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Published in:
ArXiv pre-print

Publication date:
2018

Document Version
Early version, also known as pre-print

Link to publication

Citation for published version (HARVARD):
Sarrazin, M & Stasser, C 2018, 'Bounds on ortho-positronium and J/psi - Y quarkonia invisible decays and constraints on hidden braneworlds in a SO(3,1)-broken 5D bulk', ArXiv pre-print.

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Bounds on ortho-positronium and J/ψ − Υ quarkonia invisible decays and constraints on hidden braneworlds in a SO(3,1)-broken 5D bulk

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While our visible Universe could be a 3-brane, some cosmological scenarios consider that other 3-branes could be hidden in the extra-dimensional bulk. Matter disappearance toward a hidden brane is mainly discussed for neutron – both theoretically and experimentally – but other particles are poorly studied. Recent experimental results offer new constraints on positronium or quarkonium invisible decays. In the present work, we show how a two-brane Universe allows for such invisible decays. We put this result in the context of the recent experimental data to constrain the brane energy scale $M_B$ (or effective brane thickness $M_B^{-1}$) and the interbrane distance $d$ for a relevant two-brane Universe in a SO(3,1)-broken 5D bulk. Quarkonia present poor bounds compared to results deduced from previous passing-through-walls-neutron experiments for which scenarios with $M_B < 2.5 \times 10^{17}$ GeV and $d > 0.5$ fm are excluded. By contrast, positronium experiments can compete with neutron experiments depending on the matter content of each brane. To constrain scenarios up to the Planck scale, positronium experiments should be able to reach $\text{Br}(\alpha\text{-Ps} \to \text{invisible}) \approx 10^{-8}$.

PACS numbers: 11.27.+d, 11.25.Wx, 11.10.Kk, 36.10.Dr

Many cosmological scenarios consider the existence of hidden braneworlds, specifically to explain the origin of dark matter and dark energy or as alternatives to cosmic inflation [1–10]. Among the cosmological models of interest, universes involving SO(3,1)-broken 5D bulks containing at least two braneworlds are under consideration [8–10]. In the last decade, it has been shown that such scenarios imply matter exchange between branes, which is a possible way to test these models [11–16]. In this context, neutron disappearance (reappearance) toward (from) a hidden brane has been widely discussed both theoretically [11–13] and experimentally [14–16]. Nevertheless, other kinds of particles have been somewhat neglected, although with good reason. In the present paper, we consider recent experimental results about constraints on positronium or quarkonium invisible decays [17, 18]. Showing how a hidden braneworld allows for such invisible decays, we discuss about constraints on the distance between our visible brane and a hidden one, and on the brane energy scale (or brane thickness) in the bulk. Highlighting the significant efficiency of the neutron to probe the existence of hidden braneworld, we put theoretical constraints on the expected branching ratio for positronium and quarkonium invisible decays.

When considering a two-brane Universe, whatever its full high-energy description (i.e. whatever the number or properties of bulk scalar fields responsible for particle trapping on branes, the number of extra dimensions or the bulk metric, etc.), the fermion dynamics on both branes at a sub-GeV energy scale is the same as the dynamics of fermions in a $M_4 \times Z_2$ space-time in the context of the non-commutative geometry (see Refs. [11, 12] for details). There are then two copies of the Standard Model, each sector being localized in each brane. Assuming that each braneworld is a $M_B^{-1}$-thick domain wall – where $M_B$ is the brane energy scale – the two sectors are mutually invisible to each other at the zeroth-order approximation for processes with energies below $M_B$. By contrast, matter fields in separate branes can mix at a first-order approximation mainly through the Lagrangian $L_c = ig\bar{\psi}_+\gamma^\mu\psi_+ + ig\bar{\psi}_-\gamma^\mu\psi_-$, where $\psi_\pm$ are the Dirac fermionic fields in each braneworld – denoted (+) and (−) [11, 12]. The interbrane coupling $g$ is given by [11]:

$$g = (m^2/M_B) \exp(-nd),$$

and depends on the distance $d$ between branes, on their effective thickness $M_B^{-1}$ and on the mass $m$ of the particle of the Standard Model under consideration, i.e. a constituent quark [11, 19–22] or a lepton. More specifically for bound quarks, $g$ presents a cut-off at $d \approx 0.5$ fm beyond which it cancels [11].

