

Unparticle effects in rare $t \rightarrow cgg$ decay

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Abstract

Rare $t \rightarrow cgg$ decay can only appear at loop level in the Standard Model (SM), and naturally they are strongly suppressed. These flavor changing decays induced by the mediation of spin-0 and spin-2 unparticles, can appear at tree level in unparticle physics. In this work the virtual effects of unparticle physics in the flavor-changing $t \rightarrow cgg$ decay is studied. Using the SM result for the branching ratio of the $t \rightarrow cgg$ decay, the parameter space of $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}$, where the branching ratio of this decay exceeds the one predicted by the SM, is obtained. Measurement of the branching ratio larger than 10^{-9} can give valuable information for establishing unparticle physics.

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1 Introduction

Rare decays, induced by the flavor-changing neutral current (FCNC) transitions, is quite a promising research area and attracts theoretical interest as a potential testing ground for checking predictions of the Standard model (SM) at loop level, as well as search for new physics effects (NP) beyond the SM. The effects of NP in rare decays can appear in two different ways: i) via new contributions to the Wilson coefficients existing in the SM, ii) via appearance of new operators with new Wilson coefficients which are absent in the present SM.

The impressive and exciting results on the FCNC decays in B-meson sector were observed at two B-meson factories KEK and BELLE [1–3], and at CLEO [4], which are in good agreement with the SM prediction.

The interest to the study of FCNC decays in t-quark sector can be explained on the following reasons: i) In many models beyond the SM the new physics scale is closer to the t-quark mass, and ii) many two-body t-quark FCNC decays, like $t \rightarrow cV$ ($V = g, \gamma, Z$) and $t \rightarrow cH$ are highly suppressed in the SM due to the GIM mechanism and their branching ratios are of the order $10^{-11} \div 10^{-14}$ [5, 6]. These branching ratios are practically impossible to measure at LHC [7] or at International Linear Collider (ILC) [8]. But many models of NP predict that the branching ratios of the above-mentioned FCNC decays are much larger compared to that obtained in the SM ([9] and references therein).

The t-quark three-body FCNC decays like $t \rightarrow cWW, cZZ, bWZ$ are also discussed in the framework of SM [10–12] and its beyond [13]. It is shown in [10–12] that the rate of higher order three-body FCNC decay $t \rightarrow cgg$ exceeds the rate of lower order $t \rightarrow cgg$ decay.

As has already been noted, FCNC processes are very sensitive to the new physics effects. One such model is the so-called unparticles is recently proposed by H. Georgi [14]. The main idea in this model is that at very high energies the SM fields and the Banks-Zaks (BZ) fields with a nontrivial infrared fixed point interact. The interaction between these two sectors is due to the exchange of particles with a large mass scale M_U . The interaction below this scale is nonrenormalizable and is suppressed by a power of M_U . The renormalizable couplings of BZ fields then produce the dimensional transmutation and the scale invariant unparticle emerges below the scale Λ_U , and the unparticle stuff with scale dimension d_U , looks like massless invisible particles with noninteger number d_U . For this reason, production of unparticles might be detectable in missing energy processes. Phenomenology of unparticle physics is studied extensively in the literature [15–27]. In the present work we study $t \rightarrow cgg$ ($t \rightarrow c\gamma\gamma$) decay in the framework of an unparticle physics. Organization of this paper is as follows: In section-2 calculation of the $t \rightarrow cgg$ and $t \rightarrow c\gamma\gamma$ decays are presented. In section-3 numerical results and discussion are given.

2 Formalism

In this section we calculate the branching ratio of $t \rightarrow cgg$ ($t \rightarrow c\gamma\gamma$) decay in unparticle physics. As we have noted already, below $\Lambda_U = 1 \text{ TeV}$ the interaction between SM fields

and BZ fields become an effective operator, i.e., it has the following form:

$$\mathcal{L}_{int} = \frac{1}{\Lambda_{\mathcal{U}}} O_{SM} O_{\mathcal{U}} . \quad (1)$$

Obviously, High-dimension operators should be suppressed by inverse power of $\Lambda_{\mathcal{U}}$. Therefore, we should choose the appropriate operators with the lowest dimension. Also, the effective interaction should satisfy the SM gauge symmetry. The effective Lagrangian of scalar and tensor unparticle operators with SM fields are given in [28]:

a) scalar unparticle

$$\begin{aligned} & \lambda_0 \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \bar{f} f O_{\mathcal{U}} , \\ & \lambda_0 \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \bar{f} i \gamma_5 f O_{\mathcal{U}} , \\ & \lambda_0 \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \bar{f} \gamma_{\mu} (\gamma_5) f \partial_{\mu} O_{\mathcal{U}} , \\ & \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \left[\lambda_0 G_{\alpha\beta} G^{\alpha\beta} + \lambda'_0 G_{\alpha\beta} \tilde{G}^{\alpha\beta} \right] O_{\mathcal{U}} , \end{aligned} \quad (2)$$

b) tensor unparticle

$$\begin{aligned} & -\frac{1}{4} \lambda_2 \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \bar{\psi} i \left[\gamma_{\mu} (\gamma_5) \overleftrightarrow{\mathcal{D}}_{\nu} + \gamma_{\nu} \gamma_5 \overleftrightarrow{\mathcal{D}}_{\mu} \right] \psi O_{\mathcal{U}}^{\mu\nu} , \\ & \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \left[\lambda_2 G_{\mu\nu} G^{\mu\nu} + \lambda'_2 G_{\mu\alpha} \tilde{G}^{\alpha}_{\nu} \right] O_{\mathcal{U}}^{\mu\nu} , \end{aligned} \quad (3)$$

