosmological constants can be due to different contributions to casimir energies in the different direction of space-time [38, 39]. In the absence of a four dimensional cosmological constant, the \((\mu, \nu)\)-component of the Einstein’s equation gives

\[
-3\sigma''(y) + 6\sigma'^2(y) = -\frac{\Lambda_0}{4M^{m+3}} - \sum_i \frac{v_0^i}{4M^{m+3}} \delta(y - y_i)
\]

The \((y,y)\)-component of the Einstein’s equation is given by

\[
6\sigma'(y)^2 = -\frac{\Lambda_0}{4M^{m+3}}
\]
The \((\theta, \theta)\)-component of the Einstein’s equation is given by

\[
-4\sigma''(y) + 10\sigma'^2(y) = -\frac{\Lambda_\theta}{4M^{m+3}} - \sum_i \frac{v_i^0}{4M^{m+3}} \delta(y - y_i) \tag{7}
\]

From Eq.(5,6,7), using the orbifold symmetry along \(y\)-direction, we find

\[
\sigma(y) = \sqrt{\Lambda_0 24 M^{m+3}} (|y| - y_c) \tag{8}
\]

and \(\Lambda_\theta = \frac{5}{3} \Lambda_0, \quad v_\theta^i = \frac{4}{3} v_i^0\). This solution is valid only if brane tensions and cosmological constants are related as follows

\[
\Lambda_0 = -24 M^{m+3} k^2, \quad v_0^1 = -v_0^2 = 24 M^{m+3} k \tag{9}
\]

The four-dimensional reduced Planck mass \(M_{pl}\) can be written as

\[
M_{pl}^2 = M^{(m+3)} \exp \left( 2ky_c \frac{(2\pi R)^m}{k} \left[ 1 - \exp(-2ky_c) \right] \right) \tag{10}
\]

where \(M = \frac{M_*}{(32\pi)^{(m+3)}}\). If we demand \(ky_c \simeq 35, M_* = 1 \text{TeV} \) and \(R < (\text{TeV})^{-1}\) we find that \(k < M_*\). We can take \(y_c\) of order \(\text{fm}\) without violating the existing observational constraints from astrophysics and accelerator physics and then we find \(k\) is of order \(\text{GeV}\). We assume that some unknown mechanism stabilizes the radius of extra dimensions.

We next determine the spectrum of the KK-modes of this theory in the low energy limit. Decomposition of KK states can be carried out by considering the following graviton perturbations:

\[
ds^2 = e^{-2\sigma} [\eta_{\mu\nu} + h_{\mu\nu}(x, y, \theta)] + dy^2 + R^2 \delta_{ij} d\theta_i d\theta_j \tag{11}
\]

where \(R\) is the radius of compact extra dimensions and \(\sigma = k(|y| - y_c)\). Here we are not considering fluctuations around extra dimensions. We expand the linear fluctuations around the flat 4-dimensional metric in a complete set of radial wavefunctions and fourier modes:

\[
h_{\mu\nu}(x, y, \theta) = \sum_{n,l} h_{\mu\nu}^{(n,l)}(x) \phi_{(n,l)}(y) e^{il\theta} \tag{12}
\]
Working in a gauge $\eta^{\alpha\beta}\partial_\alpha h_{\beta\gamma} = \eta^{\alpha\beta}h^{(n,l)}_{\alpha\beta} = 0$ and $\partial^2 h^{(n,l)}_{\mu\nu}(x) = m_n^2 h^{(n,l)}_{\mu\nu}(x)$, the differential equation for radial part is given by

$$\partial_y \left( e^{-4\sigma} \partial_y \phi_{n,l}(y) \right) = -e^{-2\sigma} \left( m_n^2 - \frac{l^2}{R^2} e^{-2\sigma} \right) \phi_{n,l}(y)$$

