ON THE ANOMALOUS CP VIOLATION AND NONCONTRACTIBILITY OF THE PHYSICAL SPACE

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There is a growing evidence for the anomalously large semileptonic CP asymmetry in the B meson system measured at the Tevatron. The noncontractible space, as an alternative symmetry-breaking mechanism to the Higgs mechanism, can change standard field theoretic calculations of the physical processes mediated through quantum loops for large external momenta or large internal masses. The presence of the W bosons and t quarks in loops of the B meson mixing can enhance the corresponding semileptonic CP asymmetry when the loop integration is up to the universal Lorentz and gauge invariant UV-cutoff. We show that the enhancement is roughly 13%, thus the possible deviation is measurable at the Tevatron, LHCb, SuperKEKB and SuperB facilities.

1. Introduction and motivation

Any theory beyond the Standard Model (SM) in particle physics must fulfil very severe theoretical and experimental requirements. Let us enumerate the most important ones: (1) presence of massive neutrinos and a very probable CP violation in the lepton sector, (2) it must contain a candidate for a cold Dark Matter particle, (3) the theory should explain the existence of only three fermion families, mass hierarchy and very small neutrino masses, (4) the SM CP violation in the baryon sector is insufficient to generate a very large excess of baryons over antibaryons in the Universe, (5) the theory of elementary particles should have some fundamental relationship to the theory of gravity.

One can also find the theoretical attempts, where the Dark Matter problem is resolved with the modified theories of gravity without any Dark Matter particle. However, the rotational curves of galaxies, gravity lensing by dark

halos, fluctuations of the CMBR, large scale galaxy and cluster correlations, Sunyaev–Zeldovich effect, age of the Universe, *etc.* are phenomena that are very difficult to simultaneously explain without the introduction of Dark Matter [1].

Recent Tevatron results [2] for the semileptonic CP asymmetry of B_d and B_s mesons reveal much larger baryon CP violation than the SM expectations. They induced a lot of theoretical attempts with beyond the SM physics in order to reproduce the experimental data.

In this paper, we want to investigate the impact of the new symmetry-breaking mechanism on the most sensitive electroweak observables. The theory proposed in [3] is formulated to understand and contrive the relations between gauge, conformal and discrete symmetries in particle physics and solving the problem of the unitarity (SU(2) global anomaly) and the problem of the UV (zero distance) infinity.

The acronym BY denotes precisely the attributes of the theory with respect to the two previously mentioned issues: A theory contains Nambu–Goldstone scalars without carrying the lepton number, while B theory contains Nambu–Goldstone scalars with the lepton number. X denotes a theory with the Higgs mechanism, Y denotes a theory of the noncontractible space, *i.e.*, the existence of the universal UV-cutoff. Thus, the four possible combinations are: AX, BX, AY and BY. The AX theory is the SM, while our favourite theory BY is trying to elucidate two essentially nonperturbative problems: the SU(2) global anomaly and the UV singularity.

The universal gauge and Lorentz invariant UV-cutoff in the space-like domain of the Minkowski spacetime $\Lambda = \frac{\hbar}{cd} = \frac{2}{g}\frac{\pi}{\sqrt{6}}M_W \simeq 326$ GeV, $g =$ weak coupling, $M_W =$ weak boson mass, is fixed by the trace anomaly [3]. It is the only parameter of this symmetry-breaking mechanism.

Let us briefly summarize the particle content and the most important phenomenological consequences of the BY theory:

— Light Majorana neutrinos as a hot Dark Matter component and heavy Majorana neutrinos as a cold Dark Matter component of the Universe are SU(3) singlet fermions of the SU(3) conformal unification scheme of the BY theory. The absence of the Higgs scalar appears crucial for the cosmological stability of the heavy Majorana neutrinos as cold Dark Matter particles [4] which decay to the pairs of the weak interaction gauge bosons [5].