where

$$\mathcal{D}_{\mu} = \partial_{\mu} + ig \frac{\tau^a}{2} W_{\mu}^a + ig' \frac{V}{2} B_{\mu} ,$$

is the covariant derivative in the SM, $G^{\alpha\beta}$ is the gauge field strength tensor (gluon, photon, as well as weak gauge bosons),

$$\tilde{G}_{\alpha\beta} = \frac{1}{2} \varepsilon_{\mu\nu\alpha\beta} G^{\mu\nu} ,$$

f is a standard model fermion, and ψ stands for a SM doublet or singlet fermion, and λ_i are the dimensionless effective couplings. Note that we will neglect the third term in scalar unparticle case because this term contain an extra $1/\Lambda_{\mathcal{U}}$ factor.

It follows from Eqs. (2) and (3) that the flavor violating $t \rightarrow cgg$, $t \rightarrow c\gamma\gamma$ decays can take place at tree level in unparticle physics, while they exist at loop level in the SM, and this is the main reason why we consider them in unparticle physics.

Scale invariance determines the form of the propagators within normalization factor. The propagators corresponding to scalar and tensor unparticles are

$$\mathcal{D}(q^2) = \frac{A_{d_{\mathcal{U}}}}{2 \sin(d_{\mathcal{U}}\pi)} (-q^2)^{d_{\mathcal{U}}-2} , \quad \text{and}, \quad (4)$$

$$\Delta_{\mu\nu\rho\sigma} = \mathcal{D}(q^2) \mathcal{P}_{\mu\nu\rho\sigma} . \quad (5)$$

For the transverse and traceless tensor operators $O_{\mu\nu}$ [?] we have

$$\mathcal{P}_{\mu\nu\rho\sigma} = \frac{1}{2} \left\{ \Pi_{\mu\nu}\Pi_{\rho\sigma} + \Pi_{\mu\sigma}\Pi_{\nu\rho} - \frac{2}{3}\Pi_{\mu\rho} + \Pi_{\nu\sigma} \right\}, \quad (6)$$

while in conformal field theories [29]

$$\begin{aligned} T_{\mu\nu\rho\sigma} = & \frac{1}{2} \left[(g_{\mu\rho}g_{\nu\sigma} + g_{\mu\sigma}g_{\nu\rho}) + \frac{4 - d_{\mathcal{U}}(d_{\mathcal{U}} + 1)}{2d_{\mathcal{U}}(d_{\mathcal{U}} - 1)} g_{\mu\nu}g_{\rho\sigma} \right. \\ & - 2 \left(\frac{d_{\mathcal{U}} - 2}{d_{\mathcal{U}}} \right) \left(g_{\mu\rho} \frac{k_{\nu}k_{\sigma}}{k^2} + g_{\mu\sigma} \frac{k_{\nu}k_{\rho}}{k^2} + g_{\nu\rho} \frac{k_{\mu}k_{\sigma}}{k^2} + g_{\nu\sigma} \frac{k_{\mu}k_{\rho}}{k^2} \right) \\ & \left. + 4 \frac{d_{\mathcal{U}} - 2}{d_{\mathcal{U}}(d_{\mathcal{U}} - 1)} \left(g_{\mu\nu} \frac{k_{\rho}k_{\sigma}}{k^2} + g_{\rho\sigma} \frac{k_{\mu}k_{\nu}}{k^2} \right) + 8 \frac{(d_{\mathcal{U}} - 2)(d_{\mathcal{U}} - 3)k_{\mu}k_{\nu}k_{\rho}k_{\sigma}}{d_{\mathcal{U}}(d_{\mathcal{U}} - 1)(k^2)^2} \right], \quad (7) \end{aligned}$$

where

$$\Pi_{\mu\nu} = -g_{\mu\nu} + a \frac{q_{\mu}q_{\nu}}{q^2}, \quad (8)$$

where

$$a = \begin{cases} 1, & \text{for transverse vector operator, and,} \\ \frac{2(d_{\mathcal{U}} - 2)}{(d_{\mathcal{U}} - 1)}, & \text{in conformal field theory.} \end{cases}$$

In the present work we follow the Georgi's approach [14], namely, Feynman propagators of the unparticle operator $O_{\mathcal{U}}$ is determined by the scalar invariance.

The factor $A_{\mathcal{U}}$ in Eqs. (4) and (5) is

$$A_{\mathcal{U}} = \frac{16\pi^{5/2}}{(2\pi)^{2d_{\mathcal{U}}}} \frac{\Gamma(d_{\mathcal{U}} + 1/2)}{\Gamma(d_{\mathcal{U}} - 1)\Gamma(2d_{\mathcal{U}})}.$$

In order to calculate the decay rate $t \rightarrow cgg$ and $t \rightarrow c\gamma\gamma$ decays we also need $gg\mathcal{U}$, $\gamma\gamma\mathcal{U}$ and fermion–fermion unparticle interaction vertices. From Eqs. (2) and (3) we get the following expressions for the above–mentioned vertices:

a) fermion–fermion scalar unparticle

$$\frac{\lambda_0}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \bar{c} [C_S + i\gamma_5 C_P] t,$$

b) gluon–gluon scalar unparticle

$$\frac{\lambda_0}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} 2 \left\{ \frac{\lambda^a}{2} \frac{\lambda^b}{2} \right\} [(k_1 \cdot k_2)g_{\mu\nu} - k_{1\mu} \cdot k_{2\nu} - 2\epsilon_{\mu\nu\alpha\beta} k_{1\alpha} k_{2\beta}] \varepsilon_{\mu}^a(k_1) \varepsilon_{\nu}^b(k_2),$$