(13)

with normalization condition $(2\pi R)^m \int_{y_c}^{y_c} dy e^{-2\sigma} \phi_n^* \phi_{n'} = \delta_{nn'}$. After changing variables to $z_n = \frac{M_n}{k} e^{\sigma}$ and $f_{n,l} = e^{-2\sigma} \phi_{n,l}$, Eq.(13) can be written as (for $y \neq 0, y_c$)

$$z_n^2 \frac{d^2 f_{n,l}}{dz_n^2} + z_n \frac{df_{n,l}}{dz_n} + \left[ z_n^2 - \left( 4 + \frac{l^2}{k^2 R^2} \right) \right] f_{n,l} = 0$$

(14)

In the low energy limit we can set $l=0$ since these modes will have energy larger than $1/R$. Then in this limit the theory reduces to RS model except for the fact that the size of the warped extra dimension is taken to be of order $1$ fm which is about three orders of magnitude larger than the fundamental length scale of $(1 TeV)^{-1}$. The masses of these KK-modes are quantized in mass scale of the order of GeV and their couplings to the standard model fields which lives on the brane at $y = y_c$ are the same as in RS case. As there are a large number of modes present, we can ignore the few low lying modes and the sum over the remaining modes can be done by changing the summation to an integration as in the ADD case.

We also consider another model proposed by Lykken and Randall [37] in which the size of the extra dimension can be infinite. The model was originally proposed with only one extra dimension. The model involves three branes, one located at $y = 0$, second at $y = \infty$ and the third, which is the physical brane, located at some position $y = y_0$ such that $0 < y_0 < \infty$. We generalize this model by introducing $m$ additional compactified dimensions in exact analogy to the above construction in the case of RS model. We find that the low energy predictions of the model remain unchanged due to the addition of $m$ compact extra dimensions. Hence we find that we can extend both the RS and the Lykken-Randall model such that they contain $m$ extra dimensions compactified on a length scale of the order of inverse TeV and one large extra dimension of length scale larger than (or order of) 1 fm. We can now construct p-brane solutions in these models such that the brane is completely wrapped around the $m$ extra compact dimensions.

The parton level brane production cross section $\hat{\sigma}_{brane}$ for p-brane of mass
\[ M_p = \sqrt{s} \] is given by,
\[ \hat{\sigma}_{\text{brane}}(\hat{s}) = \Sigma(\hat{s}; n, m, p \leq m)\hat{\sigma}_{\text{BH}}(\hat{s}) \] (15)
where \( \hat{\sigma}_{\text{BH}} \) is the production cross section for black hole of mass \( M_{BH} = \sqrt{\hat{s}} \).

The black hole production cross section, \( \hat{\sigma}_{\text{BH}} \), is given by \([23, 24]\)
\[ \hat{\sigma}_{\text{BH}} \approx \pi r_S^2 \] (16)
where \( r_S \) is the Schwarzschild radius of a \( 4 + n \) dimensional black hole,
\[ r_S = \frac{1}{\sqrt{\pi}} \left( \frac{M_{BH}}{M_*} \right)^{\frac{n}{2}} \left[ \frac{8\Gamma(\frac{3+n}{2})}{2 + n} \right]^{\frac{1}{1+n}} \] (17)

The black hole and the brane production processes are expected to give dominant contributions when \( s >> M_*^2 \).

The total cross section can be computed using
\[ \sigma(\nu N \rightarrow \text{brane}) = \sum_i \int_{x_{\text{min}}}^{1} dx \hat{\sigma}_{\text{brane},i}(xs)f_i(x, Q), \] (18)
where \( f_i \) is the parton distribution function corresponding to the \( i \)th parton, \( x_{\text{min}} = M_p^2/s \) and \( Q \) is the typical momentum scale of the collision which is taken to be the brane mass. We use the CTEQ parton distributions [40] for our calculation. Since the parton distributions are not known beyond \( Q = 10 \) TeV, we set \( Q \) equal to this value if the brane mass exceeds 10 TeV. The minimum value of \( x \) is obtained by assuming that branes are produced with masses \( M_p \) greater than \( M_* \).

In fig. 1 we plot the neutrino-nucleon cross section for some representative choice of parameters, \( p \) and \( n \). The ratio of length scales \( L/L_\ast \) has been set equal to 0.25. We find that for \( n = 7 \) and \( p = 6 \) the cross section rises to roughly 50 mb at neutrino energy of \( 10^{11} \) GeV. As \( p \) becomes smaller the cross section is much smaller. Similarly the cross section is considerably reduced as we lower \( n \). The cross section is also found to be very sensitive to the precise value of \( L/L_\ast \). The dependence of total cross section on \( L/L_\ast \) is shown in fig. 2 for some representative choice of parameters, \( p \) and \( n \). The neutrino energy has been taken to be \( 10^{11} \) GeV for this plot.

It is clear from fig. 1 and fig. 2 that there exists a large range of parameter space where the neutrino-nucleon cross section is of the order of 100 mb or
Figure 1: The neutrino-nucleon cross section $\sigma_{\nu p}$ including the p-brane production for several representative values of $p$ and the number of extra dimensions $n$. The ratio of the length scales $L/L_*$ of the small to large extra dimensions has been set equal to 0.25 and the scale of quantum gravity $M_* = 1$ TeV. The highest energy HERA data point and the cross section obtained in Standard Model (SM) are also shown for comparison.

above for the primary neutrino energy of $10^{11}$ GeV. These are comparable to the nucleon-nucleon cross section at these energies and is typically the cross section values required in order that the neutrino can considered as a candidate for generating the observed air showers above the GZK bound. We point out that in the present case the neutrino delivers its entire energy in a single collision. This is in contrast to our earlier work [41], using a t-channel exchange in which the neutrino loses only a fraction of its energy in a single collision. In that case, the neutrino collides with several air nuclei along its path, and the resulting air shower tends to be longer in developing compared to an equally-energetic proton shower. Nevertheless, the characteristics of air showers do not rule out this model, given fluctuations and uncertainties in energy [41]. In the present case with high inelasticity, the position of the shower maximum will be determined dominantly by the value of the neutrino-air cross section. Since there exists a large range of parameter space where the neutrino-nucleon cross section exceeds 100 mb, the shower maximum could exist at the same or higher altitude than that produced by
Figure 2: The $L/L_*$ dependence of the neutrino-nucleon cross section $\sigma_{\nu p}$ for several representative value $p$ and the number of extra dimensions $n$. The scale of quantum gravity $M_*$ = 1 TeV and the incident neutrino energy is taken to be $10^{11}$ GeV.

a proton or even an iron nucleus of equivalent energy. Hence the position of the shower maximum may be indistinguishable from a hadron primary of equivalent energy.

We point out that our cross section does not violate any bounds imposed by unitarity. The models and parameter ranges considered in this paper are not ruled out by the low energy data. The cross section estimate is based on the calculation of the area of the classical brane solution, which essentially can be regarded as an estimate of the size of the region where the gravitational interaction between the two particles is strong. Therefore within these models, the total cross section has to be atleast as large as the estimate obtained by assuming that it is dominated by p-brane production. Furthermore as argued in Ref. [41] it is not possible to determine whether a cross section violates unitarity simply by looking at its rate of growth with energy. For example, the large $\sigma_{\nu-p}$ ($> 100$mb) obtained in Ref. [41] by assuming that it rises as $s^2$ with energy, is well within the bounds obtained by Goldberg and Weiler [42] by a model independent application of dispersion relations and unitarity. The growth of p-brane production cross section with energy is relatively mild in any case, with the total cross section rising at most.
as $E_{\nu}^{0.6}$ at large energy for the parameter ranges considered in this paper.

After formation the p-brane will decay. In analogy with black hole evaporation, we expect that branes will decay into gravitons and matter fields. The emitted gravitons will not be detected. Hence only a fraction of the particles that are produced will be responsible for generating the shower and the primary energy may be significantly underestimated. This implies that for a shower of energy $E$ we need a primary neutrino of energy $E'$ which is somewhat larger than $E$. The precise amount of energy which goes into gravitons is model dependent.

