A New Strong Interaction Sector as the origin for the Dark Energy and Dark Matter

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(Dated: February 4, 2022)

A new strong interaction sector is proposed as a possible origin for the Dark Energy and Dark matter. It is described by an unbroken gauge group $SU(2)_Z$ which grows strong at a scale $\Lambda_Z \sim 10^{-3}eV$ provided its coupling is of the order of the electroweak coupling at some high energy scale (either $\sim 10^{16}$ GeV or $\sim O(\text{TeV})$). The present vacuum energy which is consistent with a cosmological constant is related to $\Lambda_Z$ and one of the $SU(2)_Z$ “matter fields” provides a candidate for the Dark Matter in the form of a weakly interacting massive particle (WIMP) with an annihilation cross section naturally of the order of the electroweak one at the time of its decoupling.

The nature of the dark matter and dark energy (responsible for an accelerating universe) is one of the deepest problems in contemporary cosmology. With the equation of state relating the pressure to the energy density being $p=\rho w$, the latest CTIO Lensing Survey performed in conjunction with CMB and Type Ia supernovae data up to redshift $z \sim 0.4$ found a value of $w$ which is consistent with $-1$. This, in turn, is consistent with the scenario in which the present total energy density is dominated by the vacuum energy or cosmological constant. However, it is well-known that this vacuum energy is surprisingly tiny, namely $\rho_V \approx (10^{-3} eV)^4$. We shall take the point of view that this could represent a new scale of physics on the same footing as other known scales ($\Lambda_{EW} \sim 250 GeV$, $\Lambda_{QCD} \sim 200 MeV$) rather than one in which the associated vacuum energies so finely cancel down to $(10^{-3} eV)^4$. (We assume that the electroweak and QCD vacuum energies are set to zero after the associated phase transitions are completed.) The purpose of this paper is to propose a new interaction which grows strong at $\Lambda_Z \sim 10^{-3} eV$ with interesting implications for the Dark Energy and Dark Matter.

The model we would like to propose is based on a very simple observation: For an unbroken gauge group (with a coupling comparable in magnitude to the $SU(2)_L$ gauge coupling at high energy) to become confining at energy scales many orders of magnitude below the QCD scale, its one-loop beta function has to be small compared with that of $SU(3)_c$. We propose:

- An unbroken vector-like gauge group $SU(2)_Z$.
- The initial value of the $SU(2)_Z$ gauge coupling is similar to that of the SM $SU(2)_L$ coupling at some high energy scale such as the early unification scale (PUT) in the TeV range (or even the GUT scale of $O(10^{16}$ GeV). This assumption makes the model fairly predictive as we shall see below (although a more general case can be easily entertained and will be dealt with elsewhere).
- Fermions (or chiral superfields in the SUSY case) transform as adjoints of $SU(2)_Z$ and singlets under the SM. We denote them by $\psi^{(Z)}_{L,R} = (1, 1, 0, 3)$ under $SU(3) \otimes SU(2)_L \otimes U(1)_Y \otimes SU(2)_Z$ for the non-SUSY case, and similarly $\Phi^{(Z)}_{L,R} = (1, 1, 0, 3)$ for the chiral superfield. These are all electrically neutral and the fermions would be qualified for the designation “sterile neutrinos”.
- Messenger fields transform either as fundamentals or as adjoints under $SU(2)_Z$ and possess SM quantum numbers. They are denoted by $\varphi^{(Z)}_{1/2} = (\varphi^{(Z)}_{1/2}, 0, \varphi^{(Z)}_{1/2}, 0, \varphi^{(Z)}_{1/2}, 0, \varphi^{(Z)}_{1/2}, 0, \varphi^{(Z)}_{1/2}) = (1, 2, Y_\varphi = -1, 2)$ or $\varphi^{(Z)}_1 = (\varphi^{(Z)}_1, 0, \varphi^{(Z)}_1, 0, \varphi^{(Z)}_1, 0, \varphi^{(Z)}_1, 0, \varphi^{(Z)}_1) = (1, 2, Y_\varphi = -1, 3)$ under $SU(3) \otimes SU(2)_L \otimes U(1)_Y \otimes SU(2)_Z$ for the non-SUSY case and $\chi_{L,R} = (1, 1, Y_\chi, 2)$ for the SUSY case (the triplet case is not allowed here if we want the model to be asymptotically free).

We split the discussion into two parts: Particle Physics and Cosmology.

(A) PARTICLE PHYSICS:

The basic one-loop RG equation is $\frac{d \lambda}{d \ln a} = -\beta_{\lambda}/2\lambda^2$ where $\beta_{0,N_S} = (11C_2(G)/3 - 2/3 \sum T_i(R) - 1/3 \sum T_S(R))$ for the non-SUSY case and $\beta_{0,S} = (3C_2(G) - \sum T_i(R))$ for the SUSY case, with $T_i(R), T_S(R), T_\chi(R)$ for chiral fermion, scalar, and chiral superfield respectively. From the particle content

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listed above, one has $C_2(G) = T_{\psi,z} = T_\Phi = 2$, $T_\psi = T_\chi = 1/2$ or $T_\psi = 2$ (for the Georgi-Glashow type of situation), giving $\beta_0_{NS} = (22 - 8n_\psi - n_\chi)/3$ or $\beta_0_{NS} = (22 - 8n_\psi - 4n_\phi)/3$, for the non-SUSY case, and $\beta_0_{NS} = (6 - 4n_\psi - n_\chi)$, for the SUSY case. Here, $n_\psi$, $n_\phi$, and $n_\chi$ already include both left and right-handed fields. We wish to make $\beta_0$ as small as possible while keeping it positive. This leads to the following minimal set: (I) $n_\psi = 2$; $n_\phi = 1$, for the non-SUSY case and (II) $n_\psi = 1$; $n_\chi = 1$, for the SUSY case. Above all mass thresholds, one has $\beta_0_{NS} = 5/3$ or $\beta_0_{NS} = 2/3$, $\beta_0_{S} = 1$, to be compared with $\beta_0_{SU(2)_L} = 10/3$ and $\beta_0_{SU(3)} = 7$ in the SM with three families.

Below, we will present a few numerical results, leaving a general analysis for a longer version of the paper.

(I) Non-SUSY $SU(2)_Z$:

Since $\varphi$ also carries the SM quantum number, it is safe and not unreasonable to assume that it has a mass of the order of the electroweak scale $\Lambda_{EW} \approx 250 GeV$. We will assume that the potential for $\varphi$ is such that $\phi < \varphi < 0$ in order for $SU(2)_Z$ to be unbroken.

For the fermions $\psi_i^{L,R}$ which are singlets with respect to the SM model, one can write down a gauge-invariant mass term of the type $\sum_{i=1}^2 m_i \psi_i^{Z_1} \bar{\psi}_i^{(2)}$. For the purpose of this paper, we will deal with the case where $m_1 \leq m_2$.

Let $M$ be the starting scale (early unification scale or GUT). The four basic equations which we will use are the following: 1) $\alpha_2(\Lambda_{EW}) = 02(M)/\{1 + (\alpha_2(M)/2\pi)(K_2 + K_3) \ln(\Lambda_{EW}/M)\}$ where $K_2 = 5/3$ for $\psi^{1/2}$ and $K_2 = 2/3$ for $\psi^{(2)}$; 2) $\alpha_2(m_2) = \alpha_2(\Lambda_{EW})/[1 + (\alpha_2(\Lambda_{EW})/\pi)\ln(m_2/\Lambda_{EW})]$; 3) $\alpha_2(m_1) = \alpha_2(m_2)/[1 + (7\alpha_2(m_2)/3\pi)\ln(m_1/m_2)]$; 4) $\alpha_2(m_1) = \alpha_2(m_2)/[1 + (7\alpha_2(m_2)/3\pi)\ln(m_1/m_2)]$; 5) $\Lambda_2 = m_2e^{[(-3\pi/11\alpha_2(m_1)) - (1 - \alpha_2(m_1))] - \alpha_2(m_2)}$.

For comparison, let us recall that the $SU(2)_L$ coupling at $M$ is $\alpha_2^{-1}(M) \approx 29.6$. At $2 TeV$, it is roughly (depending on unknown particle threshold) $\alpha_2^{-1}(2 TeV) \approx 28.2$. At $M = 2 \times 10^{10} GeV$, $\alpha_2^{-1}(M) \approx 45$ for the non-SUSY case and $\alpha_2^{-1}(M) \approx 24$ for the SUSY case with three generations.

In Table I, we list the results coming from the requirement: $\alpha_2(\Lambda) \approx 1$ at $\Lambda \approx 10^{-3} eV$, for doublet (bi-fundamental) $\varphi^{1/2}$ and triplet $\varphi^{(2)}$.

From Table I, one notices how important it is to have a "large" value for $m_2$, a non-zero value for $m_1$ and to have the presence of $\varphi^{(2)}$ if we wish $\alpha_2(M)$ to have a value close to the SM one.

(II) SUSY $SU(2)_Z$:

We will assume the mass of the $\chi$ superfield is of the order of the electroweak scale. We will further assume that SUSY breaking is such that the scalar field inside $\Phi^{(2)}_{L,R}$ has a mass $m_S$ while the fermionic partner remains massless. We obtain: $\{\alpha_2^{-1}(M) = 22.5; m_S = 1.5 GeV\}$, $\{\alpha_2^{-1}(M) = 1/23.5; m_S = 150 GeV\}$, for $M = 2 \times 10^{16} GeV$, $\Lambda_Z \approx 1.4 \times 10^{-3} eV$, to be compared with the typical SUSY GUT scale $\alpha_2^{-1}(M) \approx 24$.

(III) Comments:

Although the above results are but a sample of a more comprehensive analysis, a pattern clearly emerges: The $SU(2)_Z$ gauge coupling can start out at high energy with a magnitude close to the value of the $SU(2)_L$ coupling and subsequently becomes large at a scale which is 11 orders of magnitude smaller than the QCD scale. When the lighter of the two $\psi^{(2)}$ has a mass at around 1 eV and the heavier one with a mass of O(100 GeV), the initial $SU(2)_Z$ gauge coupling is remarkably close in value to that of $SU(2)_L$. If one simply starts out with the assumption $\alpha_2(\Lambda) \sim 2$ (4), the sought-after scale $\Lambda_Z \approx 10^{-3} eV$ arises naturally. Notice that $\alpha_2(M)$ is much closer to $\alpha_2(M)$ for the early unification scenario 6 than for GUT as can be seen from Table I. Furthermore, this scenario is closer in spirit and in construction to the early unification picture 4 where the gauge couplings of the various $SU(2)_S$'s are assumed to be equal at the PUT unification scale of the order of a few TeVs.

(IV) Messenger fields and their implications:

The manner in which the $SU(2)_Z$ matter communicates with normal matter will depend on whether $\varphi^{(2)}$ is a doublet or a triplet of $SU(2)_Z$ for the non-SUSY scenario.

For $\varphi^{(2)}$, one can write down a gauge-invariant Yukawa coupling between $\varphi^{(2)}$, the normal leptons, and $\psi^{(2)}$ (and not $\psi^{(2)}$, the reason of which will become clear below) as:

$$L_{\varphi^{(2)}} = \sum_i g_{Y_{i1}} \bar{l}_i \varphi^{(2)} \psi_{2,R}^{(2)} + h.c.,$$

where $l_i$ is a normal lepton doublet and the sum is over three generations. (This could be accomplished by assuming a discrete symmetry in such a way that a similar coupling to $\psi^{(2)}$ is forbidden.) A scattering process of the type $l_i + \psi \rightarrow \psi + l_j$ via a heavy $\varphi_1$ has an

| TABLE I: Correlations between $m_1$, $m_2$ and $\alpha_2^{-1}(M)$ |
|----------------------|---------|---------|---------|
|                      | $m_1$   | $m_2$   | $\alpha_2^{-1}(M)$ |
| GUT ($\varphi^{(2)}_1$): $M = 2 \times 10^{16} GeV$ | 0       | 30 MeV  | 30       |
|                      | 1.5 eV  | 30 MeV  | 33       |
|                      | 0       | 200 GeV | 34       |
|                      | 1 eV    | 200 GeV | 36.5     |
| GUT ($\varphi^{(2)}_2$): $M = 2 \times 10^{16} GeV$ | 1 eV    | 200 GeV | 31.5     |
| PUT ($\varphi^{(2)}_1$): $M = 2 TeV$               | 0       | 30 GeV  | 25       |
|                      | 1 eV    | 30 GeV  | 27.8     |
|                      | 0       | 120 GeV | 25.6     |
|                      | 1 eV    | 120 GeV | 28.4     |
| PUT ($\varphi^{(2)}_2$): $M = 2 TeV$               | 0       | 200 GeV | 25.6     |
|                      | 1 eV    | 200 GeV | 28.2     |
amplitude proportional to $g_Y^2 g_Y / m_{\phi(2)}^2$. What the constraints on the arbitrary Yukawa couplings $g_Y$ might be is of great interest for the Dark Matter search.

For $\phi^{(2)}$ (as well as for the SUSY $\chi_{L,R} = (1, 1, Y_{\chi, 1}, 2)$), no such Yukawa coupling can be written down and the interaction of $SU(2)_Z$ matter with normal matter proceeds through its anomalous magnetic moment (obtained at two loops) $\mu^{(2)}_Z \sim [2m_e m_{1,2}/m_{\phi(2)}^2] [\alpha^2_Z / 16\pi^2] \times \log$ terms. After $\mu_B$, where $\mu_B$ is the Bohr magneton. (Recall that, in this case, $SU(2)_Z$ matter interacts with the messenger field by exchanging $SU(2)_Z$ “gluons”. Since $m_1 \sim 1$ eV, one can see that $\mu^{(2)}_Z$ is completely negligible. However, for $\mu^{(2)}_Z$ can be as large as $10^{-8}\mu_B$ (in the regime where $\alpha_Z = 1$). The Lagrangian for the interaction between $\psi^{(2)}_{1,2}$ and the photon can be written as

$$\mathcal{L}_{\psi^{(2)}_{1,2}} = \mu^{(2)}_Z \psi^{(2)}_{1,2} (\sigma_{\mu\nu}/2) \psi^{(2)}_{1,2} F^{\mu\nu} + h.c.$$ Notice that $\psi^{(2)}_{2}$ cannot decay into lighter normal SM particles, neither through Eq. (2) nor through its anomalous magnetic moment. As a result it is absolutely stable.

The messenger field $\phi^{(2)}$ carrying both $SU(2)_Z$ and SM quantum numbers can be searched for at colliders such as the LHC. As shown below, $\psi^{(2)}_{2}$ has all the properties for being a Dark Matter candidate in the form of a WIMP. The previous discussion will be useful in the search for the Dark Matter, the phenomenology of which will be presented below.

(B) COSMOLOGICAL IMPLICATIONS:

(I) $SU(2)_Z$ “brief thermal history”:

At $T \gg m_i$, where $m_i$ is a generic particle mass, all normal matter and matter that carries $SU(2)_Z$ quantum numbers are in thermal equilibrium and are characterized by a common temperature. The fact that $SU(2)_Z$ matter is in thermal equilibrium with normal matter is because the messenger fields, e.g. $\phi^{(2)}$, carry both SM and $SU(2)_Z$ quantum numbers and, therefore, can interact with normal matter as well as with the $SU(2)_Z$ “gluons” and fermions $\psi^{(2)}$. (Let us recall that $\phi^{(2)}$ also serves the purpose of slowing the evolution of the $SU(2)_Z$ coupling so that it can grow strong at $\sim m eV$ if $\alpha_Z(M) \sim \alpha_2(M)$.)

When $T < m_{\phi(2)}$, the now-non-relativistic messenger field $\phi^{(2)}$ begins to annihilate each other with their number density decreasing like $n \sim (m_{\phi(2)} T)^{3/2} \exp(-m T)$. For the charged messenger $\phi^{(2)+}$, the dominant annihilation cross section is the electromagnetic one which goes like $\sigma \sim \alpha^2/T^2$ which grows with decreasing $T$. Their relic density would be negligible. For $\phi^{(2)0}$, it can decay into either $2 Z$ bosons or into lighter SM particles via $2 Z$ bosons. Again, its present relic density would be negligible.

For $\psi^{(2)}_{2}$ with mass $m_2$, the situation becomes interesting when $T < m_2$. Notice that, since $m_{\phi(2)} > m_2$, $\psi^{(2)}_{2}$ is a stable particle. The dominant annihilation cross section (into $2 \psi^{(2)}$ or $2 SU(2)_Z$ gluons) behaves as $\sigma_Z \sim \alpha_2^2(T)/m_2^2$. As we have seen above, the crucial feature of the $SU(2)_Z$ model is that its coupling basically almost “parallels” that of the SM $SU(2)_L$. From the above analysis, one can make the approximation $\alpha_Z(T \sim M_Z) \sim \alpha_2(T \sim M_Z)$. This implies that $\sigma_Z$ is naturally of the order of the weak cross section for $m_2 = O(100 \text{GeV})$. This will have an interesting implication concerning the nature of the dark matter as shown below.

After $\phi^{(2)}$ and $\psi^{(2)}_{2}$ decouple, the only relativistic particles left in the $SU(2)_Z$ sector are the $SU(2)_Z$ “gluons” and $\psi^{(2)}_1$. As we have mentioned earlier, these particles interact very weakly with normal matter and decouple soon after $\phi^{(2)}$ and $\psi^{(2)}_1$ went out of equilibrium. Their temperature $T_Z$ would go like $R^{-1}$. It is important to compare $T_Z$ with $T$ for the normal matter at the time of Big Bang Nucleosynthesis (BBN). The SM number of degrees of freedom (including $\phi^{(2)}$) for three families above the top quark mass is $g_* = 427/4 + 8 = 459/4$ and between $m_e$ and $m_\mu$ $g_* = 43/4$. Because of entropy conservation, the relationship between $T_Z$ and $T$ at $m_e < T < m_\mu$ is simply $T_Z = (43/459)^{1/3} T \sim 0.45 T$. During that period, $g_*$ $SU(2)_Z = 33/2$ and $\rho_{SU(2)_Z}/\rho_{SM} = (66/43)(43/459)^{1/3} \sim 0.07$ and one can see that the presence of $SU(2)_Z$ “matter” will have a negligible effect on BBN. After $e^\pm$ decoupling, one obtains $T_Z = [(43/459)/(4/11)]^{1/3} T \sim T/3$.

(II) Dark Energy:

We now discuss two possible scenarios: one having to do with a first-order phase transition from an unbroken chiral symmetry to a broken one, and another with an axion-like potential for $SU(2)_Z$. If the present universe is dominated by a vacuum energy of the order of $(A Z)^3 \sim (10^{-3} eV)^4$ and is entering an inflationary stage, how long will it last? The first-order phase transition proceeds by the nucleation of bubbles of the true vacuum with a rate $\Gamma = A \exp(-5 S_{E_{\lambda}})$, where $A$ is a prefactor and $S_{E_{\lambda}}$ is the Euclidean action. If this rate is small, the inflationary period can last for a long time. (Notice also that $SU(2)_Z$ provides a rationale for the scale in front of the flat potential of $\Phi$.)

(a) Chiral phase transition:

Extensive studies of chiral phase transition in QCD have been performed and they appeared to suggest that it is of first-order. One might expect a similar first-order phase transition for $SU(2)_Z$. There have also been studies of the chiral phase transition by bubble nucleation, e.g. one in which the linear $\sigma$ model is used, and there it is seen that, at $T = T_0$, the bubble nucleation rate is negligible. A similar study for the $SU(2)_Z$ case is under investigation. Since one scenario for the messenger field involves $\phi^{(2)}$ in the adjoint representation, one can also investigate the phase structure of the Georgi-Glashow (Polyakov) model in 2+1 dimensions.

(b) Axion-like potential:
There has been interesting proposals to use the QCD (or QCD-like) axion potential \( V(a) = V_0[1 - \cos \frac{aZ}{v}] - \eta \cos \left( \frac{aZ}{v} + \gamma \right) \) as a model for the early inflationary scenario \([12]\). We would like to propose a similar scenario but, this time, as an explanation for the origin of the dark energy. The \( SU(2)_Z \) axion comes from a \( U(1)_{-\rho Q'q'} \)-like symmetry of the \( SU(2)_Z \) fermions. For the purpose of this paper, we will only need to consider the case of the \( SU(2)_Z \) instanton potential with \( N = 1 \). The \( SU(2)_Z \) axion \( a_Z \) (where \( \phi_Z = \int_a (v + p) \exp(i \frac{aZ}{\Lambda^4}) \)) has a periodicity \( a_Z \rightarrow a_Z + 2\pi v = a_Z + 2\pi f_a \), where we assume that the axion decay constant \( f_a = O(v) \). (For this analysis, it is irrelevant whether or not \( f_a = v \) as long as they are of the same order.) We write

\[
V(a_Z) = \Lambda^4_Z [1 - \cos \frac{a_Z}{v}] - \epsilon \frac{a_Z^4}{2\pi^2 v^4}.
\]

In Eq. \(2\), the first term on the right hand side represents the potential with two degenerate vacua and the second term represents a soft-breaking term lifting that degeneracy. The difference in energy density between the two vacua is \( \epsilon \). We will assume that \( \epsilon \leq \Lambda^4_Z \) but not too different from it and at the same time have a small barrier between the two vacua. In other words we assume \( \epsilon = O(\Lambda^4_Z) \) and try to see its consequence on the cosmology of the model. The false vacuum of the “axion” potential is at \( a_Z = 2\pi v \) and the true vacuum is at \( a_Z = 0 \). The scenario goes as follows.

(i) As \( T_Z \) drops below \( T_c \sim \Lambda_Z \sim 10^{-3} eV \), the universe is trapped in the false vacuum with \( a_Z = 2\pi v \). The total energy density is dominated by \( \epsilon = O(\Lambda^4_Z) \).

(ii) The first order phase transition to the true vacuum at \( a_Z = 0 \) proceeds by bubble nucleation. The Euclidean action \( S_E \), in the thin wall limit, can be computed by looking at \( \tilde{S} = \int_{a_Z=0}^{a_Z=2\pi v} \sqrt{2\Lambda^4_Z [1 - \cos \frac{a_Z}{v}]} da_Z = 8\pi \Lambda^2_Z v \), giving

\[
S_E = \frac{27\pi^2\tilde{S}^4}{2\epsilon^3} \geq 5 \times 10^5 \left( \frac{v}{\Lambda_Z} \right)^4.
\]

The scale \( v \) of \( U(1)_{-\rho Q'q'} \) breaking is unknown but even when it is of the order of the electroweak scale, one would still get \( S_E \geq 5 \times 10^{64} \) which is huge. This will assure us that the transition will take a very long time from the present epoch to complete.

(III) Dark Matter:

There are several possibilities concerning Dark Matter candidates in our model. To be more specific, we will concentrate on the non-SUSY case.

A combination of various data (WMAP, etc.) gives a constraint for the non-baryonic dark matter density as follows: \( \Omega_M h^2 = 0.135^{+0.008}_{-0.009} \), for \( h \approx 0.72 \) giving \( \Omega_M \sim 0.26 \). (See \([13]\) for an excellent recent review.) Among the candidates for Dark Matter, there is an interesting one that goes by the name of WIMP (for weakly interacting massive particle). An approximate solution to the Boltzmann equation gives \([14]\) \( \Omega_\chi h^2 \sim 3 \times 10^{-27} s^{-1/2} < \sigma_A v \), where \( \chi \) denotes a generic WIMP, \( \sigma_A \) the annihilation cross section and \( v \) the relative velocity. It is generally noticed \([14]\) that \( \Omega_\chi h^2 \) is of order unity if \( \sigma_A \) is of a typical weak cross section, e.g. \( \sigma_A \sim \alpha^2/m^2 \) with \( \alpha \sim O(0.01) \) and \( m \sim O(100\, GeV) \). In brief, a stable WIMP with a weak cross section might fulfill the requirement of being a plausible candidate for Dark Matter.

The \( \psi_2 \) particle with mass \( m_2 \sim O(100\, GeV) \) is just such a candidate. It is stable (see the comment made in the particle physics discussion). Its annihilation cross section is typically of the order of the weak cross section as we have discussed above. Let us recall that the fact that it is so is because \( \alpha_Z(T \sim m_2) \sim \alpha_0(m_2) \), a notable feature of our model.

Acknowledgments

I would like to thank Paul Frampton and Marc Sher for useful discussions. This work is supported in parts by the US Department of Energy under grant No. DE- A505-89ER40518.

[1] S. Perlmutter et al., Astrophys. J. 517, 565 (1999); A. Riess et al., Astron. J. 116, 1009 (1998).
[2] For a nice overview and an extended list of references, see W. L. Freedman and M. S. Turner, astro-ph/0308418.
[3] M. Jarvis et al., astro-ph/0502243.
[4] A. J. Buras and P. Q. Hung, Phys. Rev. D 68, 035015 (2003); A. J. Buras, P. Q. Hung, Ngoc-Khanh Tran, A. Puschenerieder, E. Wyszomirski, Nucl. Phys. B 699, 253 (2004).
[5] S. N. Gninenko and N. V. Krasnikov, Phys. Lett. B450, 165 (1999).
[6] C. Athanassopoulos et al. [LSND collaboration], Phys. Rev. Lett. 75, 2650 (1995).
[7] E. W. Kolb and M. S. Turner, The Early Universe, Addison-Wesley Publishing Company (1990).
[8] P. Q. Hung, hep-ph/0010120.
[9] See e.g. O. Scavenius, A. Dumitru, E. S. Fraga, J. T. Lenaghan, and A. D. Jackson, Phys.Rev.D63, 116003 (2001).
[10] See e.g. Dmitri Antonov, hep-th/0207224.
[11] H. Quinn and R. Peccei, Phys. Rev. Lett. 38, 1440 (1977); M. Dine, W. Fischler, and M. Srednicki, Phys. Lett. B104, 199 (1981); P. Sikivie, Phys. Rev. Lett. 48, 1156 (1982).
[12] K. Freese, J. A. Friedman and A. V. Olinto, Phys. Rev. Lett. 65, 3233 (1990); K. Freese, J. Liu, and D. Spolyar, hep-ph/0502177.
[13] G. Bertone, D. Hooper and J. Silk, hep-ph/0404175.