Photon–photon scalar unparticle vertex can be obtained from gluon–gluon scalar unparticle by making the replacement

$$\left\{ \frac{\lambda^a}{2} \frac{\lambda^b}{2} \right\} \rightarrow 1,$$

and omitting color indices in ε_μ^a , and hence we get:

c) fermion–fermion tensor unparticle

$$\frac{1}{4\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \left\{ \lambda_2 \left[\gamma_\mu (p_c + p_t)_\nu + \gamma_\nu (p_c + p_t)_\mu \right] + \lambda'_2 \left[\gamma_\mu \gamma_5 (p_c + p_t)_\nu + \gamma_\nu \gamma_5 (p_c + p_t)_\mu \right] \right\} ,$$

d) gluon–gluon tensor unparticle

$$\begin{aligned} \frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \left[\lambda_2 \left(\left\{ \frac{\lambda^a \lambda^b}{2} \right\} K_{\mu\nu\rho\sigma}^S + \left[\frac{\lambda^a \lambda^b}{2} \right] K_{\mu\nu\rho\sigma}^A \right) \right. \\ \left. + \lambda'_2 \left(\left\{ \frac{\lambda^a \lambda^b}{2} \right\} F_{\mu\nu\rho\sigma}^S + \left[\frac{\lambda^a \lambda^b}{2} \right] F_{\mu\nu\rho\sigma}^A \right) \right] \varepsilon_\mu^a(k_1) \varepsilon_\nu^b(k_2) , \end{aligned}$$

where

$$\begin{aligned} K_{\mu\nu\rho\sigma}^{S(A)} = \frac{1}{2} \left\{ (k_1 \cdot k_2) g_{\mu\rho} g_{\nu\sigma} + g_{\mu\nu} k_{1\rho} k_{2\sigma} - g_{\nu\sigma} + g_{\mu\nu} k_{1\rho} k_{2\mu} - g_{\mu\rho} + g_{\mu\nu} k_{1\nu} k_{2\sigma} \right. \\ \left. \pm \left[(k_1 \cdot k_2) g_{\mu\nu} g_{\rho\sigma} + g_{\mu\rho} k_{1\nu} k_{2\sigma} - g_{\rho\sigma} + g_{\mu\rho} k_{1\nu} k_{2\mu} - g_{\mu\nu} + g_{\mu\rho} k_{1\rho} k_{2\sigma} \right] \right\} , \end{aligned}$$

$$K_{\mu\nu\rho\sigma}^{S(A)} = \frac{1}{2} \left(k_{1\rho} k_{2\beta} \varepsilon_{\mu\nu\beta\sigma} - k_{1\alpha} k_{2\beta} g_{\mu\rho} \varepsilon_{\sigma\alpha\beta\nu} \mp k_{1\beta} k_{2\rho} \varepsilon_{\mu\nu\beta\sigma} \mp k_{1\beta} k_{2\alpha} g_{\rho\nu} \varepsilon_{\sigma\alpha\beta\mu} \right) .$$

Photon–photon tensor unparticle vertex can easily be obtained from gluon–gluon tensor unparticle vertex by making the following replacements:

$$\left\{ \frac{\lambda^a \lambda^b}{2} \right\} \rightarrow 1 , \quad \left[\frac{\lambda^a \lambda^b}{2} \right] \rightarrow 0 ,$$

and omitting color indices in ε_μ^a .

Now we are ready to calculate the branching ratio of the $t \rightarrow cgg$ and $t \rightarrow c\gamma\gamma$ decays. In calculation of the branching ratios of these decays there appear infrared and collinear divergences. There are three possible sources of these singularities:

- 1) One gluon (photon) flying parallel to the c–quark,
- 2) two gluons (photons) flying parallel to each other, and,
- 3) one of the gluons (photons) is soft.

First and second cases are related with the collinear singularity, while the last case is related to the infrared singularity.

In order to avoid the singularity in case–1, it is enough to take into account mass of the c–quark in calculations. There are two different ways to prevent the singularities in cases–2 and –3, one of them is to put cut–off factor in “dangerous” integration limit where singularities are present (see [30]).

Using the Feynman rules for the matrix element of the $t \rightarrow cgg$ decay exchanging the scalar and tensor unparticles, we get respectively,

$$M_S = T_{\mu\nu}^+ \left\{ \frac{\lambda^a \lambda^b}{2} \right\} \bar{c} \left[C_S + C_P \gamma_5 \right] t \varepsilon_\mu^a(k_1) \varepsilon_\nu^b(k_2) ,$$

$$\begin{aligned} M_T = \left(T_{\mu\nu\rho\sigma}^+ \left\{ \frac{\lambda^a \lambda^b}{2} \right\} + T_{\mu\nu\rho\sigma}^- \left[\frac{\lambda^a \lambda^b}{2} \right] \right) \mathcal{P}_{\rho_1, \sigma_1, \rho\sigma} \bar{c} \left\{ \lambda_2 \left[\gamma_{\rho_1} (p_c + p_t)_{\sigma_1} + \gamma_{\sigma_1} (p_c + p_t)_{\rho_1} \right] \right. \\ \left. + \lambda'_2 \left[\gamma_{\rho_1} \gamma_5 (p_c + p_t)_{\sigma_1} + \gamma_{\sigma_1} \gamma_5 (p_c + p_t)_{\rho_1} \right] \right\} t \varepsilon_\mu^a \varepsilon_\nu^b , \end{aligned}$$