In conclusion we have shown that in low scale gravity models the neutrino-nucleon cross section is of the order of or larger than 100 mb for a large range of parameter space allowed by current experimental constraints. This makes the neutrino a candidate for explaining the cosmic ray events which appear to violate the GZK bound. Future cosmic ray observatories will be able to confirm or rule out this hypothesis. In particular, the ability of neutrino events to point back to their sources, which can be established by well-defined statistical correlations [43], has the potential to establish or to rule out neutrinos as primaries in the GZK-violation mystery. A very clear discussion on the possible origin of the highest energy cosmic rays is given in Ref. [44].

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References

[1] N. Arkani-Hamed, S. Dimopolous and G. Dvali, Phys. Lett. B\textbf{429}, 263 (1998); I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, Phys. Lett. B \textbf{436}, 257 (1998).

[2] L. Randall and R. Sundrum, Phys. Rev. Lett. \textbf{83}, 3370 (1999); Phys. Rev. Lett. \textbf{83}, 4690 (1999).

[3] T. Han, J.D. Lykken and R.-J. Zhang, Phys. Rev. D \textbf{59}, 105006 (1999); G.F. Giudice, R. Rattazzi and J.D. Wells, Nucl. Phys. \textbf{544}, 3 (1998); J. Hewett, Phys. Rev. Lett. \textbf{82}, 4765 (1999).
[4] S. Nussinov and R. Shrock, *Phys. Rev. D* **59**, 105002 (1999).

[5] P. Jain, D. McKay, S. Panda and J. Ralston, *Phys. Lett.* **B484**, 267 (2000), hep-ph/0001031.

[6] G. Domokos and S. Kovési-Domokos, *Phys. Rev. Lett.* **82**, 1366 (1999).

[7] M. Kachelrieß and M. Plümacher, *Phys. Rev. D* **62**, 103006 (2000).

[8] F. Cornet, J. I. Illana and M. Masip, *Phys. Rev. Lett.* **86**, 4235 (2001), hep-ph/0102065; S. Cullen, M. Perelstein and M. Peskin, *Phys. Rev. D* **62**, 055012 (2000).

[9] R. Emparan, M. Masip, and R. Rattazzi, *Phys. Rev. D* **65**, 064023 (2002), hep-ph/0109287.

[10] J. L. Feng and A. D. Shapere, hep-ph/0109106, *Phys. Rev. Lett.* **88**, 021303 (2002).

[11] L. Anchordoqui and H. Goldberg, *Phys. Rev. D* **65**, 047502 (2002), hep-ph/0109242.

[12] K. Greisen, *Phys. Rev. Lett.* **16**, 748 (1966); G. T. Zatsepin and V. A. Kuzmin, Sov. Phys. JETP Lett. **4**, 78 (1966).

[13] J. Puget, F.W. Stecker and J. Bredekamp, *Astrophys. J.* **205**, 638 (1976); F.W. Stecker and M.H. Salamon, *Astrophys. J.* **512**, 521 (1999); J. Wdowczyk, W. Tkaczyk and A. Wolfendale, *J. Phys. A* **5**, 1419 (1972).

[14] M. Nagano and A. Watson, *Rev. Mod. Phys.* **72**, 689 (2000).

[15] G. Sigl, Lect. Notes Phys. **556**, 259 (2000), astro-ph/0008364.

[16] Early versions of the idea of enhanced UHE neutrino-nucleon cross sections include V. Berezinsky and G. Zatsepin, *Phys. Lett. B* **28**, 423 (1969); G. Domokos and S. Nussinov, *Phys. Lett.* **B187**, 372 (1987); G. Domokos, B. Elliot, S. Kovési-Domokos, and S. Mrenna, *Phys. Rev. Lett.* **63**, 844 (1989); J. Bordes, H.-M. Chan, J. Faridani, J. Pfaudler and S.-T. Tsou, *Astropar. Phys.* **8**, 135 (1998).
[17] C. Tyler, A. Olinto and G. Sigl, *Phys. Rev. D* 63, 055001 (2001).