— A study of Dyson–Schwinger equations of the Abelian model in the first nontrivial approximation within UV-finite theory gives mass functions of fermions with the magnitudes close to that observed in nature. Since any mass function has the form $m(q^2, \alpha_g, \Lambda) = \Lambda f(q^2, \alpha_g)$, where q^2 is Lorentz invariant squared momentum, $\alpha_g = g^2/4\pi$ is the coupling and

$\Lambda = 326$ GeV is the UV-cutoff defined by the weak boson masses and the weak coupling, if the universal UV-cutoff exists, it must be fixed by weak interactions [6]. The appearance of multinode solutions is a general phenomenon of bootstrap equations that helps to resolve the number and the mass gaps between the fermion families [6].

— The lepton CP violating phase and the dynamic of the heavy Majorana neutrinos, which are strongly coupled to the Nambu–Goldstone scalars, suffice to provide enough power to generate lepton asymmetries in the early Universe. Light Majorana neutrinos, on the other hand, induce vorticity of the Universe with right-handed chirality [7] within the Einstein–Cartan cosmology. The resulting angular momentum of the Universe can play the role of the dark energy within the nonsingular Einstein–Cartan cosmology [8].

Since the BY theory has very appealing features such as the existence of the mixed light neutrinos and the Dark Matter particle, we enumerate some additional phenomenological problems based on recent experiments that could be explained within the UV-finite BY theory.

The Tevatron experiments reported already [9] larger cross sections than predicted by the standard field theoretic QCD (see also Tevatron papers quoted in [10]). Although the shape of the cross sections at large scales ($\mu > 250$ GeV) could be reproduced by the SM QCD, the magnitude of the cross section quotient $\sigma(630\,\text{GeV})/\sigma(1800\,\text{GeV})$ is more than 10% away from the SM prediction. This is in accordance with the prediction of the BY theory that the QCD in the noncontractible space is not an asymptotically free gauge theory $\lim_{\mu\to\infty}\alpha_{\text{s}}^{\Lambda}(\mu) \neq 0$ [10]. To one loop order, the quotient between the BY (Λ) and the SM (∞) strong couplings can be evaluated [10]: $\alpha_{\text{s}}^{\Lambda}/\alpha_{\text{s}}^{\infty}(\mu = 1; 2; 3.5; 7\,\text{TeV}) = 1.23; 1.31; 1.38; 1.47$, respectively.

The observed forward–backward t-quark pair asymmetry at the Tevatron [11] deviates substantially from the theoretical prediction [12]. A larger QCD coupling of the BY theory and the deviation of the corresponding box diagram from the standard field theoretic QCD estimate, could improve the agreement with the experimental value because the asymmetry is proportional to α_{s}. The systematic errors are reduced because the charge asymmetry is defined as a quotient of the difference and the sum of the forward and backward integrated cross sections. The observed enhancement [11] is nonresonant with larger deviations from the SM QCD for larger invariant masses of the t-quark pairs. The huge enhancement of the asymmetry within the BY theory is confirmed in [13].

The branching ratio for the rare decay $B_s \to \mu\mu$ appears to be lower for more than 30% in the BY theory compared with the SM [14]. The LHCb could measure this mode very soon [15].

High energy hadron colliders, such as the LHC and the Tevatron, require very demanding analyses of data with incorporated SM physics. Rediscovering the SM physics, the LHC is faced with the phenomena that can be attributed to the leading order as problems with the power of the QCD coupling at certain scales such as: energy flow [16], B^+ production [17], bottomonium production [18], J/Ψ production [19], high energy dijets [20], multiplicities of charged particles [21], diphoton production [22], *etc.* These problems call for the reevaluation of the "background" physics with the BY theory.

In the next section, we are concentrated on the evaluation of the loop dominated mixing of the neutral mesons within the BY theory: SU(3) × SU(2)$_\mathrm{L}$ × U(1) gauge theory within noncontractible space ($\Lambda < \infty$) without the Higgs scalar.