where

$$T_{\mu\nu}^+ = \frac{\lambda_0}{\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}-1}} \frac{A_{\mathcal{U}}}{\sin(d_{\mathcal{U}}\pi)} \frac{1}{(q^2)^{2-d_{\mathcal{U}}}} \left\{ \lambda_0 [k_{1\nu} k_{2\mu} - g_{\mu\nu} (k_1 \cdot k_2)] + \lambda'_0 \epsilon_{\mu\nu\alpha\beta} k_{1\alpha} k_{2\beta} \right\},$$

$$T_{\mu\nu\rho\sigma}^{\pm} = \frac{1}{4\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}}} \frac{A_{\mathcal{U}}}{2\sin(d_{\mathcal{U}}\pi)} \frac{1}{(q^2)^{2-d_{\mathcal{U}}}} \left(\lambda_2 K_{\mu\nu\rho\sigma}^{S(A)} + \lambda'_2 F_{\mu\nu\rho\sigma}^{S(A)} \right). \quad (9)$$

In further analysis we take into account the following fact. It is well known that [31, 32] if in the problem under consideration there appear two or more external gluons whose polarization sum is $\sum_{\lambda} \varepsilon_{\mu}^*(k, \lambda) \varepsilon_{\nu}(k, \lambda) = -g_{\mu\nu}$, gauge invariance is violated. In our calculation we choose the following expression for the polarization sum of the gluons, simultaneously which are transverse to massless vector boson momenta k_1 and k_2 ,

$$P_{\mu\nu} = \sum_{\lambda=1,2} \varepsilon_{\mu}^*(k, \lambda) \varepsilon_{\nu}(k, \lambda),$$

$$= -g_{\mu\nu} + \frac{k_{1\mu} k_{2\nu} + k_{1\nu} k_{2\mu}}{k_1 \cdot k_2},$$

which leads to the gauge invariant result for on-shell massless vector mesons.

Using the matrix element for the $t \rightarrow cxx$ ($x = g, \gamma$) decay, in the rest frame system of the decaying t-quark, we get for the differential decay width

$$d\Gamma = \frac{1}{256m_t\pi^3} C_X |M'_X|^2 dE_C dE_1,$$

where prime means summation over gluon (photon) is performed, and C_X is the color factor whose values are presented in the table.

	$t \rightarrow cgg$		$t \rightarrow c\gamma\gamma$
	Antisymmetric	Symmetric	
C_X	$\frac{N^2-1}{2}$	$\frac{N^2-1}{2N^2}$ (Singlet) $\frac{(N^2-1)(N^2-2)}{2N^2}$ (Adjoint)	1

Table 1:

In order to calculate the branching ratio, we take into account that the $t \rightarrow bW$ decay is the dominant channel of the t-quark, and use $\Gamma(t \rightarrow bW) = 1.55 \text{ GeV}$.

3 Numerical analysis

In this section we study the sensitivity of the branching ratio on the scaling dimension parameter $d_{\mathcal{U}}$, energy scale $\Lambda_{\mathcal{U}}$ and the coupling constants. In numerical analysis we choose

the scaling dimension d_U in the range $1 < d_U < 2$. The main reason for choosing $d_U > 1$ is that in this region the decay rate is free from the nonintegrable singularity [14]. As has already been mentioned, there appear singularities for $d_U > 2$. Therefore we will consider the above-mentioned restricted domain of d_U . The values of the off-diagonal t - c unparticle coupling constants C_S and C_P are chosen in the range $10^{-1} \div 10^{-3}$. For the parameter alpha we choose three different values $\alpha = 0.1; 0.5; 1.0$. Note that the branching ratio of the $t \rightarrow cgg$ decay in the SM is calculated in [9] which predicts $\mathcal{B}(t \rightarrow cgg) \simeq 1.02 \times 10^{-9}$, when the cut-off parameter C is taken $C = 10^{-3}$. Our numerical calculations shows that when the cut-off parameter C varies in the range $c = 0.001 \div 0.1$ for a given set of the fixed values of C_P and C_S , no substantial change in the value of the branching ratio is observed, the variation being about three times. The above-mentioned value of the branching ratio of the $t \rightarrow cgg$ decay in the SM is too small to be observable in the forthcoming LHC experiments. For this reason any experimental observation of the $t \rightarrow cgg$ decay will definitely indicate the appearance of the new physics beyond the SM. Therefore, the observability limit of the $t \rightarrow cgg$ decay can be assumed to be $\mathcal{B}(t \rightarrow cgg) = 10^{-9}$. In this connection there follows the question about the range of values of d_U for which the branching ratio is larger than 10^{-9} , at the value $\Lambda_U = 1 \text{ TeV}$ of the cut-off parameter and at fixed values of the effective couplings C_P and C_S (in the presence of the scalar unparticle operator).

Our numerical analysis predicts the following results:

- at $C_S = C_P = 10^{-1}$, $d_U < 1.5$ (< 1.53 , < 1.55), and when $C = 0.1(10^{-2}, 10^{-3})$;
- at $C_S = C_P = 10^{-2}$, $d_U < 1.2$ (< 1.24 , < 1.25), and when $C = 0.1(10^{-2}, 10^{-3})$;
- at $C_S = C_P = 10^{-3}$,

the corresponding branching ratios are larger compared to the the SM result.

It follows from the above-presented results that the restrictions to the values of d_U in both decays, for which the branching ratio exceeds 10^{-9} , are practically the same.