[18] L. Anchordoqui, H. Goldberg, T. McKauley, T. Paul, S. Reucroft and J. Swain, *Phys. Rev. D* 63, 124009 (2001), hep-ph/0011097; Anchordoqui, H. Goldberg, J. MacLeod, T. McCauley, T. Paul, S. Reucroft and J. Swain, hep-ph/0104114.

[19] S. Nussinov and R. Shrock, hep-ph/0103043, Phys. Rev. D 64, 047702 (2001).

[20] R. Emparan, hep-th/0104009, Phys. Rev. D 64, 024025 (2001).

[21] G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B 630, 293 (2002), hep-ph/0112161.

[22] T. Banks and W. Fischler, hep-th/9906038.

[23] S. B. Giddings and S. Thomas, Phys. Rev. D 65, 056010 (2002), hep-ph/0106219; S. B. Giddings, hep-ph/0110127.

[24] S. Dimopoulos and G. Landsberg, hep-ph/0106295, Phys. Rev. Lett. 87, 161602 (2001).

[25] D. M. Eardley and S. B. Giddings, gr-qc/0201034.

[26] S. Hossenfelder, S. Hofmann, M. Bleicher, H. Stöcker, hep-ph/0109085; M. Bleicher, S. Hofmann, S. Hossenfelder and H. Stöcker, hep-ph/0112186.

[27] K. Cheung, Phys. Rev. Lett. 88, 221602 (2002), hep-ph/0110163.

[28] R. Casadio and B. Harms, hep-th/0110255.

[29] A. Ringwald and H. Tu, hep-ph/0111042, Phys. Lett. B 525, 135 (2002).

[30] S. C. Park and H. S. Song, hep-ph/0111069.

[31] T. G. Rizzo, hep-ph/0111230.

[32] L. A. Anchordoqui, J. L. Feng, H. Goldberg, and A. D. Shapere, Phys. Rev. D 65, 124027 (2002), hep-ph/0112247.
[33] Y. Uehara, Prog. Theor. Phys. 107, 621 (2002), hep-ph/0110382.

[34] M. Kowalski, A. Ringwald and H. Tu, Phys. Lett. B529, 1 (2002), hep-ph/0201139.

[35] M. B. Voloshin, hep-ph/0107119, Phys. Lett. B 518 137 (2001); hep-ph/0111099, Phys. Lett. B 524, 376 (2002).

[36] Eun-Joo Ahn, M. Cavaglia, A. V. Olinto, hep-th/0201042.

[37] J. Lykken and L. Randall, JHEP 0006, 014 (2000), hep-th/9908076.

[38] I. I. Kogan and N. A. Voronov, JETP Lett. 38, 311 (1983).

[39] P. Candelas and S. Weinberg, Nucl. Phys. B 237, 397 (1984).

[40] H. L. Lai et al, Phys. Rev. D 51, 4763 (1995).

[41] A. Jain, P. Jain, D. McKay, and J. Ralston, hep-ph/0011310, IJMPA 17, 533 (2002).

[42] H. Goldberg and T. Weiler, Phys. Rev. D 59, 113005 (1999).

[43] Y. Uchihori et al, Astropart. Phys. 13, 151 (2000); G. Farrar and P. Biermann, Phys. Rev. Lett. 81, 3571, (1998); A. Virmani et al. astro-ph/0010235, Astropart. Phys. (to appear); G. Sigl, D. Torres, L. Anchordoqui, G. Romero, Phys. Rev. D 63, 081302 (2001); P. G. Tinyakov and I. I. Tkachev, Pisma Zh. Eksp. Teor. Fiz. 74, 499 (2001), astro-ph/0102476.

[44] J. Elbert and P. Sommers, Astrophys. J. 441, 151 (1995).