2. CP violation and the B-meson mixing

The BY theory differs from the SM in the lepton sector already at the tree level because of the heavy Majorana neutrinos [3], but in the quark sector one has to search for electroweak processes mediated by quantum loops such as rare decays or CP violated transitions. It would be instructive to compare possible modifications to the CP violated mixings of the K and B mesons, which are due to the universal Lorentz and gauge invariant UV-cutoff $\Lambda \simeq 326$ GeV.

The formalism of meson mixing is very well known from the old studies of the strange quark physics, but here we refer to the updated analysis of B_s–$\bar{B}_s$ mixing of Lenz and Nierste [23] and references therein.

B_s–$\bar{B}_s$ oscillations can be studied by Schrödinger equation

$$\imath \frac{d}{dt} \left(\begin{array}{c} |B_s(t)\rangle \\ |\bar{B}_s(t)\rangle \end{array} \right) = \left(M^s - \frac{\imath}{2}\Gamma^s \right) \left(\begin{array}{c} |B_s(t)\rangle \\ |\bar{B}_s(t)\rangle \end{array} \right), \tag{1}$$

where the mass matrix is M^s and the decay matrix Γ^s. The physical eigenstates $|B_H\rangle$ and $|B_L\rangle$ are achieved by diagonalization of the matrix $M^s - \frac{\imath}{2}\Gamma^s$. We can write, to a very good approximation, for the mass and width differences [23]

$$\Delta M_s = M_H^s - M_L^s = 2|M_{12}^s|, \qquad \Delta\Gamma_s = \Gamma_L^s - \Gamma_H^s = 2|\Gamma_{12}^s|\cos\phi_s, \tag{2}$$

$$\text{CP phase } \phi_s = \arg(-M_{12}^s/\Gamma_{12}^s).$$

It is precisely the off-diagonal mass matrix element M_{12}^s that is potentially the most sensitive quantity to New Physics because it is defined by the quantum loop, *i.e.* a box diagram. The off-diagonal part of the decay matrix Γ_{12}^s is dominated by the tree level processes.

The SM prediction for M_{12}^s is well known [23]

$$M_{12}^s = \frac{G_{\mathrm{F}}^2 M_{B_s}}{12\pi^2} M_W^2 (V_{tb} V_{ts}^*)^2 \hat{\eta}_B S_0(x_t) f_{B_s}^2 B \,, \tag{3}$$

where $G_{\mathrm{F}} = $ Fermi constant, $V_{ij} = $ CKM matrix elements, $M_{B_s} = $ mass of B_s meson, $M_W = W$ boson mass, $\hat{\eta}_B = $ QCD correction factor, $S_0(x_t) = $ Inami–Lim function, $x_t = m_t^2/M_W^2$, $m_t = t - $ quark mass,

$$\langle B_s|Q|\bar{B}_s\rangle = \tfrac{8}{3} M_{B_s}^2 f_{B_s}^2 B \,,$$
$$Q = \bar{s}_\alpha \gamma_\mu (1 - \gamma_5) b_\alpha \bar{s}_\beta \gamma^\mu (1 - \gamma_5) b_\beta \,, \qquad \alpha, \beta = 1, 2, 3 = \text{colour indices} \,.$$

Lenz and Nierste [23] improved essentially the theoretical prediction for Γ_{12}^s by introducing more a natural operator basis and resumming charm quark contributions.

Now, we inspect where we can expect the largest deviation of the BY theory from the SM. Apart from the nonperturbative estimates of the matrix elements with hadrons, one encounters also QCD corrections, but at the scale of $\mu \simeq m_b$, the difference between the SM and BY strong couplings is negligible [10]. The largest deviation is expected at short distance contributions when calculating box diagrams with the heaviest internal particles.