As an illustration of our analysis, we present in Fig. (1) the dependence of the branching ratio of the $t \rightarrow cgg$ decay on d_U , at $C_S = C_P = 10^{-2}$, $C = 10^{-2}$, when scalar unparticle operator is the mediator. Here the parameter α is defined as $\alpha = \lambda'_0/\lambda_0$, and we set $\lambda_0 = 1$. From this figure we see that up to $d_U = 1.1$ the perpendicular spin polarization exceeds the parallel spin polarization for two-gluon system at $\alpha = 1$.

These results are quite interesting since they give valuable information about the scaling parameter d_U , as well as information about gluon-gluon unparticle coupling constants.

For the tensor operator case we obtain the restrictions $d_U < 1.4$ (< 1.55 , < 1.58) at $C = 0.1(10^{-2}, 10^{-3})$ for the $t \rightarrow cgg$ decay, for which the branching ratio exceeds 10^{-9} , at $\Lambda_U = 1 \text{ TeV}$.

Depicted in Fig. (2) is the dependence of the branching ratio for the $t \rightarrow cgg$ decay on d_U at $\Lambda_U = 1 \text{ TeV}$, when the mediator is the tensor particle. In this figure β is defined as $\beta = \lambda'_2/\lambda_2$. Similar to spin-0 unparticle mediator case, we set $\lambda_2 = 1$ in numerical calculations. It follows from this figure that, when the coupling constants of two gluon system with perpendicular and parallel spin orientations are equal, the branching ratio of the spin-perpendicular configuration exceeds the spin-parallel configuration of the two-gluon system up to $d_U = 1.15$.

Note that all above-presented results are obtained at $\Lambda_{\mathcal{U}} = 1 \text{ TeV}$. In this connection the question, how restrictions on $d_{\mathcal{U}}$ depend on the cut-off parameter $\Lambda_{\mathcal{U}}$, should be considered. In other words, at which parametric region of $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}$ the branching ratio is larger than 10^{-9} . In order to answer this question, we present in Figs. (3) and (4) the parametric plot of the branching ratio with respect to $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}$ which gives $\mathcal{B} = 10^{-9}$, for the $t \rightarrow cgg$ decay, at fixed the values of $C_S = C_P = 10^{-1}, 10^{-2}$ and $C = 10^{-2}$, in the presence of the scalar operator. The region on the right side of each curve should be excluded, since $\mathcal{B} < 10^{-9}$ in this domain. We observe that stringent constraints due to $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}$ are obtained for the $C_P = C_S = 10^{-2}$ case.

Figs. (5)–(7) depict the the same analysis for the tensor operator. It follows from these figures that the branching ratio is reachable to be investigated up to $\Lambda_{\mathcal{U}} = 10 \text{ TeV}$ and up to $d_{\mathcal{U}} = 1.5$.

In conclusion, we analyze the rare $t \rightarrow cgg$ decay, that can exist at tree level in unparticle physics. Note that these decays can take place only at loop level in the SM. For this reason the branching ratio of these decays in unparticle physics can exceed the ones predicted by the SM. The experimental measurement of the branching ratios larger than 10^{-9} can give valuable information about the existence of the new physics beyond the SM, in particular, about the unparticle physics.

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Figure captions

Fig. 1 The dependence of the branching ratio of the $t \rightarrow cgg$ decay on $d_{\mathcal{U}}$, at the values $C_P = C_S = 10^{-2}$ of the t-c unparticle coupling constants, at $C = 10^{-2}$ of the cut-off parameter, and at $\Lambda_{\mathcal{U}} = 1 \text{ TeV}$, when the scalar unparticle is the mediator.

Fig. 2 The same as in Fig. (1), but when the tensor unparticle is the mediator.

Fig. 3 The parametric plot of the dependence of $\Lambda_{\mathcal{U}}$ on the scaling parameter $d_{\mathcal{U}}$ at $C = 10^{-2}$ and $C_P = C_S = 10^{-1}$, when branching ratio for the $t \rightarrow cgg$ decay $\mathcal{B}(t \rightarrow cgg) = 1.2 \times 10^{-9}$, and when the scalar unparticle is the mediator.

Fig. 4 The same as in Fig. (3), but at $C_P = C_S = 10^{-2}$.

Fig. 5 The same as in Fig. (3), but at $C = 10^{-1}$, in the presence of tensor unparticles.

Fig. 6 The same as in Fig. (5), but at $C = 10^{-2}$.

Fig. 7 The same as in Fig. (5), but at $C = 10^{-3}$.