From $L_c$, one can show that a particle could oscillate between two states, one localized in our brane, the other localized in the hidden world. In fact, the oscillation would be driven by the effective magnetic field $B_\perp = g(A_\perp - A_\perp)$ transverse-like to the branes, where $A_\perp$ are the magnetic vector potentials in each brane. Specifically, the interaction Hamiltonian $H_c$ between the Pauli spinors of the visible and hidden worlds is given by [11–13]:

$$H_c = h\Omega \begin{pmatrix} 0 & \varepsilon \\ \varepsilon^\dagger & 0 \end{pmatrix},$$

where $\varepsilon = -i\sigma \cdot B_\perp/B_\perp$ is a unitary matrix acting on the spin with $n = (\sin \theta \cos \phi, \sin \theta \sin \phi, \cos \theta)$.

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and $\hbar \Omega = \mu_n B_\perp$ with $\mu_n$ the magnetic moment of the
particle. Here vector potentials $\mathbf{A}_\pm$ are dominated by the huge ($\gtrsim 10^9 \, \text{T m}$) overall astrophysical magnetic
vector potential $\mathbf{A}_{\text{amb}}$ related to the magnetic fields of all the astrophysical objects (planets, stars, galaxies, etc.) [13, 14, 24–28] and such that $\mathbf{A}_+ - \mathbf{A}_- \approx \mathbf{A}_{\text{amb}}$ [13, 14].

In the present work we consider bound states $\bar{p}p$ constituted by a fermion $p$ and its anti-particle $\bar{p}$. These
states are characterized by a finite lifetime and some decay channels. More precisely, we consider the long-living
$1^3S_1$ state (ortho-state) and the short-living $1^1S_0$ state (para-state) of the $\bar{p}p$ bound pair. Following the standard
description, the ortho-state wave function can be expressed as [29]:

$$
|\psi_1\rangle = |p, \uparrow\rangle \otimes |\bar{p}, \uparrow\rangle , \quad (3)
$$

$$
|\psi_{-1}\rangle = |p, \downarrow\rangle \otimes |\bar{p}, \downarrow\rangle , \quad (4)
$$

$$
|\psi_0\rangle = (1/\sqrt{2}) \left( |p, \uparrow\rangle \otimes |\bar{p}, \downarrow\rangle + |p, \downarrow\rangle \otimes |\bar{p}, \uparrow\rangle \right) , \quad (5)
$$

and similarly for the para-state wave function, we have:

$$
|\psi_P\rangle = (1/\sqrt{2}) \left( |p, \uparrow\rangle \otimes |\bar{p}, \downarrow\rangle - |p, \downarrow\rangle \otimes |\bar{p}, \uparrow\rangle \right) , \quad (6)
$$

where $|p, \uparrow\rangle$ and $|\bar{p}, \downarrow\rangle$ relate to the wave functions of the fermion and its anti-particle respectively both taking into
account the spin state.

Let us now consider the relevant interbrane coupling, which is simply the sum of two $\mathbf{H}$. Hamiltonian operators
concerning to the contributions of the particle and its anti-particle [29]:

$$
W = -i\hbar \Omega \left( \begin{array}{cc} 0 & (\sigma_p - \sigma_{\bar{p}}) \cdot \mathbf{n} \\ - (\sigma_p - \sigma_{\bar{p}}) \cdot \mathbf{n} & 0 \end{array} \right) , \quad (7)
$$

where the minus sign ($-\sigma_{\bar{p}}$) arises from the opposite magnetic moment of these particles. Considering the non-diagonal terms of the coupling term $W$, and the $\bar{p}p$ wave functions, we get:

$$
(\sigma_p - \sigma_{\bar{p}}) \cdot \mathbf{n} |\psi_1\rangle = -\sqrt{2} e^{i\phi} \sin \theta |\psi_P\rangle , \quad (8)
$$

$$
(\sigma_p - \sigma_{\bar{p}}) \cdot \mathbf{n} |\psi_0\rangle = 2 \cos \theta |\psi_P\rangle , \quad (9)
$$

$$
(\sigma_p - \sigma_{\bar{p}}) \cdot \mathbf{n} |\psi_{-1}\rangle = \sqrt{2} e^{-i\phi} \sin \theta |\psi_P\rangle , \quad (10)
$$

$$
(\sigma_p - \sigma_{\bar{p}}) \cdot \mathbf{n} |\psi_{-1}\rangle = 2 \cos \theta |\psi_0\rangle - \sqrt{2} e^{-i\phi} \sin \theta |\psi_1\rangle + \sqrt{2} e^{i\phi} \sin \theta |\psi_{-1}\rangle . \quad (11)
$$

It is noticeable that an ortho-state in a brane convert into a para-state only in the second brane and vice versa.

Then, considering only the ortho-state in our brane, the relevant wave function to consider for matter exchange can be written as:

$$
|\Psi\rangle = \left( \begin{array}{c} |\Psi_+\rangle \\ |\Psi_-\rangle \end{array} \right) = \left( \begin{array}{c} |\psi_1\rangle \\ |\psi_0\rangle \\ |\psi_{-1}\rangle \\ |\psi_P\rangle \end{array} \right) , \quad (12)
$$

where $|\Psi_{\pm}\rangle$ are the wave functions in each brane ($+$ or $-$), and $|\Psi_{O/P}\rangle$ the wave functions for the ortho-state
and para-state. The two-brane wave function $|\Psi(t)\rangle$ then follows:

$$
\text{i}\hbar \partial_t |\Psi(t)\rangle = (\mathbf{H}_0 + W) |\Psi(t)\rangle \quad (13)
$$

where the Hamiltonian $W$ can be recast as:

$$
W = \hbar \Omega \left( \begin{array}{cc} 0 & |C\rangle \\ \langle C| & 0 \end{array} \right) \quad (14)
$$

where $H_0 = \text{diag}(E_0 - i\hbar \Gamma_O, E_0 - i\hbar \Gamma_O, E_0 - i\hbar \Gamma_O, E_0 - i\hbar \Gamma_O)$. $E_0$ (respectively $E_P$) is the eigenenergy and $\Gamma_O$ (respectively $\Gamma_P$) is the decay rate of the ortho-state (respectively of the para-state).

$|\Psi(t)\rangle$ is obtained from the Lippmann–Schwinger equation:

$$
|\Psi(t)\rangle = |\Psi(0)\rangle + \int_{-\infty}^{+\infty} G^{(0)}(t-t') W |\Psi(t')\rangle \, dt' \quad (15)
$$

where the propagator follows $\text{i}\hbar \partial_t - H_0)G^{(0)}(t-t') = \delta(t-t')$ such that $G^{(0)}(t-t') = -i\hbar^{-1} e^{-\hbar^{-1}H_0(t-t')} \Theta(t-t')$, with $\Theta(t)$ the Heaviside step function. $|\Psi(0)\rangle = \text{i}hG^{(0)}(t)|\Psi^{(0)}\rangle$ is the solution of Eq. 13 when $W = 0$ with

$$
|\Psi^{(0)}\rangle = \left( \begin{array}{c} |\psi_{O}\rangle \\ 0 \end{array} \right) \quad (16)
$$

with $|p|^2 + |q|^2 + |r|^2 = 1$. $|\Psi^{(0)}\rangle$ defines the initial state when the particle is created. Assuming that $\hbar \Omega \ll E_0, E_P$ we use an up to the second-order expansion of the Lippmann–Schwinger equation such that:

$$
|\Psi(t)\rangle \sim |\Psi^{(0)}(t)\rangle \quad (17)
$$

$$
+ \int_{-\infty}^{+\infty} G^{(0)}(t-t')W |\Psi^{(0)}(t')\rangle \, dt' + \int_{-\infty}^{+\infty} \int_{-\infty}^{+\infty} G^{(0)}(t-t') W G^{(0)}(t'-t'') W |\Psi^{(0)}(t'')\rangle \, dt'' \, dt' \quad (18)
$$

We are now going to calculate the whole decay rate $\Gamma$ of the $\bar{p}p$ bound state in the two-brane Universe. We set $\Gamma = (1/2) \tau^{-1}$ where $\tau$ is the mean lifetime of the $\bar{p}p$ state given by $\tau = \int_{0}^{\infty} t f(t) \, dt$ where the related distribution function for the probability to observe the particle is $f(t) = -(d/dt) P$ with $P = (\langle \Psi(t) | \Psi(t) \rangle)$. Then, we get:

$$
\Gamma^{-1} = 2 \int_{0}^{\infty} \langle \Psi(t) | \Psi(t) \rangle \, dt . \quad (19)
$$

Using Eq. 17, we can now
Table I: Summary of positronium and charmonium masses, full widths and invisible decay branching ratios from theoretical and experimental data from Refs. [17–23]. $\varepsilon$ and $\varepsilon'$ describe the mass difference between the ortho-state and the para-state of the positronium, such that $|\varepsilon - \varepsilon'| = 0.84$ meV.

| System | Mass | Full width $2\Gamma$ | $\text{Br}(o-p\bar{p} \rightarrow \text{invisible})$ |
|--------|------|---------------------|---------------------------------|
| Positronium o-Ps | $1.022 - \varepsilon$ MeV | $7.211 \times 10^6$ s$^{-1}$ (th) | $< 5.9 \times 10^{-4}$ at 90 % C.L. |
| | p-Ps | $1.022 - \varepsilon'$ MeV | $8.033 \times 10^6$ s$^{-1}$ (th) | - |
| Charmonium $J/\psi$ | $3996.900 \pm 0.006$ MeV | $92.2 \pm 2.8$ keV (exp) | $< 7 \times 10^{-4}$ at 90 % C.L. |
| | $\eta$ | $2983.4 \pm 0.5$ MeV | $31.8 \pm 0.8$ MeV (exp) | - |
| | $\eta_c$ | $9399.0 \pm 2.3$ MeV | $10^{+3}_{-1}$ MeV (exp) | - |

Equation 19 applies by using the relevant parameters (see Table I). Here the difference $E_O - E_P$ includes the following contributions:

$$E_O - E_P = (m_O - m_P)\varepsilon^2 + m_O V_+ - m_P V_-$$

where $m_O$ and $m_P$ are the masses of the ortho-state and of the para-state respectively, and $V_{\pm}$ are the gravitational potentials in each brane. Doing this, we consider the following hypotheses:

1. **Null electric dipole moment hypothesis**

   In our previous works [11, 15, 16], it has been underlined that the interaction between the magnetic vector potential and the charge of a particle precludes the particle swapping between branes. As a consequence, since we consider electron/positon pairs or quark/antiquark pairs, we must consider these pairs as strictly chargeless and without charge structure. Let $E_{bind}$ be the typical binding energy of the fermion/antifermion pair. Then, the time $\Delta t \sim h/E_{bind}$ during which quantum fluctuations allow for an instantaneous electric dipole moment (EDM) must be smaller than the period $T \sim h/(E_O - E_P)$ of the Rabi oscillation during which the $p\bar{p}$ state oscillates between each brane. This allows a time averaged EDM equal to zero. That means that $E_O - E_P \ll E_{bind}$ must be verified. This condition is verified in quarkonium where the mass-energy difference between the ortho and para states dominates (i.e. $E_O - E_P \sim (m_O - m_P)\varepsilon^2$). Indeed, $E_{bind} \approx 1$ GeV and $\approx 0.7$ GeV while $E_O - E_P \approx 61$ MeV and $\approx 114$ MeV for the bottonium and the charmonium respectively [23]. By contrast, for the positronium the gravitational potential is now dominant (i.e. $E_O - E_P \sim m_{Ps} (V_+ - V_-)$) and $E_{bind} \approx 6.8$ eV [29]. Then, positronium is a relevant probe if $|V_+ - V_-| \lesssim 10^{-3}\varepsilon^2$. The consequence of this is discussed later in this paper.

2. **Global gravitational potential hypothesis**

   The gravitational potential $V_{\pm}$ felt by the particles incorporates many contributions. Some are local and for a neutron with a given polarization state such that $|u\rangle = \alpha |\uparrow\rangle + \beta |\downarrow\rangle$ (with $|\alpha|^2 + |\beta|^2 = 1$), considering the off-diagonal terms of $H$, proportional to $\sigma \cdot n$, we get $\sigma \cdot n |u\rangle = \alpha'|\uparrow\rangle + \beta'|\downarrow\rangle$. Although $(\alpha', \beta') \neq (\alpha, \beta)$ in most cases (i.e. the polarization is modified when the neutron leaps into the hidden brane), one gets: $|\alpha'|^2 + |\beta'|^2 = 1$ such that the swapping amplitude does not depend on the polarization state of the neutron. It is not longer the case of the $p\bar{p}$ states here under consideration due to the structure of their wave functions.
A can be then roughly constrained. Galactic magnetic
velocities used are \(m_b = 4730 \text{ MeV}\) and \(m_c = 1550 \text{ MeV}\) [19–
expressed in the context of large-scale superstructures [30, 31].
TABLE II: Derived constraints on \(\hbar \Omega\) for positronium and
\(\Omega_m = 0.3\) and \(\Omega_c = 0.7\) [19–

(iii) Global magnetic vector potential hypothesis

The overall ambient astrophysical magnetic vector potential \(A_{amb}\) was previously discussed in the literature in a different context [24–28]. \(A_{amb}\) is the sum of all of the magnetic vector potential contributions related to the magnetic fields of astrophysical objects (planets, stars, galaxies, etc.) since \(B(\mathbf{r}) = \nabla \times \mathbf{A}(\mathbf{r})\). At large distances from sources (for instance, close to the Earth), \(A_{amb}\) is almost uniform (i.e. \(\nabla \times A_{amb} \approx 0\)) and cannot be cancelled out with magnetic shields [28]. As a rule of thumb, \(A \approx DB\) where \(D\) is the distance from the astrophysical source and \(B\) is the typical field induced by the object. The expected order of magnitude of \(A_{amb}\) can be then roughly constrained. Galactic magnetic field variations on local scales towards the Milky Way core \(A_{amb} \approx 2 \times 10^9 \text{ T m}\) are usually assumed [24–28].

By contrast, the Earth’s magnetic field leads to 200 T m while the Sun contributes 10 T m [28]. By contrast, the magnitude of intergalactic contributions are less clear [13]). Then, for now \(A_{amb}\) can be fairly bounded by the lower estimated value: \(A_{amb} = 10^9 \text{ T m}\).

FIG. 1: (Color online). Bounds on the brane energy scale \(M_B\) and on the interbrane distance \(d\) derived from each particle kind: bottonium (yellow), charmonium (magenta), positronium (cyan) and neutron (purple). Each domain on the left of a colored band is excluded. Each colored band widens to the right when sensitivity increases: up to the Planck scale for neutron and positronium (Br(\(\alpha-Ps \rightarrow \text{invisible}\)) = \(1.4 \times 10^{-6}\), upper to Br(charmonium \(\rightarrow \text{invisible}\)) = \(10^{-6}\) for quarkonia.

As a result, from table I and Eq. 19 we deduce the results shown in table II. Table II gives the upper bounds for the coupling constants \(\langle \hbar \Omega \rangle\) or \(g\) both for bound states and some fundamental particles of the Standard Model. From table II data and Eq. 1, we can now derive lower bounds on the brane energy scale \(M_B\) and upper bounds on the interbrane distance \(d\) – for a \(SO(3,1)\)-broken \(M_4 \times R_1\) bulk – which are independent of the particle under consideration. The results are summarized on Fig. 1. Each domain on the left of a colored band is excluded: If matter disappearance can occur in the present braneworld scenario, these values of \(\langle M_B, d \rangle\) are not relevant. One notes that bounds are much poorer for quarkonia than for positronium or neutron. For the latter, it is noticeable that the constraints are just one order of magnitude below the reduced Planck energy scale, except if one accepts as possible a fine tuning of the distance \(d\) between 1 fm or 1 Å allowing for significantly low values for \(M_B\) (see Fig. 1). To be more quantitative, table III (left column) gives the bounds of the brane energy scale \(M_B\) for each particle.
Considering the same two-brane scenario and assuming now that the brane energy scale is the reduced Planck energy scale, we can derive the expected branching ratio for the positronium and quarkonia invisible decays (see right column in table III), which need to be reached to observe the phenomenon. While the expected branching ratio for quarkonia is far beyond any current experimental skill, it is not the case for positronium which could compete with neutron to constrain braneworld scenarios [14–16]. In addition, in Fig. 1, each colored band widens to the right when sensitivity increases: up to the Planck scale for neutron and positronium (\(\text{Br}(\text{o-Ps} \rightarrow \text{invisible}) = 1.4 \times 10^{-6}\)), up to \(\text{Br}(\text{quarkonium} \rightarrow \text{invisible}) = 10^{-6}\) for quarkonia.

As a major result, pretty low-energy experiments – involving positronium or neutron – appear then more suitable to probe the Planck scale than experiments needing colliders to produce quarkonia states. Even for a brane energy scale \(M_B\) at the Planck energy, low-energy disappearance phenomena can occur and could be observed in future low-energy experiments while the Planck scale is directly unreachable with colliders. However, to occur, positronium oscillations imply that the gravitational potentials in each brane must verify \(|V_+ - V_-| \lesssim 10^{-6}c^2\), i.e. a value one order of magnitude lower that the upper gravitational constraint mentioned above. As a result, positronium and neutron experiments are complementary as they could allow to discriminate between two kinds of gravitational environment.

### Acknowledgements

Coraline Stasser’s PhD is supported by the Special Research Fund (FSR) of the University of Namur. The authors thank Nicolas Reckinger for reading the manuscript.

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**Table III:** Left column: Various bounds on the brane energy scale \(M_B\) derived from each particle kind for interbrane distances below 0.5 fm. Right column: Expected branching ratios for a brane energy scale at the reduced Planck energy (2.43 \(\times 10^{18}\) GeV).

| Particle | \(M_B\)  | \(\text{Br}(\text{o-Ps} \rightarrow \text{invisible})\) |
|----------|----------|--------------------------------------------------|
| o-Ps     | \(> 1.2 \times 10^{17}\) GeV | \(1.4 \times 10^{-6}\) |
| \(J/\psi\) | \(> 1.0 \times 10^{12}\) GeV | \(1.3 \times 10^{-16}\) |
| \(Y\)     | \(> 2.4 \times 10^{12}\) GeV | \(2.9 \times 10^{-16}\) |
| neutron  | \(> 2.5 \times 10^{17}\) GeV | |

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[1] D. Battefeld, P. Peter, Phys. Rept. **571** (2015) 1-66, arXiv:1406.2790 [astro-ph.CO].
[2] T. Koivisto, D. Wills, I. Zavala, JCAP **06** (2014) 036, arXiv:1312.2397 [hep-th].
[3] R. Maartens, K. Koyama, Living Rev. Relativity **13** (2010) 5, arXiv:1004.3962 [hep-th].
[4] J.-L. Lehners, Phys. Rept. **465** (2008) 223-263, arXiv:0806.1245 [astro-ph].
[5] P. Brax, C. van de Bruck, A.-C. Davis, Rep. Prog. Phys. **67** (2004) 2183-2231, arXiv:hep-th/0404011.
[6] J. Khoury, B.A. Ovrut, P.J. Steinhardt, N. Turok, Phys. Rev. D **64** (2001) 123522, arXiv:hep-th/0103239.
[7] N. Arkani-Hamed, S. Dimopoulos, G. Dvali, N. Kaloper, JHEP **0112** (2000) 010, arXiv:hep-ph/9911386.
[8] D.J.H. Chung, K. Freese, Phys. Rev. D **67** (2003) 103505, arXiv:astro-ph/0202066.
[9] D.J.H. Chung, E.W. Kolb, A. Riotto, Phys. Rev. D **65** (2002) 083516, arXiv:hep-th/0008126.
[10] D.J.H. Chung, K. Freese, Phys. Rev. D **62** (2000) 063513, arXiv:hep-ph/9910235.
[11] D.J.H. Chung, K. Freese, Phys. Rev. D **62** (2000) 063513, arXiv:hep-ph/9910235.
[12] M. Sarrazin, F. Petit, Phys. Rev. D **81** (2010) 035014, arXiv:0903.2408 [hep-th].
[13] M. Sarrazin, F. Petit, Eur. Phys. J. C **72** (2012) 2230, arXiv:1208.2014 [hep-ph].
[14] M. Sarrazin, et al., Phys. Lett. B **758** (2016) 14, arXiv:1604.07861 [hep-ex].
[15] M. Sarrazin, et al., Phys. Rev. D **91** (2015) 075013, arXiv:1501.06468 [hep-ph].
[16] M. Sarrazin, G. Pignon, F. Petit, V.V. Nesvizhevsky, Phys. Lett. B **712** (2012) 213, arXiv:1201.3949 [hep-ph].
[17] V. Vigo, L. Gerchow, L. Liszkgay, A. Rubbia, P. Crivelli, Phys. Rev. D **97**, 092008 (2018), arXiv:1803.05744 [hep-ex].
[18] C. Patrignani et al (Particle Data Group), Chin. Phys. C **40**, 090001 (2016).
[19] D. Griffiths, *Introduction to Elementary Particles*, Wiley-VCH Verlag GmbH & Co. KGaA (2008).
[20] W. Lucha, F.F. Schöberl, D. Gromes, Phys. Rep. **200** (1991) 127-240.
[21] H.J. Lipkin, Phys. Lett. B **233** (1989) 446-450.
[22] I. Cohen, H.J. Lipkin, Phys. Lett. **93B** (1980) 56-60.
[23] H. Satz, J. Phys. G **32**, R25 (2006), arXiv:hep-ph/0512217.
[24] D.D. Ryutov, Phys. Rev. Lett. **103** (2009) 201803.
[25] E. Adelberger, G. Dvali, A. Gruzinov, Phys. Rev. Lett. **98** (2007) 010402, arXiv:hep-ph/0306245.
[26] J. Luo, L.-C. Tu, Z.-K. Hu, E.-J. Luan, Phys. Rev. Lett. **90**, 081801 (2003).
[27] J. Luo, C.-G. Shao, Z.-Z. Liu, Z.-K. Hu, Phys. Lett. A **270** (2000) 288.
[28] R. Lakes, Phys. Rev. Lett. **80** (1998) 1826.
[29] O. Halpern, Phys. Rev. **94** (1954) 904.
[30] P. Schneider, *Extragalactic Astronomy and Cosmology*, Springer (2015).
[31] P. Peter, J.-P. Uzan, *Primordial Cosmology*, Oxford Graduate Texts (2013).