The SM formulas of the box diagrams (two virtual W bosons and two virtual u, c or t quarks) responsible for K–$\bar{K}$, B_d–$\bar{B}_d$ or B_s–$\bar{B}_s$ mixings are provided by Inami and Lim [24]. A pedagogical derivation of the box diagrams in R_ξ gauges for vanishing external masses and momenta can be found in [25]. The final gauge invariant result looks like the following

$$\text{Box diagram (Fig. 17.2; Ref. [19])} = \frac{g^4}{4} (\bar{s}_1 \Gamma^\mu d_1)(\bar{s}_2 \Gamma_\mu d_2)$$

$$\times \sum_{\alpha=u,c,t} \sum_{\beta=u,c,t} \lambda_\alpha \lambda_\beta \int \frac{d^4 k}{(2\pi)^4} \frac{1}{D_\alpha D_\beta D_W^2} \left(\frac{k^2 m_\alpha^2 m_\beta^2}{4 M_W^4} + k^2 - \frac{2 m_\alpha^2 m_\beta^2}{M_W^2} \right) , \tag{4}$$

$$D_\alpha \equiv k^2 - m_\alpha^2 \,, \qquad D_W \equiv k^2 - M_W^2 \,, \qquad \Gamma_\mu \equiv \gamma_\mu \tfrac{1}{2}(1 - \gamma_5) \,,$$

$$\lambda_\alpha \equiv V_{\alpha s}^* V_{\alpha d} \quad \text{for} \quad K^0 - \bar{K}^0 \quad \text{system} \,;$$
$$\lambda_\alpha \equiv V_{\alpha b}^* V_{\alpha d} \quad \text{for} \quad B_d^0 - \bar{B}_d^0 \quad \text{system} \,;$$
$$\lambda_\alpha \equiv V_{\alpha b}^* V_{\alpha s} \quad \text{for} \quad B_s^0 - \bar{B}_s^0 \quad \text{system} \,.$$

Resolving the UV convergent integral in the Feynman fashion, one finds [25]

$$\text{Box diagram} = \frac{\imath g^4}{64\pi^2 M_W^2} (\bar{s}_1 \Gamma^\mu d_1)(\bar{s}_2 \Gamma_\mu d_2) \sum_{\alpha=u,c,t} \sum_{\beta=u,c,t} \lambda_\alpha \lambda_\beta F(x_\alpha, x_\beta) \,, \tag{5}$$

$$F(x_\alpha, x_\beta) \equiv \frac{1}{(1-x_\alpha)(1-x_\beta)}\left(\frac{7x_\alpha x_\beta}{4} - 1\right) + \frac{x_\alpha^2 \ln x_\alpha}{(x_\beta - x_\alpha)(1-x_\alpha)^2}$$

$$\times \left(1 - 2x_\beta + \frac{x_\alpha x_\beta}{4}\right) + \frac{x_\beta^2 \ln x_\beta}{(x_\alpha - x_\beta)(1-x_\beta)^2}\left(1 - 2x_\alpha + \frac{x_\alpha x_\beta}{4}\right),$$

$$x_\alpha \equiv \frac{m_\alpha^2}{M_W^2}.$$

This form can be simplified by the unitarity of the CKM matrix $\lambda_u + \lambda_c + \lambda_t = 0$ and with the approximation $m_u = 0$, we get [25]

$$\begin{aligned}
\text{Box diagram} &= \frac{-\imath G_F^2 M_W^2}{2\pi^2}(\bar{s}_1 \Gamma^\mu d_1)(\bar{s}_2 \Gamma_\mu d_2)\mathcal{F}_0, \qquad (6)\\
\mathcal{F}_0 &= \lambda_c^2 S_0(x_c) + \lambda_t^2 S_0(x_t) + 2\lambda_c\lambda_t S_0(x_c, x_t),\\
S_0(x_c, x_t) &= -F(x_c, x_t) - F(0,0) + F(0, x_c) + F(0, x_t),\\
S_0(x_c) &= \lim_{x_t \to x_c} S_0(x_c, x_t),\\
S_0(x) &= \frac{x}{(1-x)^2}\left[1 - \frac{11x}{4} + \frac{x^2}{4} - \frac{3x^2 \ln x}{2(1-x)}\right].
\end{aligned}$$

The dominant t-quark short distance contribution in the effective quark theory for the mass difference M_{12}^s enters in the form of the Inami–Lim function $S_0(x_t)$.

The Nambu–Goldstone scalars in the BY theory are decoupled from quarks at the tree level [3] and they carry lepton number. Thus, one has to perform calculations of the box diagrams in the unitary gauge. The form of the subintegral function in Eq. (4) is the same as in the SM because of the gauge invariance ($\xi \to 0$ in the unitary gauge: $\Delta_{\mu\nu} = -\imath(g_{\mu\nu} - \frac{p_\mu p_\nu}{M^2})\frac{1}{p^2 - M^2} - \imath\frac{p_\mu p_\nu}{M^2}\frac{1}{p^2 - M^2/\xi}$, $\Delta = \frac{\imath}{p^2 - M^2/\xi}$). If it is easier to calculate an observable in R_ξ rather than in the unitary gauge, it is allowed to do so even in the BY theory because of the gauge invariance (*i.e.* independence on the ξ parameter) of the observable. The perturbation theory with coupling as a perturbation parameter in the strong coupling system of the heavy Majorana neutrinos and the Nambu–Goldstone scalars is useless and Dyson–Schwinger equations within the nonsingular BY theory are a framework for the analysis.

Let us divide the UV convergent integral of the box diagram Eq. (4) into two UV convergent integrals introducing the UV-cutoff after Wick's rotation. These integrals can be evaluated numerically or analytically by elementary functions (albeit with lengthy expressions) to recheck their sum evaluated previously in Eq. (5)

$$F(x_\alpha, x_\beta) \equiv F^\infty(x_\alpha, x_\beta) = F^\Lambda(x_\alpha, x_\beta) - \Delta F^\Lambda(x_\alpha, x_\beta), \tag{7}$$

$$F^\Lambda(x_\alpha, x_\beta) = -M_W^2 \int_0^{\Lambda^2} dz\, z \left(z + m_\alpha^2\right)^{-1} \left(z + m_\beta^2\right)^{-1} \left(z + M_W^2\right)^{-2}$$

$$\times \left(z \left(1 + \frac{m_\alpha^2 m_\beta^2}{4M_W^4}\right) + \frac{2m_\alpha^2 m_\beta^2}{M_W^2} \right),$$

$$\Delta F^\Lambda(x_\alpha, x_\beta) = +M_W^2 \int_0^{1/\Lambda^2} dw \left(1 + m_\alpha^2 w\right)^{-1} \left(1 + m_\beta^2 w\right)^{-1} \left(1 + m_W^2 w\right)^{-2}$$

$$\times \left(1 + \frac{m_\alpha^2 m_\beta^2}{4M_W^4} + 2\frac{m_\alpha^2 m_\beta^2}{M_W^2} w \right).$$

Thus, the BY theory with the universal UV-cutoff defined by the weak boson mass $\Lambda = 326$ GeV predicts for M_{12}^s ($\hat{\eta}_B$ QCD correction function in the BY theory remains almost unchanged with respect to the SM: $\mu \simeq m_b \ll \Lambda$) owing to the electroweak box diagram (see Eq. (3) and Eqs. (5)–(7))

$$M_{12}^s(\text{BY}) = \frac{S_0^\Lambda(x_t)}{S_0^\infty(x_t)} M_{12}^s(\text{SM}), \tag{8}$$

$$S_0^\Lambda(x_c, x_t) \equiv -F^\Lambda(x_c, x_t) - F^\Lambda(0,0) + F^\Lambda(0, x_c) + F^\Lambda(0, x_t),$$

$$S_0^\Lambda(x_t) = \lim_{x_c \to x_t} S_0^\Lambda(x_c, x_t).$$

The final relation for the CP asymmetry in the flavour-specific $B_s \to f$ decays has the form [23]

$$a_{\text{fs}}^s = \frac{|\Gamma_{12}^s|}{|M_{12}^s|} \sin \phi_s. \tag{9}$$

In the last chapter, we present results and concluding remarks.

3. Results and conclusions

Lenz and Nierste [23] reported in 2007 2σ deviation for the CP violating phase ϕ_s despite of the fact that the mass difference M_{12}^s (Eq. (3)) contains poorly known V_{CKM} matrix elements and $f_{B_s}^2 B$ form factor. Hadron models still generate the largest uncertainty in studying semileptonic and hadronic

processes, where matrix elements are extracted: $\mid V_{td} \mid = (8.4 \pm 0.6) \times 10^{-3}$, $\mid V_{ts} \mid = (38.7 \pm 2.1) \times 10^{-3}$, $\mid V_{tb} \mid = 0.88 \pm 0.07$ [26]. For the detailed error budget analysis the reader can consult Ref. [23].

We can easily estimate the deviation of the BY theory prediction from the SM one. From the preceding chapter, one concludes that only the short distance electroweak part of M_{12}^s, hidden in the modified Inami–Lim function, can enhance the magnitude of the semileptonic CP asymmetry parameter [23]

$$
\begin{aligned}
a_{\mathrm{sl}}(\mathrm{SM}) &\simeq 0.582 a_{\mathrm{sl}}^d + 0.418 a_{\mathrm{sl}}^s \,, \\
M_W &= 80.4 \text{ GeV} \,, \qquad m_c = 1.3 \text{ GeV} \,, \qquad m_t = 172 \text{ GeV} \\
&\Rightarrow \kappa \equiv a_{\mathrm{sl}}(\mathrm{BY})/a_{\mathrm{sl}}(\mathrm{SM}) = S_0^\infty(x_t)/S_0^\Lambda(x_t) = 1.13 \,. \quad (10)
\end{aligned}
$$

In Fig. 1, the reader can visualize the enhancement factor as a function of the UV-cutoff, however one should bear in mind that Λ is not a free parameter in the BY theory.
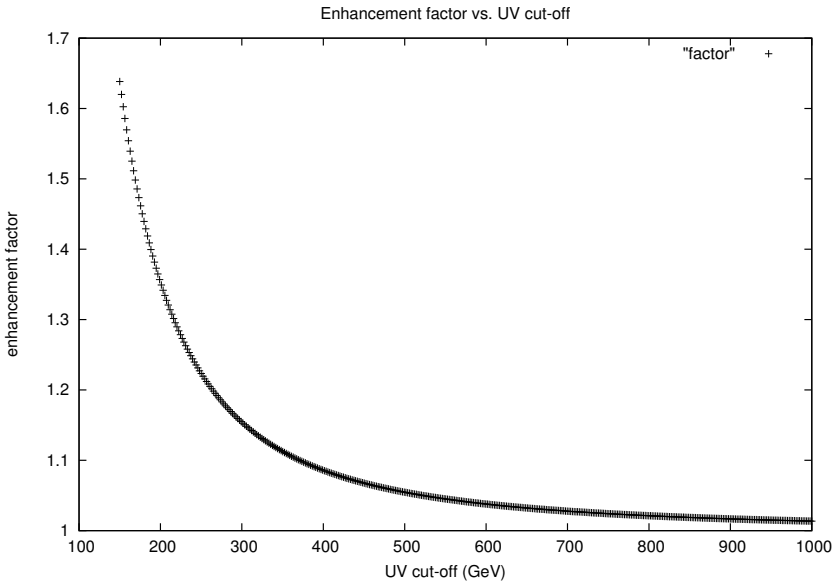

Fig. 1. Enhancement factor κ as a function of the UV-cutoff (GeV).

The larger magnitude of the asymmetry parameter for the BY theory by 13% is still significantly smaller than the median value of the experiment [2]. In numerical evaluations of the CP, asymmetry parameter of the B_d system more intermediate experimental values are incorporated than for B_s system. However, the BY theory predicts the same enhancement for both systems.

It is important to notice that the mixing of the neutral kaons is also affected by the virtual heavy t quark. We can estimate the deviation from the SM (see chapter 17 of Ref. [25] for the theory and Ref. [27] for Wolfenstein parameters)

$$
\begin{aligned}
M_{12} \propto{} & \Re(\lambda_c^* \lambda_u)[\eta_1 S_0(x_c) - \eta_3 S_0(x_c, x_t)] \\
& + \Re(\lambda_t^* \lambda_u)[\eta_3 S_0(x_c, x_t) - \eta_2 S_0(x_t)],
\end{aligned}
\tag{11}
$$

QCD correction factors: $\eta_1 = 1.38$, $\eta_2 = 0.57$, $\eta_3 = 0.47$,

$$
V_{\text{CKM}} = \begin{pmatrix} 1 - \lambda^2/2 & \lambda & A\lambda^3(\rho - \imath\eta) \\ -\lambda & 1 - \lambda^2/2 & A\lambda^2 \\ A\lambda^3(1 - \rho - \imath\eta) & -A\lambda^2 & 1 \end{pmatrix} + \mathcal{O}\left(\lambda^4\right),
$$

Wolfenstein parameters for V_{CKM} : $\lambda = 0.2257$, $A = 0.814$, $\rho = 0.135$, $\eta = 0.349$,

$$
\begin{aligned}
S_0^\infty(x_t)/S_0^\Lambda(x_t) &= 1.13, \\
S_0^\infty(x_c, x_t)/S_0^\Lambda(x_c, x_t) &= 1.014, \\
S_0^\infty(x_c)/S_0^\Lambda(x_c) &= 1.000, \\
\Rightarrow |M_{12}(\text{SM})/M_{12}(\text{BY})| &= 1.055.
\end{aligned}
\tag{12}
$$

The most recent analysis of Lenz [28] based on the new LHCb measurements [29] rules out large SM deviation suggested by Tevatron data. Moreover, it strengthens the validity of the heavy quark expansion corrections (HQE) to the amplitudes [28]

$$
\left(\frac{\Delta\Gamma_s}{\Delta M_s}\right)^{\text{Exp}} \bigg/ \left(\frac{\Delta\Gamma_s}{\Delta M_s}\right)^{\text{SM}} = 1.12 \pm 0.27
\tag{13}
$$

and

$$
\frac{\Delta\Gamma_s^{\text{Exp}}}{\Delta\Gamma_s^{\text{SM}}} = 1.15 \pm 0.32.
\tag{14}
$$

The theory relies on the HQE, the QCD corrections and lattice gauge theory evaluations of the hadron matrix elements. The total discrepancy between the theory and the experiment at the level of 30% [28] does not exclude small short distance corrections of the BY theory. Since the BY theory is UV finite, the extensive use of the Dyson–Schwinger and Bethe–Salpeter amplitudes and equations can reassure lattice and HQE calculations. The dependence of QCD amplitudes in the strong coupling regime on the UV-cutoff is only logarithmic or log–log when solving the bootstrap equations.

The essential dependence of the solutions appears in the parametrization of the infrared sector [30] that does not differ from the SM one.

To conclude, the noncontractible spacetime as a symmetry breaking mechanism in the BY theory of Ref. [3] changes the SM processes essentially via quantum loops if internal masses or external momenta and masses are large [10]. We show that the enhanced baryon CP violation in the BY theory cannot explain Tevatron results [2], but its prediction can help to resolve some present and possible future discrepancy of the unitarity relations of the mixing matrix. It seems that new measurements at the LHCb [29] of $B_s \to J/\Psi f_0(980)$ and Φ do not support a large deviation from the SM prediction. It could be more favourable to search for the electroweak deviations of the SM in the $B_s \to \mu\mu$ decay [14], where the hadron physics uncertainty is smaller.

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