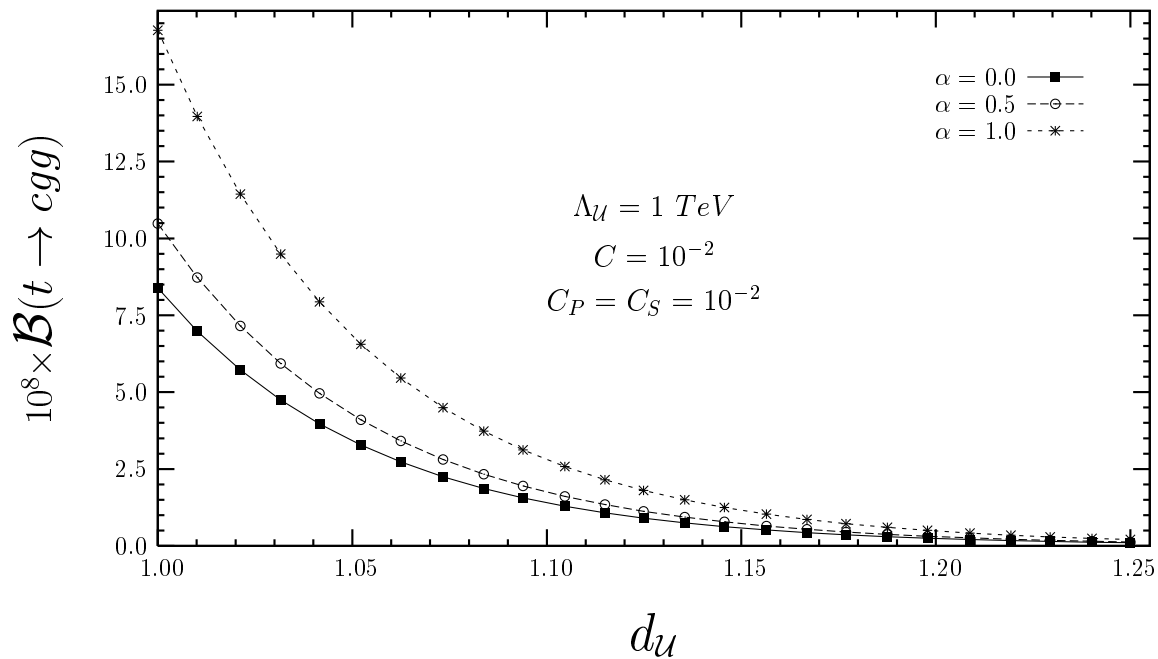


Figure 1:

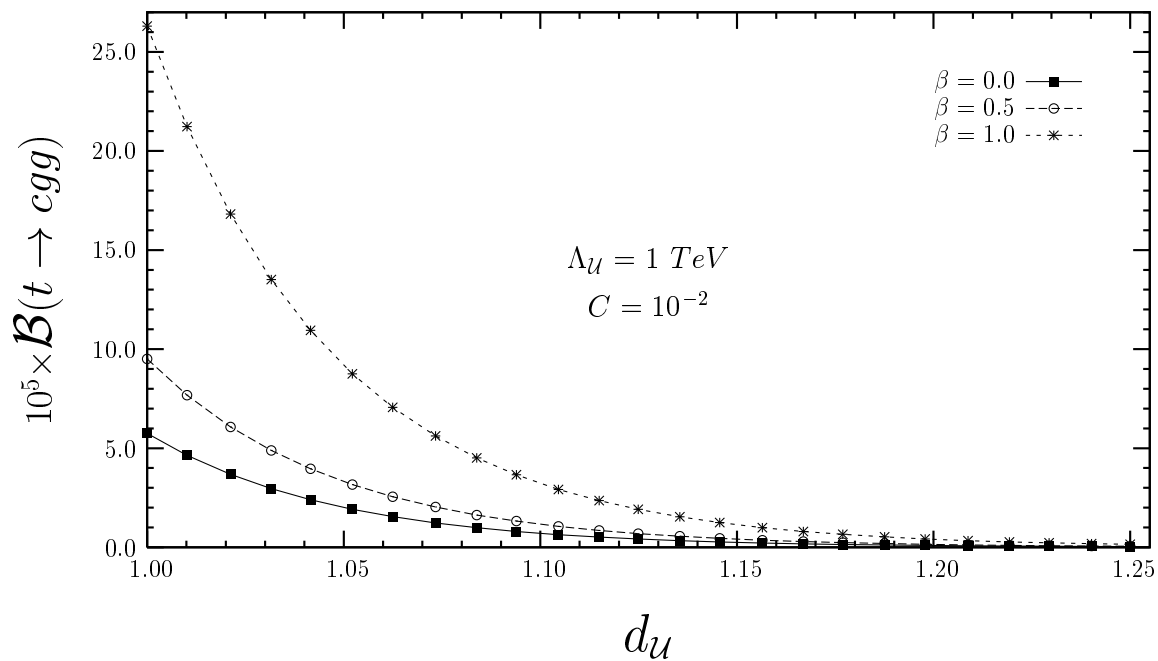


Figure 2:

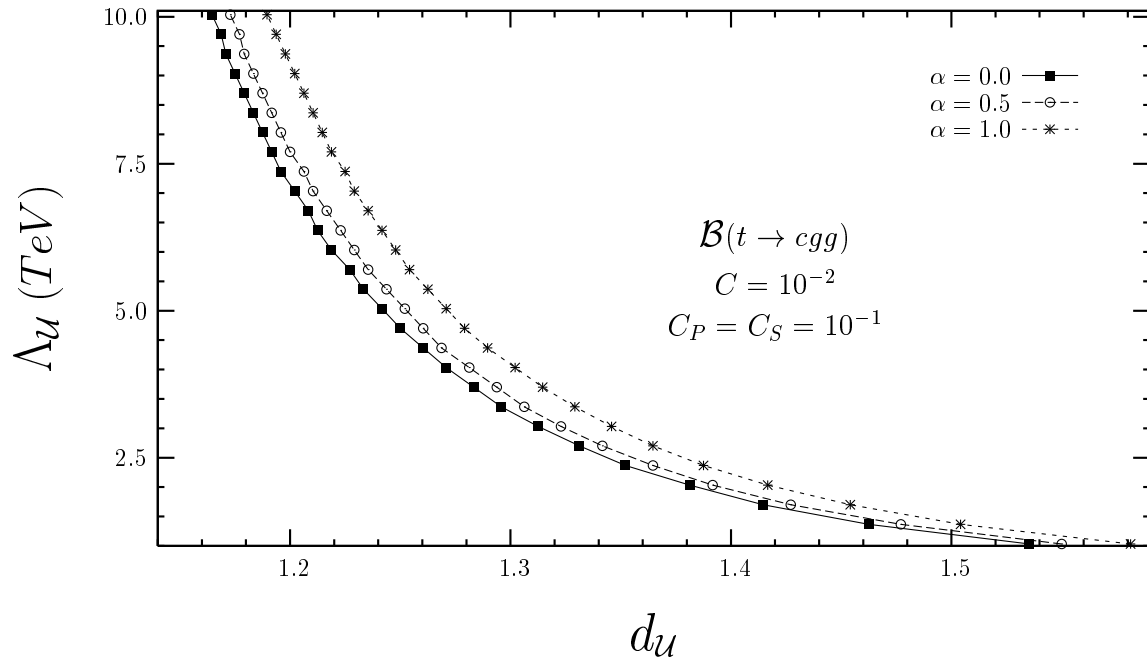


Figure 3:

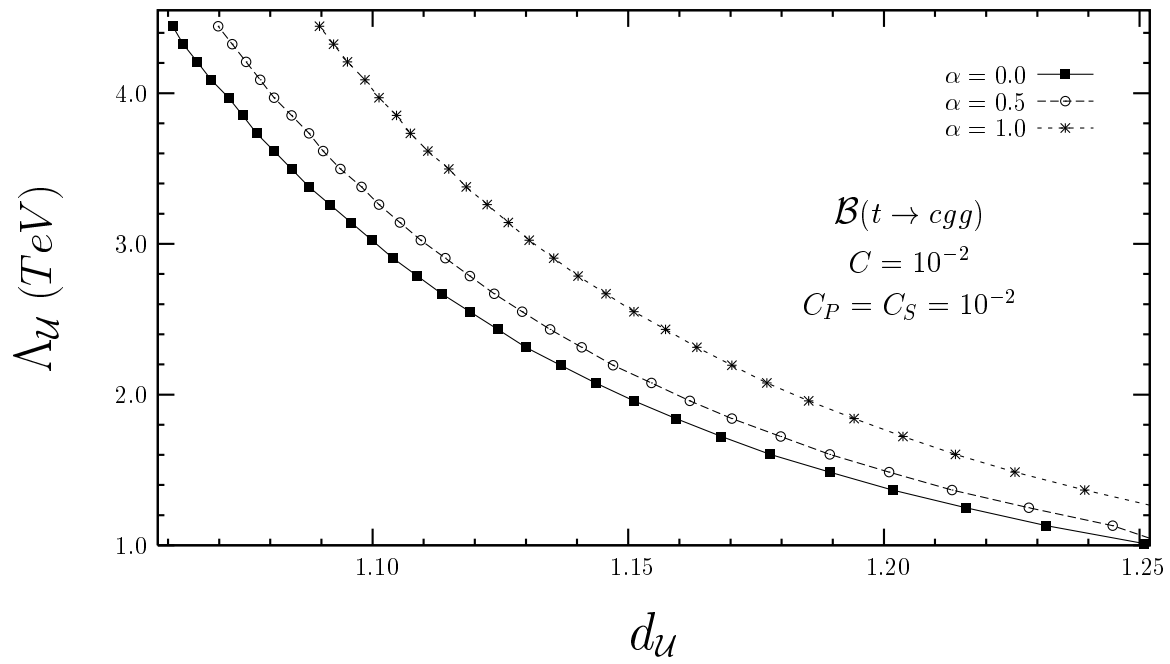


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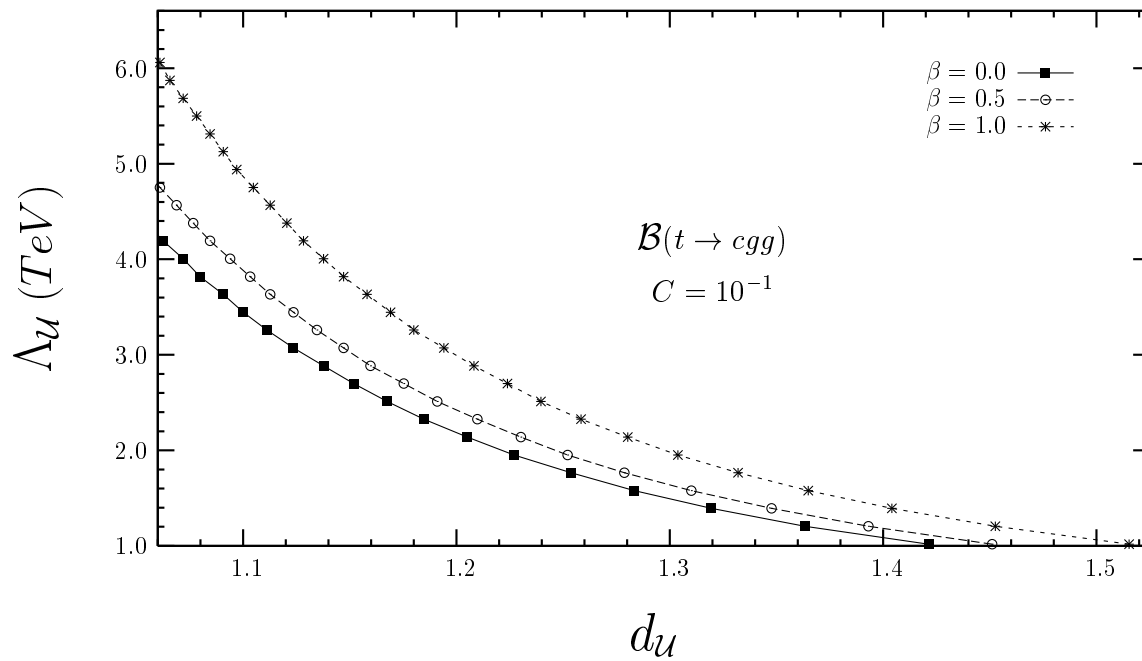


Figure 5:

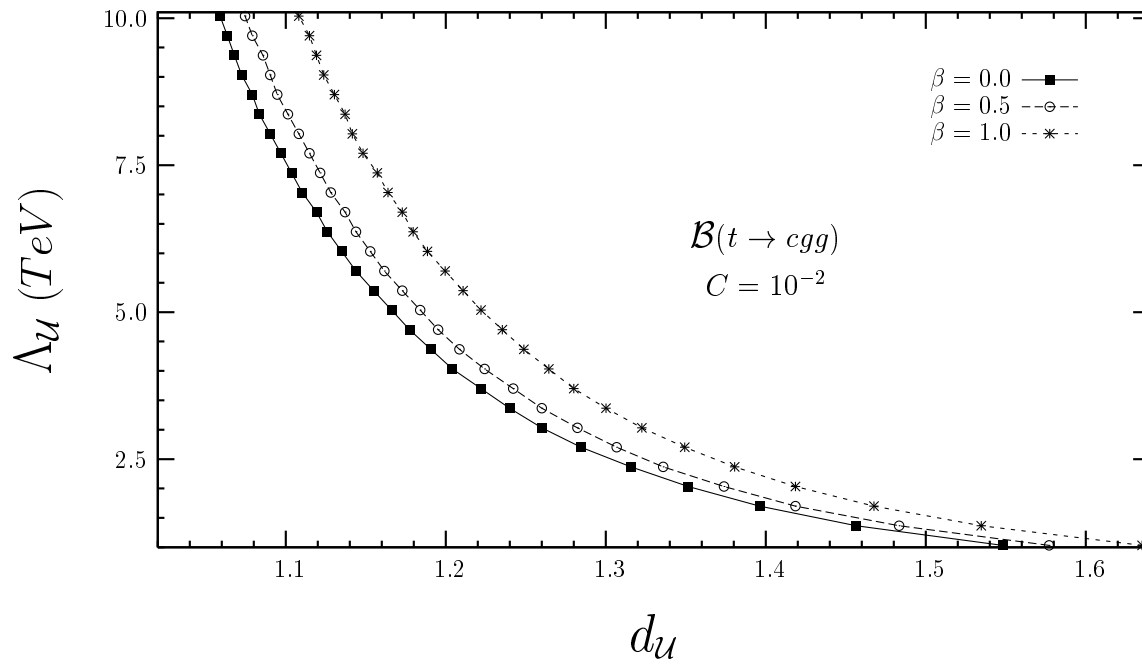


Figure 6:

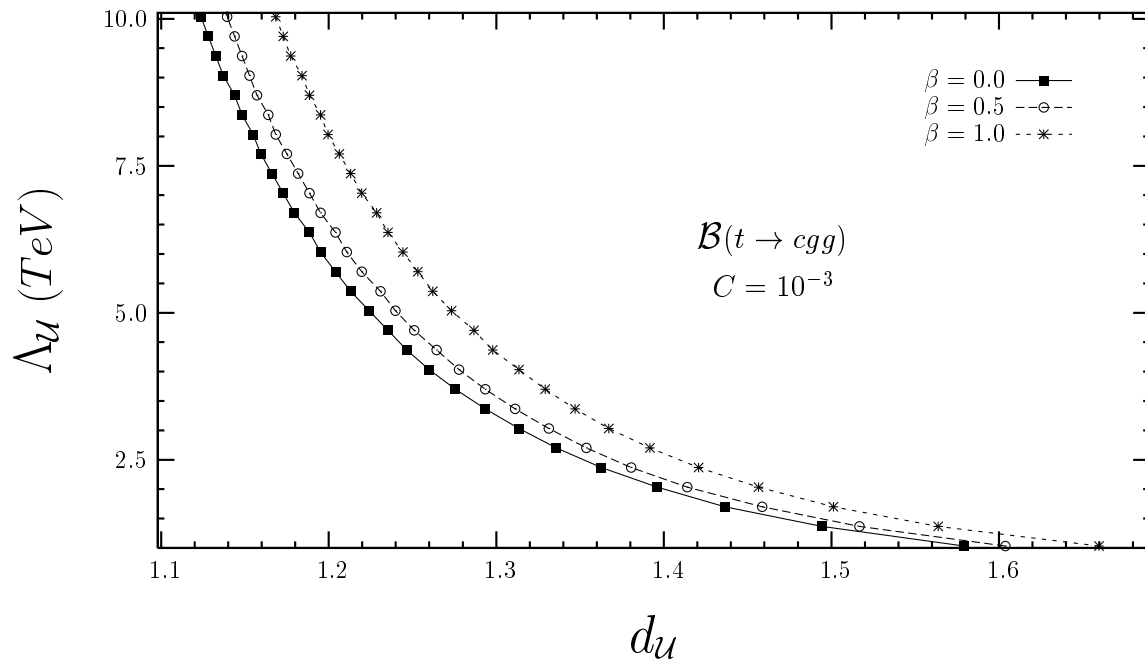


Figure 7: