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Unparticle physics effects in $\Lambda_b \to \Lambda + \textit{missing energy}$ processes

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Abstract

We study unparticle physics effects in $\Lambda_b \to \Lambda + \textit{missing energy}$ decay with polarized Λ_b and Λ baryons. The sensitivity of the branching ratio of this decay and polarizations of Λ_b and Λ baryons on the scale dimension $d_{\mathcal{U}}$ and effective cut-off parameter $\Lambda_{\mathcal{U}}$ are discussed.

1. Introduction

Flavor changing neutral current (FCNC) decays induced by the $b \to s$ transition are promising decays for checking predictions of the Standard Model (SM) at quantum loop level, since they are forbidden at tree level in the SM. These transitions are also very suitable in looking for new physics beyond the SM.

In the SM the $b \to s\nu\bar{\nu}$ decay receives special attention due to the theoretical advantage that uncertainties in this decay are much smaller compared to other FCNC decays due to the absence of photonic penguin diagrams and hadronic long distance effects. However, in spite of these theoretical advantages experimental measurement of this inclusive channel seems to be very difficult, because it requires a construction of all X_s. Therefore, experimentalists focus only on exclusive channels like $B \to K(K^*)\bar{\nu}\nu$. This channel studied extensively on theoretical grounds in many works (see for example [1–4]). Another class of decays, which is described by the $b \to s\bar{\nu}\nu$ transition at inclusive level, is the baryonic $\Lambda_b \to \Lambda\bar{\nu}\nu$ decay.

It should be noted that it is impossible to analyze the helicity structure of the effective Hamiltonian in B meson decays governed by $b \to s$ transition, since the information about chiralities of the quarks is lost in the hadronization process. In contrary to the mesonic decays, baryonic decays could access the helicity structure of the effective Hamiltonian for the $b \to s$ transition [5]. Therefore the heavy baryonic decays can be very rich for studying the polarization effects.

Radiative and semileptonic decays of Λ_b baryon, such as $\Lambda_b \to \Lambda\gamma$, $\Lambda_b \to \Lambda_c\ell\bar{\nu}_\ell$, $\Lambda_b \to \Lambda\ell^+\ell^-$ and $\Lambda_b \to \Lambda\bar{\nu}\nu$, are comprehensively studied in the framework of SM in many works [5–12]. Present status of the experimental investigations of heavy baryons is discussed in [13].

As has already been noted, FCNC transitions are very sensitive to the existence of new physics beyond the SM. One such model is the so-called unparticles is proposed by H. Georgi [14]. It is assumed in this model that, at very high energies the theory contains SM fields and the fields with a non-trivial infrared fixed point so that in the infrared limit it will be an asymptotic conformal theory and be scale-invariant, which are called Banks–Zaks (BZ) fields [15]. In unparticle physics model these two sectors interacted by exchange of particle with a large mass scale μ, below this scale where non-renormalizable couplings between SM and BZ fields will be induced and renormalizable couplings between the BZ fields are then produced by dimensional transmutation, and the scale-invariant unparticle fields emerge below a scale $\Lambda_{\mathcal{U}}$. In the effective theory, below $\Lambda_{\mathcal{U}}$, BZ operators match onto the

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unparticle operators, and non-renormalizable interactions between SM and unparticle operators can be obtained. An important result in this theory is that unparticle stuff with scale dimension $d_{\mathcal{U}}$ looks like a non-integer number $d_{\mathcal{U}}$ of invisible particles [14], where production might be detectable in missing energy and momentum distributions. Various phenomenological aspects of the unparticle physics have recently been extensively discussed in literature (see [16–21], and references therein).

In the present work we study the $\Lambda_b \to \Lambda + \textit{missing energy}$ decay in unparticle physics. The Letter is organized as follows: In Section 2, we give necessary theoretical framework to describe the differential decay width of $\Lambda_b \to \Lambda + \textit{missing energy}$ in the SM and in unparticle physics. Section 3 is devoted to numerical analysis and our concluding remarks are presented in Section 4.

2. Theoretical framework

In the SM $\Lambda_b \to \Lambda + \textit{missing energy}$ channel is described by the $\Lambda_b \to \Lambda\bar{\nu}\nu$ decay. As has already been noted, unparticles can also contribute to this decay. Therefore, a comparison of the signature of the decay modes $\Lambda_b \to \Lambda\bar{\nu}\nu$ and $\Lambda_b \to \Lambda + U$ is required.

In the SM, the $\Lambda_b \to \Lambda\bar{\nu}\nu$ decay is described at quark level by the $b \to s\bar{\nu}\nu$ transition and receives contributions from Z-penguin and box diagrams, where main contributions come from intermediate top quarks. The effective Hamiltonian responsible for $b \to s\bar{\nu}\nu$ transition is described by only one Wilson coefficient C_{10} and its explicit form is

$$\mathcal{H} = \frac{G_F}{\sqrt{2}} \frac{\alpha}{2\pi} V_{tb} V_{ts}^* C_{10} \bar{s}\gamma_\mu (1-\gamma_5) b \bar{\nu}\gamma^\mu (1-\gamma_5)\nu, \tag{1}$$

where G_F and α are the Fermi constant and structure constants, respectively, V_{ij} are the elements of the Cabibbo–Cobayashi–Maskawa matrix (CKM). The Wilson coefficient C_{10} in Eq. (1), including $\mathcal{O}(\alpha_s)$ corrections, has the following form:

$$C_{10} = \frac{X(x_t)}{\sin^2\theta}, \tag{2}$$

where

$$X(x_t) = X_0(x_t) + \frac{\alpha_s}{4\pi} X_1(x_t). \tag{3}$$

$X_0(x_t)$ in Eq. (3) is the usual Inami–Lim function [22] given by:

$$X_0(x_t) = \frac{x_t}{8}\left[\frac{x_t+2}{x_t-1} + \frac{3x_t-6}{(x_t-1)^2}\ln x_t\right], \tag{4}$$

and

$$\begin{aligned} X_1(x_t) = {} & \frac{4x_t^3 - 5x_t^2 - 23x_t}{3(x_t-1)^2} - \frac{x_t^4 + x_t^3 - 11x_t^2 + x_t}{(x_t-1)^3}\ln x_t + \frac{x_t^4 - x_t^3 - 4x_t^2 - 8x_t}{2(x_t-1)^3}\ln^2 x_t \\ & + \frac{x_t^3 - 4x_t}{(x_t-1)^2}\mathrm{Li}_2(1-x_t) + 8x_t \frac{\partial X_0(x_t)}{\partial x_t}\ln x_\mu. \end{aligned} \tag{5}$$

Here,

$$\mathrm{Li}_2(1-x_t) = \int_1^{x_t} dt \frac{\ln t}{1-t}$$

is the spence function, $x_t = m_t^2/m_W^2$, $x_\mu = \mu^2/m_W^2$ and μ describes the scale dependence when leading QCD corrections are taken into account. The function $X_1(x_t)$ is calculated in [23].

Similarly, at quark level in unparticle physics, $b \to s + \textit{missing energy}$ is described by the $b \to sU$ transition, where we shall consider two types of unparticle operators:

- Scalar unparticle operators,
- Vector unparticle operators.

For the scalar and vector operators $b \to sU$ transition is described by the following matrix elements

$$\frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\left[\bar{s}\gamma^\mu (C_S + C_P\gamma_5) b\right]\partial_\mu \mathcal{O}_u, \tag{6}$$

$$\frac{1}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}}\left[\bar{s}\gamma_\mu (C_V + C_A\gamma_5) b\right]\mathcal{O}_u^\mu, \tag{7}$$

where C_i are the dimensionless effective couplings.

Before performing analytical calculations, let us present the forms of the propagators for the scalar and vector unparticle physics [16,17]

$$\mathcal{D}(q^2) = \int d^4x\, e^{iqx} \langle 0|\mathrm{T}(\mathcal{O}_u(x)\mathcal{O}_u(0)|0\rangle,$$
$$= \frac{A_{d_\mathcal{U}}}{2\sin(d_\mathcal{U}\pi)} (-q^2)^{d_\mathcal{U}-2}, \tag{8}$$

$$\mathcal{D}_{\mu\nu} = \frac{A_{d_\mathcal{U}}}{2\sin(d_\mathcal{U}\pi)} (-q^2)^{d_\mathcal{U}-2} \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}\right), \tag{9}$$

where

$$A_{d_\mathcal{U}} = \frac{16\pi^{5/2}}{(2\pi)^{2d_\mathcal{U}}} \frac{\Gamma(d_\mathcal{U}+1/2)}{\Gamma(d_\mathcal{U}-1)\Gamma(2d_\mathcal{U})}. \tag{10}$$

It is found in [14] that, using scale invariance of the unparticle physics, the phase for an unparticle operator with the scale dimension $d_\mathcal{U}$ and momentum q is the same as the phase space for $d_\mathcal{U}$ invisible massless particles

$$d\Phi_\mathcal{U}(q) = A_{d_\mathcal{U}} \Theta(q^0)\Theta(q^2)(q^2)^{d_\mathcal{U}-2} \frac{d^4q}{(2\pi)^2}. \tag{11}$$

Having the explicit forms of the effective Hamiltonian at hand for the $b \to \bar{\nu}\nu$ transition and effective interaction for the $b \to sU$ transition, our next problem is computation of the matrix element of (1), (7) and (8), between initial and final state baryons. It follows from Eqs. (1) and (7) that the we need to know the following matrix elements

$$\langle \Lambda|\bar{s}\gamma_\mu b|\Lambda_b\rangle, \qquad \langle \Lambda|\bar{s}\gamma_\mu\gamma_5 b|\Lambda_b\rangle. \tag{12}$$

These matrix elements can be parametrized in terms of the form factors a follows [11]:

$$\langle \Lambda|\bar{s}\gamma_\mu b|\Lambda_b\rangle = \bar{u}_\Lambda\left[f_1\gamma_\mu + f_2 i\sigma_{\mu\nu}q^\nu + f_3 q_\mu\right]u_{\Lambda_b}, \tag{13}$$
$$\langle \Lambda|\bar{s}\gamma_\mu\gamma_5 b|\Lambda_b\rangle = \bar{u}_\Lambda\left[g_1\gamma_\mu\gamma_5 + g_2 i\sigma_{\mu\nu}q^\nu\gamma_5 + g_3 q_\mu\gamma_5\right]u_{\Lambda_b}, \tag{14}$$

where $q = p_{\Lambda_b} - p_\Lambda$.

It follows from these expressions that $\Lambda_b \to \Lambda + \textit{missing energy}$ decay is described in terms of numerous form factors. It is shown in [5] that Heavy Effective Quark Theory (HEQT) reduces the number of independent form factors to two (F_1 and F_2) irrelevant of the Dirac structure of corresponding operators, i.e.,

$$\langle \Lambda(p_\Lambda)|\bar{s}\Gamma b|\Lambda_b\rangle = \bar{u}_\Lambda\left[F_1(q^2) + \not{v}F_2(q^2)\right]\Gamma u_{\Lambda_b}, \tag{15}$$

where Γ is the arbitrary Dirac structure and $v^\mu = p^\mu_{\Lambda_b}/m_{\Lambda_b}$ is the four velocity of Λ_b. Comparing Eqs. (13), (14) and (15) one can easily obtain relations among the form factors (see also [11])

$$f_1 = g_1 = F_1 + \frac{m_\Lambda}{m_{\Lambda_b}} F_2, \qquad f_2 = g_2 = f_3 = g_3 = \frac{F_2}{m_{\Lambda_b}}, \tag{16}$$

which we will use in our numerical analysis.

Using Eqs. (6), (7), (8), (13), (14) and (16) we get the following expression for the matrix elements of the $\Lambda_b \to \Lambda + \textit{missing energy}$ decay:

$$\mathcal{M}^{(1)} = \frac{G_F\alpha}{2\sqrt{2}\pi} V_{tb}V_{ts}^* C_{10}\bar{u}_\Lambda\left[\gamma_\mu(f_1 - g_1\gamma_5) + \not{q}\gamma_\mu(f_2 - g_2\gamma_5)\right]u_{\Lambda_b}\bar{\nu}\gamma_\mu(1-\gamma_5)\nu, \tag{17}$$

$$\mathcal{M}^{(2)} = \frac{1}{\Lambda_\mathcal{U}^{d_\mathcal{U}}} \bar{u}_\Lambda[A + B\gamma_5]u_{\Lambda_b} O, \tag{18}$$

$$\mathcal{M}^{(3)} = \frac{1}{\Lambda_\mathcal{U}^{d_\mathcal{U}-1}} \bar{u}_\Lambda\left[\gamma_\mu(A^V + B^V\gamma_5) + (C^V + D^V\gamma_5)p_{\Lambda_b\mu}\right]u_{\Lambda_b} O^\mu. \tag{19}$$

Here, $i = 1$, $i = 2$ and $i = 3$ correspond to the SM, scalar operator and vector operator contributions, respectively, and

$$A = C_S\left[(m_{\Lambda_b} - m_\Lambda)f_1 + q^2 f_3\right], \qquad B = C_P\left[-(m_{\Lambda_b} + m_\Lambda)f_1 + q^2 g_3\right],$$
$$A_V = C_V\left[f_1 - (m_{\Lambda_b} + m_\Lambda)f_2\right], \qquad B_V = C_A\left[g_1 + (m_{\Lambda_b} + m_\Lambda)g_2\right],$$
$$E_V = 2C_V f_2, \qquad D_V = 2C_A g_2. \tag{20}$$

In the derivation of Eqs. (17), (18) and (19), the neutrinos are taken to be massless and we also use $f_3 = f_2$ and $g_3 = g_2$ (see Eq. (16)).

It should be remarked here that, in the $\Lambda_b \to \Lambda + U$ decay the unparticle carries invariant momentum-square q^2, and therefore, can redecay into SM particles. In this work we assume that it is long-lived since it only couples weakly to the SM particles.

Having obtained the matrix elements for the $\Lambda_b \to \Lambda + \textit{missing energy}$, the differential decay width can be calculated straightforwardly. As has already been noted, the polarization effects for the $\Lambda_b \to \Lambda + \textit{missing energy}$, are richer than compared to the corresponding mesonic decays, since polarization of Λ_b and Λ can be measured. In this connection few words about the polarizations of baryons are in order. In the Λ_b rest frame, the unit vectors along the longitudinal, normal and transversal components of Λ polarization are defined in the following way:

$$\vec{e}_L = \frac{\vec{p}_\Lambda}{|\vec{p}_\Lambda|}, \qquad \vec{e}_N = \vec{\xi}_{\Lambda_b} \times \vec{e}_L, \qquad \vec{e}_T = \vec{e}_L \times \vec{e}_N,$$

where $\vec{p}_\Lambda$ is the momentum of Λ baryon and $\vec{\xi}_{\Lambda_b}$ is the unit vector along the Λ_b baryon spin in its rest frame. In the rest frame of Λ_b baryon the differential decay width can be written as:

$$\frac{d\Gamma^{(i)}}{dE_\Lambda} = \left(\frac{d\Gamma_0^{(i)}}{dE_\Lambda}\right)\frac{1}{4}\left[1 + \frac{I_2^{(i)}}{I_1^{(i)}}\vec{e}_L \cdot \vec{\xi}_{\Lambda_b}\right]\left[1 + \vec{\mathcal{P}}_\Lambda^{(i)} \cdot \vec{\xi}_\Lambda\right], \tag{21}$$

where $i = 1, 2, 3$, and $d\Gamma_0^{(i)}/dE_\Lambda$ describes the unpolarized differential decay width.

In Eq. (21) $\vec{\mathcal{P}}_\Lambda^{(i)}$ is determined as follows:

$$\vec{\mathcal{P}}_\Lambda^{(i)} = \frac{1}{1 + \frac{I_2^{(i)}}{I_1^{(i)}}\vec{e}_L \cdot \vec{\xi}_{\Lambda_b}}\left[\left(\frac{I_3^{(i)}}{I_1^{(i)}} + \frac{I_4^{(i)}}{I_1^{(i)}}\vec{e}_L \cdot \vec{\xi}_{\Lambda_b}\right)\vec{e}_L + \frac{I_5^{(i)}}{I_1^{(i)}}\vec{e}_T + \frac{I_6^{(i)}}{I_1^{(i)}}\vec{e}_N\right]. \tag{22}$$

Note that $\Lambda_b \to \Lambda\bar{\nu}\nu$ decay with Λ_b and Λ polarizations is studied in [11]. Explicit expressions of $d\Gamma_0/dE_\Lambda$, $\vec{\mathcal{P}}_\Lambda$, I_2 and I_1 in the SM are given in [11], and therefore we do not present them in this work.

After simple calculation, the decay width due to the scalar operator takes the following form:

$$\frac{d\Gamma_0^{(2)}}{dE_\Lambda} = \frac{1}{2m_{\Lambda_b}}\frac{A_{d_\mathcal{U}}}{(\Lambda_\mathcal{U}^{d_\mathcal{U}})^2}\left(q^2\right)^{d_\mathcal{U}-2}\frac{|\vec{p}_\Lambda|}{(2\pi)^2}I_1^{(2)}, \tag{23}$$

where $q = p_{\Lambda_b} - p_\Lambda$, and $|\vec{p}_\Lambda|$ is the magnitude of the Λ baryon three-momentum, and

$$\begin{aligned}
I_1^{(2)} &= 4m_{\Lambda_b}\left[|A|^2(E_\Lambda + m_\Lambda) + |B|^2(E_\Lambda - m_\Lambda)\right],\\
I_2^{(2)} &= -8\,\mathrm{Re}\left[AB^*\right]m_{\Lambda_b}|\vec{p}_\Lambda|,\\
I_3^{(2)} &= -8\,\mathrm{Re}\left[AB^*\right]m_{\Lambda_b}|\vec{p}_\Lambda| \equiv I_2^{(2)},\\
I_4^{(2)} &= 4|A|^2\left[-|\vec{p}_\Lambda|^2 + m_\Lambda(E_\Lambda + m_\Lambda)\right]m_{\Lambda_b} + 4|B|^2\left[|\vec{p}_\Lambda|^2 - m_\Lambda(E_\Lambda - m_\Lambda)\right]m_{\Lambda_b},\\
I_5^{(2)} &= 4|A|^2 m_{\Lambda_b}(E_\Lambda + m_\Lambda) - 4|B|^2 m_{\Lambda_b}(E_\Lambda - m_\Lambda),\\
I_6^{(2)} &= 8|\vec{p}_\Lambda|m_{\Lambda_b}\,\mathrm{Im}\left[AB^*\right].
\end{aligned} \tag{24}$$

The coefficients $\vec{e}_L$, $\vec{e}_T$ and $\vec{e}_N$, in Eq. (22) corresponding to the longitudinal, transversal and normal polarization asymmetries of Λ can also be defined as:

$$\mathcal{P}_{\Lambda,j}^{(i)} = \frac{d\Gamma^{(i)}(\vec{\xi}_\Lambda \cdot \vec{e}_j = 1) - d\Gamma^{(i)}(\vec{\xi}_\Lambda \cdot \vec{e}_j = -1)}{d\Gamma^{(i)}(\vec{\xi}_\Lambda \cdot \vec{e}_j = 1) + d\Gamma^{(i)}(\vec{\xi}_\Lambda \cdot \vec{e}_j = -1)},$$

where $j = L, T, N$.

Performing similar calculations for unpolarized decay due to the vector operator, we get:

$$\frac{d\Gamma_0^{(3)}}{dE_\Lambda} = \frac{1}{2m_{\Lambda_b}}\frac{A_{d_\mathcal{U}}}{(\Lambda_\mathcal{U}^{d_\mathcal{U}-1})^2}\left(q^2\right)^{d_\mathcal{U}-2}\frac{|\vec{p}_\Lambda|}{(2\pi)^2}I_1^{(3)},$$

and the expressions of the functions entering into Eq. (22) for the vector operator case are as follows:

$$\begin{aligned}
I_1^{(3)} = 4\frac{m_{\Lambda_b}}{q^2}\Big\{&|D_V|^2 m_{\Lambda_b}^2(E_\Lambda - m_\Lambda)\left[(E_\Lambda - m_{\Lambda_b})^2 - q^2\right] + |E_V|^2 m_{\Lambda_b}^2(E_\Lambda + m_\Lambda)\left[(E_\Lambda - m_{\Lambda_b})^2 - q^2\right]\\
&+ |A_V|^2\left[-2m_{\Lambda_b}E_\Lambda^2 + E_\Lambda\left(2m_\Lambda^2 + 2m_{\Lambda_b}^2 + q^2\right) - m_\Lambda\left(2m_\Lambda m_{\Lambda_b} + 3q^2\right)\right]
\end{aligned}$$

$$
\begin{aligned}
&+|B_V|^2\big[-2m_{\Lambda_b}E_\Lambda^2+E_\Lambda\big(2m_\Lambda^2+2m_{\Lambda_b}^2+q^2\big)-m_\Lambda\big(2m_\Lambda m_{\Lambda_b}-3q^2\big)\big]\\
&+2\,\mathrm{Re}\big[A_V E_V^*\big]m_{\Lambda_b}(E_\Lambda+m_\Lambda)\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda-m_{\Lambda_b})-q^2\big]\\
&+2\,\mathrm{Re}\big[B_V D_V^*\big]m_{\Lambda_b}(E_\Lambda-m_\Lambda)\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda+m_{\Lambda_b})+q^2\big]\big\},
\end{aligned}
$$

$$
\begin{aligned}
I_2^{(3)}=-8\frac{m_{\Lambda_b}|\vec p_\Lambda|}{q^2}&\big\{\mathrm{Re}\big[E_V D_V^*\big]m_{\Lambda_b}^2\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]+\mathrm{Re}\big[A_V B_V^*\big]\big[2m_\Lambda^2-2m_{\Lambda_b}E_\Lambda+q^2\big]\\
&+\mathrm{Re}\big[B_V E_V^*\big]m_{\Lambda_b}\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda+m_{\Lambda_b})+q^2\big]+\mathrm{Re}\big[A_V D_V^*\big]m_{\Lambda_b}\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda-m_{\Lambda_b})-q^2\big]\big\},
\end{aligned}
$$

$$
\begin{aligned}
I_3^{(3)}=-8\frac{m_{\Lambda_b}|\vec p_\Lambda|}{q^2}&\big\{\mathrm{Re}\big[E_V D_V^*\big]m_{\Lambda_b}^2\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]+\mathrm{Re}\big[A_V D_V^*\big]m_{\Lambda_b}\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda-m_{\Lambda_b})-q^2\big]\\
&+\mathrm{Re}\big[B_V E_V^*\big]m_{\Lambda_b}\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda+m_{\Lambda_b})+q^2\big]+\mathrm{Re}\big[A_V B_V^*\big]\big[2m_{\Lambda_b}(E_\Lambda-m_{\Lambda_b})-q^2\big]\big\},
\end{aligned}
$$

$$
\begin{aligned}
I_4^{(3)}=4\frac{m_{\Lambda_b}}{m_\Lambda q^2}&\big\{|D_V|^2m_{\Lambda_b}^2\big[|\vec p_\Lambda|^2-m_\Lambda(E_\Lambda-m_\Lambda)\big]\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]\\
&+|E_V|^2m_{\Lambda_b}^2\big[-|\vec p_\Lambda|^2+m_\Lambda(E_\Lambda+m_\Lambda)\big]\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]\\
&+2\,\mathrm{Re}\big[A_V E_V^*\big]m_{\Lambda_b}\big[-|\vec p_\Lambda|^2+m_\Lambda(E_\Lambda+m_\Lambda)\big]\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda-m_{\Lambda_b})-q^2\big]\\
&+2\,\mathrm{Re}\big[B_V D_V^*\big]m_{\Lambda_b}\big[|\vec p_\Lambda|^2-m_\Lambda(E_\Lambda-m_\Lambda)\big]\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda+m_{\Lambda_b})+q^2\big]\\
&+|B_V|^2\big[2m_\Lambda(E_\Lambda-m_{\Lambda_b})\big(m_{\Lambda_b}E_\Lambda-m_\Lambda^2\big)+2|\vec p_\Lambda|^2\big(m_\Lambda^2+m_\Lambda m_{\Lambda_b}-m_{\Lambda_b}(E_\Lambda-m_{\Lambda_b})\big)\\
&-q^2\big(|\vec p_\Lambda|^2-m_\Lambda(E_\Lambda-m_\Lambda)\big)\big]-|A_V|^2\big[|\vec p_\Lambda|^2\big(2m_\Lambda^2-2m_{\Lambda_b}(E_\Lambda+m_\Lambda)+2m_{\Lambda_b}^2-q^2\big)\\
&+m_\Lambda\big(2E_\Lambda^2m_{\Lambda_b}+2m_\Lambda^2m_{\Lambda_b}+m_\Lambda q^2-2E_\Lambda\big(m_\Lambda^2+m_{\Lambda_b}^2\big)+E_\Lambda q^2\big)\big]\big\},
\end{aligned}
$$

$$
\begin{aligned}
I_5^{(3)}=4\frac{m_{\Lambda_b}}{q^2}&\big\{-|D_V|^2m_{\Lambda_b}^2(E_\Lambda-m_\Lambda)\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]+|E_V|^2m_{\Lambda_b}^2(E_\Lambda+m_\Lambda)\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]\\
&-|A_V|^2\big[2m_{\Lambda_b}E_\Lambda^2+m_\Lambda\big(2m_\Lambda m_{\Lambda_b}+q^2\big)-E_\Lambda\big(2m_\Lambda^2+2m_{\Lambda_b}^2-q^2\big)\big]\\
&-|B_V|^2\big[-2m_{\Lambda_b}E_\Lambda^2+E_\Lambda\big(2m_\Lambda^2+2m_{\Lambda_b}^2-q^2\big)-m_\Lambda\big(2m_\Lambda m_{\Lambda_b}-q^2\big)\big]\\
&+2\,\mathrm{Re}\big[A_V E_V^*\big]m_{\Lambda_b}(E_\Lambda+m_\Lambda)\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda-m_{\Lambda_b})-q^2\big]\\
&-2\,\mathrm{Re}\big[B_V D_V^*\big]m_{\Lambda_b}(E_\Lambda-m_\Lambda)\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda+m_{\Lambda_b})+q^2\big]\big\},
\end{aligned}
$$

$$
\begin{aligned}
I_6^{(3)}=8\frac{m_{\Lambda_b}|\vec p_\Lambda|}{q^2}&\big\{\mathrm{Im}\big[E_V D_V^*\big]m_{\Lambda_b}^2\big[(E_\Lambda-m_{\Lambda_b})^2-q^2\big]+\mathrm{Im}\big[A_V B_V^*\big]\big[2m_{\Lambda_b}(E_\Lambda-m_{\Lambda_b})+q^2\big]\\
&-\mathrm{Im}\big[B_V E_V^*\big]m_{\Lambda_b}\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda+m_{\Lambda_b})+q^2\big]+\mathrm{Im}\big[A_V D_V^*\big]m_{\Lambda_b}\big[(E_\Lambda-m_{\Lambda_b})(m_\Lambda-m_{\Lambda_b})-q^2\big]\big\}.
\end{aligned}
$$

For the case when Λ_b is unpolarized, we get from Eq. (22) that

$$\vec{\mathcal{P}}^{(i)}_{\Lambda,L}=\alpha_\Lambda^{(i)}\vec e_L,$$

with

$$\alpha_\Lambda^{(i)}=\frac{I_3^{(i)}}{I_1^{(i)}},$$

which means that, in this case Λ polarization is purely longitudinal.

For Λ unpolarized, by performing summation over Λ spin in Eq. (22), we get

$$\frac{d\Gamma^{(i)}}{dE_\Lambda}=\left(\frac{d\Gamma_0^{(i)}}{dE_\Lambda}\right)\frac{1}{2}\big[1+\alpha_{\Lambda_b}^{(i)}\vec\xi_{\Lambda_b}\cdot\vec e_L\big],$$

where

$$\alpha_{\Lambda_b}^{(i)}=\frac{I_2^{(i)}}{I_1^{(i)}}.$$

Note that the normal component $\mathcal{P}^{(i)}_{\Lambda,N}$ of Λ polarization is a T-odd quantity and its non-zero value indicates CP violation. In the SM and considered version of unparticle physics, there is no CP violation. This is due to the fact that in both models the process is described by a single weak amplitude. Obviously, at least two different weak amplitudes are needed for CP violation, i.e., in addition to the existing mechanism there should exist a new one.

Table 1
Form factors for $\Lambda_b \to \Lambda$ transition in a three parameter fit

	$F(0)$	a_F	b_F
$F_1(0)$	0.462	−0.0182	−0.000176
$F_2(0)$	−0.077	−0.0685	0.001460

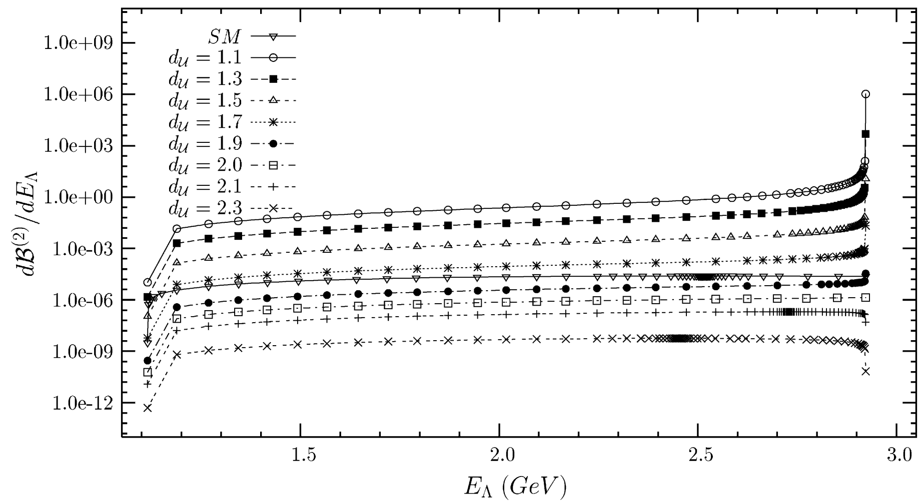

Fig. 1. The dependence of the unpolarized differential decay width of the $\Lambda_b \to \Lambda + \textit{missing energy}$ decay on Λ baryon energy at $\Lambda_{\mathcal{U}} = 1$ TeV, and at fixed values of $d_{\mathcal{U}}$ for the scalar operator.

3. Numerical analysis

In this section we calculate the numerical values of the differential branching ratio and polarizations of $\Lambda_b \to \Lambda + \textit{missing energy}$ decay in unparticle physics.

The transition form factors f_i and g_i, as well as $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}$, are the main input parameters in the numerical analysis. For the form factors we use the results of [24] in which QCD sum rules method together with HQET, which reduces the number of independent form factors to two, is used. The q^2 dependence of F_i in terms of three-parameter fit has the form [24]

$$F_i(q^2) = \frac{F_i(0)}{1 - a_F^i(q^2/m_{\Lambda_b}^2) + b_F^i(q^2/m_{\Lambda_b}^2)^2}.$$

The values of $F_i(0)$, a_F^i and b_F^i are given in Table 1.

It is emphasized in [14] that unparticles behave as a non-integer number of particles and it is shown there that the very peculiar shape of u-quark energy distribution in the $t \to c\mathcal{U}$ decay and it can serve as a good test in discovering unparticles experimentally. Along the same lines, the energy distribution of K and K^* mesons in the $B \to K(K^*) + \textit{missing energy}$ decay is analyzed [20] and it is seen that this decay, especially in the presence of vector unparticle operators, is very distinctive compared to that of the SM prediction. Similar situation can take place for the $\Lambda_b \to \Lambda + \textit{missing energy}$ decay. In what follows, we try to answer the intriguing question whether the polarization observables can be useful for the experimental observation of unparticles.

It is shown in [18] that if vector operators couple to the flavor non-diagonal current, $d_{\mathcal{U}}$ should be larger than $d_{\mathcal{U}} > 2$. On the other hand, the bound for the scalar operators turns out to be $d_{\mathcal{U}} > 1$, and these are the bounds we will use in our numerical analysis. It should be noted here that, the fundamental reason why $d_{\mathcal{U}} > 2$ for the vector particles, is due to the fact that, four-dimensional conformal group admits unitary representations only for $d_{\mathcal{U}} > 1 + j_L + j_R$, where $j_{L(R)}$ are the $SL(2, C)$ spins [25], and it follows from this expression that for vector particle $d_{\mathcal{U}} > 2$.[2]

The values of other parameters are chosen as $C_P = C_S = 2.0 \times 10^{-3}$ for scalar operators; and $C_V = C_A = 10^{-5}$ for vector operator, and $\Lambda_{\mathcal{U}} = 1$ TeV for both cases.

In Figs. 1 and 2, we present the dependence of the differential decay width as a function of the Λ baryon energy E_Λ for scalar and vector operators, for various choices of $d_{\mathcal{U}}$, respectively. From these figures we see that the distribution for the final Λ baryon energy for both operators are similar. Note that, the behavior of the dependence of the differential decay width on the energy distribution E_Λ for these operators and SM case are very similar to each other.

2 We thank the referee for calling our attention to this point.

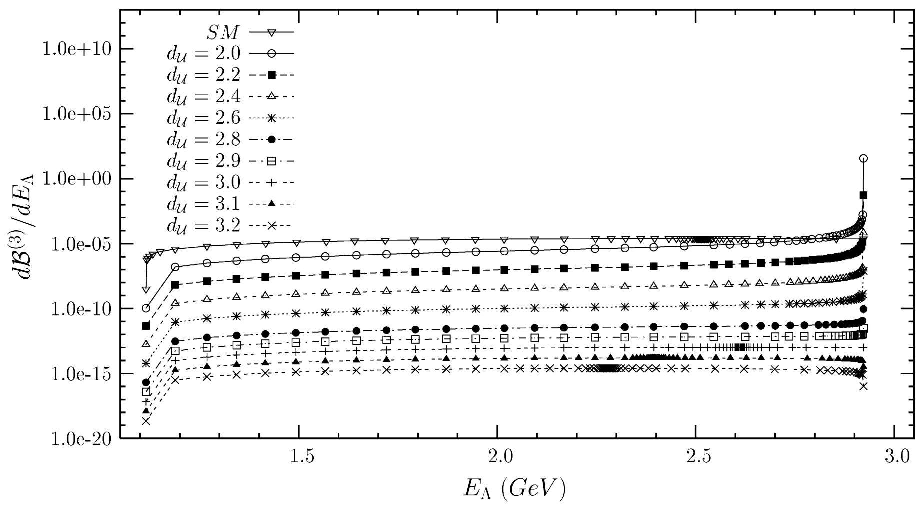

Fig. 2. The same as in Fig. 1, but for the vector operator.

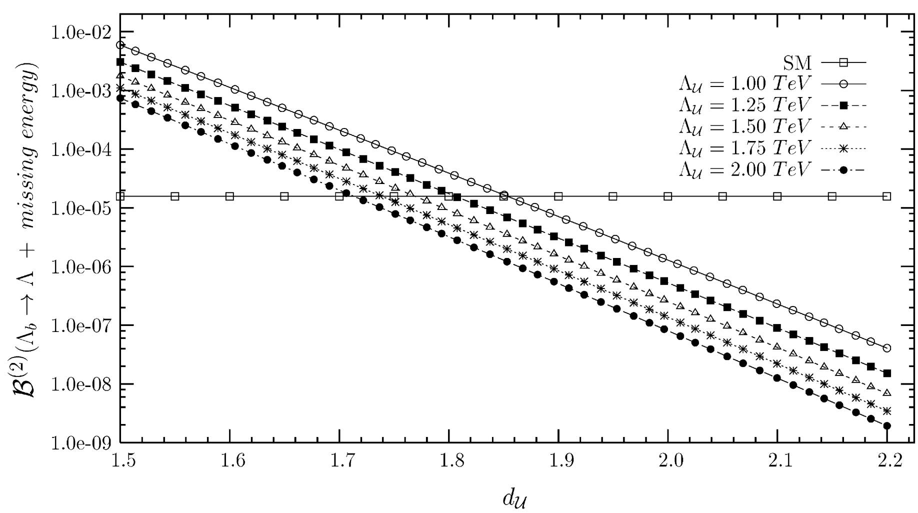

Fig. 3. The dependence of the branching ratio of $\Lambda_b \to \Lambda + missing\ energy$ decay on $d_{\mathcal{U}}$ at fixed values of $\Lambda_{\mathcal{U}}$ for the scalar operator.

In Figs. 3 and 4 we present the dependence of the branching ratios of the $\Lambda_b \to \Lambda + missing\ energy$ decay on $d_{\mathcal{U}}$ at fixed values of the effective coupling constants C_S, C_P, C_V and C_A, at different values of the cut-off scale $\Lambda_{\mathcal{U}}$, for scalar and vector operators. For completeness, in these figures we also present the SM result for the branching ratio of the $\Lambda_b \to \Lambda\bar{\nu}\nu$ decay. From these figures we see that, for $d_{\mathcal{U}} > 2$, the value of the branching ratio is smaller compared to that of the SM case in the presence of the vector operator; while the branching ratio can exceed the SM prediction in the presence of the scalar operator when $d_{\mathcal{U}} < 1.7$, whose behavior is determined by $\Lambda_{\mathcal{U}}$. Therefore, determination of the value of the branching ratio can put stringent restrictions to the values of $d_{\mathcal{U}}$ and $\Lambda_{\mathcal{U}}$, under the assumption that $C_S = C_P$ and $C_V = C_A$.

In Figs. 5 and 6, we present the dependence of the branching ratio of the $\Lambda_b \to \Lambda + missing\ energy$ decay on cut-off scale $\Lambda_{\mathcal{U}}$ at fixed values of C_S, C_P, C_V and C_A, respectively. We observe from these figures that, the branching ratios are very sensitive to the values of these effective couplings. It follows from Fig. 5 that, for the scalar operator case, under the assumption $C_S = C_P$, and up to $\Lambda_{\mathcal{U}} = 8$ TeV, the value of the branching ratio exceeds that of the SM prediction, for all choices of the fixed values of $d_{\mathcal{U}}$, which can give useful information about unparticle physics.

From all these figures we see that the branching ratios are very sensitive to the value of the parameter $d_{\mathcal{U}}$, and when $d_{\mathcal{U}} > 2$, the branching ratio of $\Lambda_b \to \Lambda + missing\ energy$ decay due to the scalar and vector operators are less than the SM prediction. Therefore at $d_{\mathcal{U}} > 2$, the measurement of the branching ratio cannot provide useful information about the existence of unparticle physics. In this connection there follows the question whether we can establish new physics by studying the effects due to the polarizations of the Λ and Λ_b baryons, or not. In other words, are there regions of the parameter $d_{\mathcal{U}}$ larger than 2, where the branching ratio is smaller compared to that of the SM prediction, while the polarization effects differ from that of the SM results?

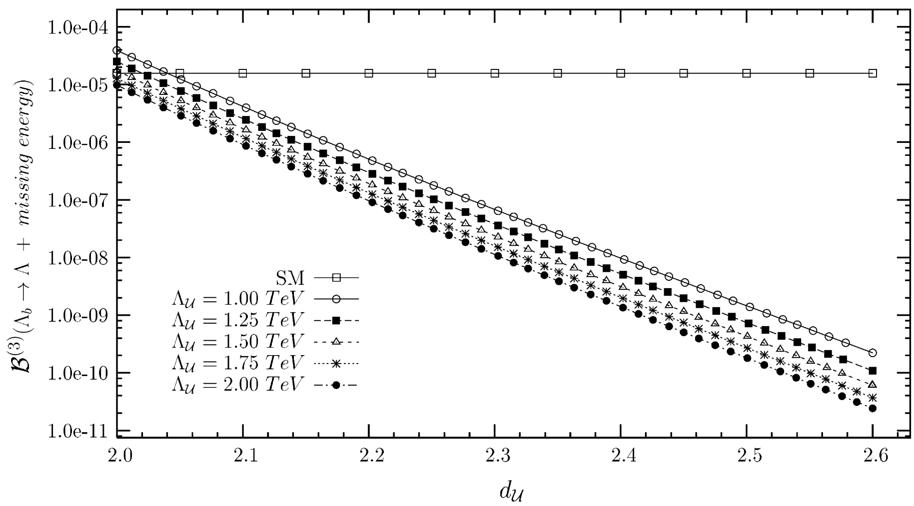

Fig. 4. The same as in Fig. 3, but for the vector operator.

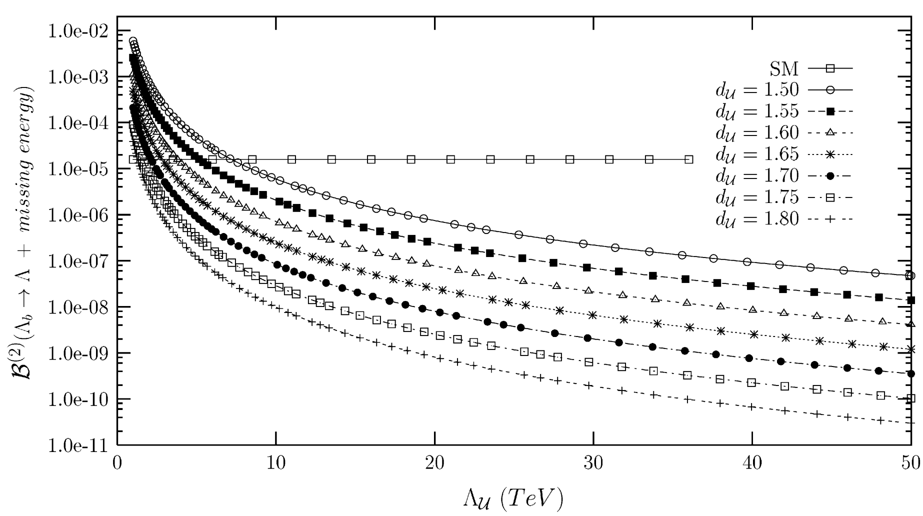

Fig. 5. The dependence of the branching ratio of $\Lambda_b \to \Lambda + missing\ energy$ decay on $\Lambda_{\mathcal{U}}$ at fixed values of $d_{\mathcal{U}}$ for the scalar operator.

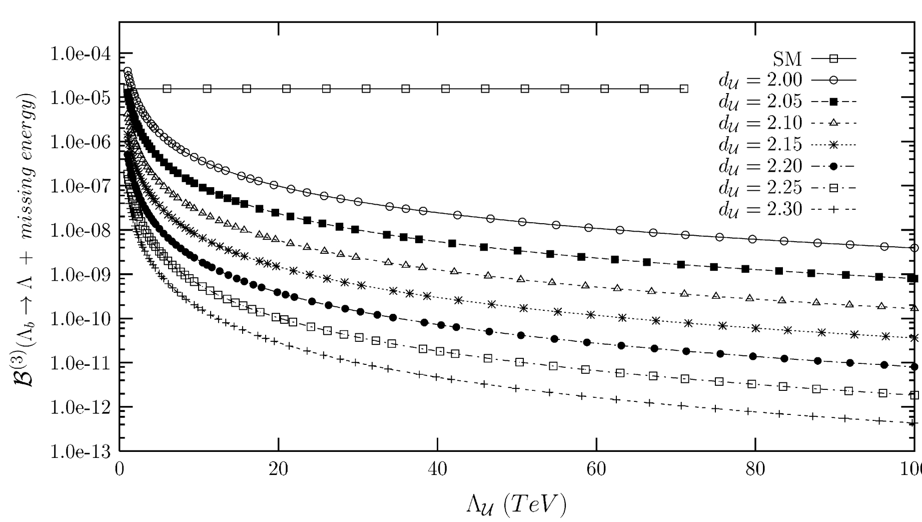

Fig. 6. The same as in Fig. 5, but for the vector operator.

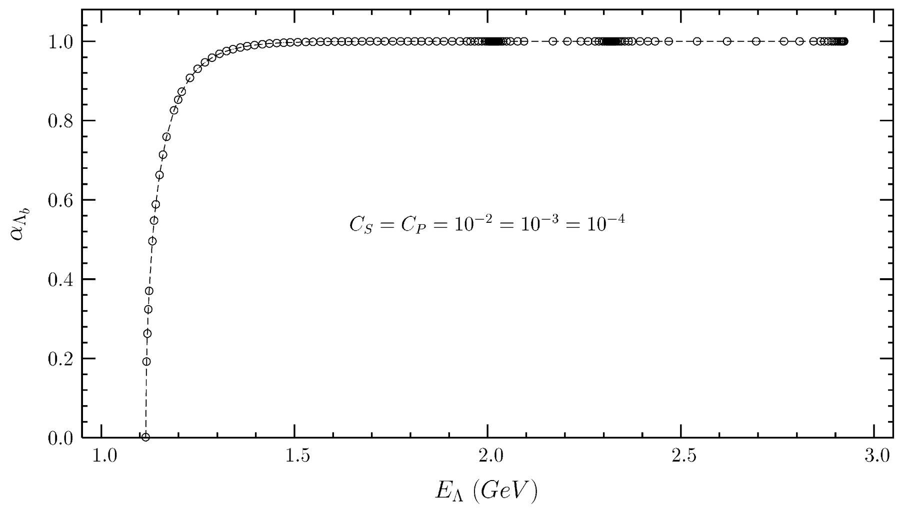

Fig. 7. The dependence of the Λ_b baryon polarization α_{Λ_b} on Λ baryon energy E_Λ at $\Lambda_{\mathcal{U}} = 1$ TeV, and at fixed values of C_S and C_P for the scalar operator.

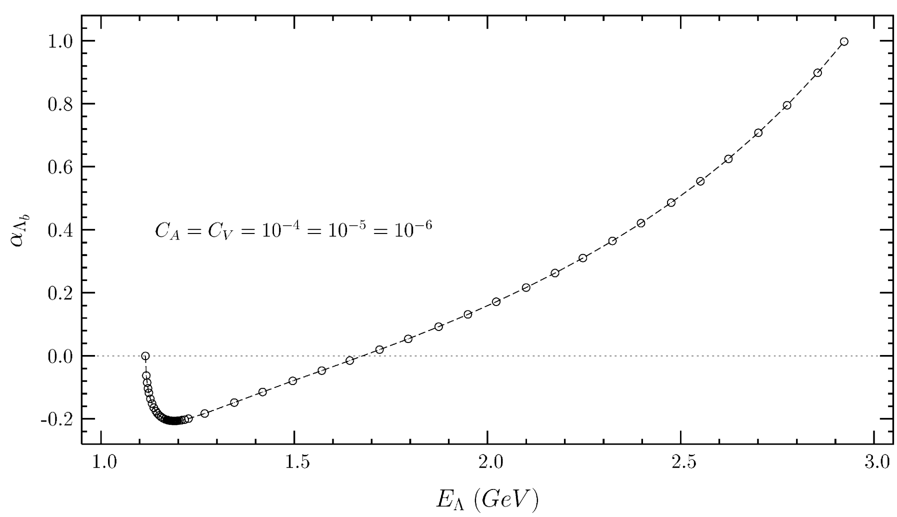

Fig. 8. The same as in Fig. 7, but for the vector operator.

In order to answer this question we study the effects due to the Λ and Λ_b baryon polarizations. As we proceed in analyzing the Λ_b polarizations we assume that Λ is not polarized, and when we analyze Λ polarizations we have assumed that $(I_2^{(i)}/I_1^{(i)})\vec{e}_L \cdot \vec{\xi}_{\Lambda_b}$ is small which can be neglected in numerical calculations.

In Figs. 7 and 8 we present the dependence of $\alpha_{\Lambda_b} = I_2/I_1$ on the Λ baryon energy, for scalar and vector unparticle operators, respectively.

In the presence of the scalar operator the magnitude of α_{Λ_b} starts from zero and approaches sharply to the value one; and from $E_\Lambda \simeq 1.25$ GeV on, its value continues to be very close to one. It is further observed that in the whole physical region of E_Λ, its value is independent of the values of the parameters C_P and C_S. Averaged value of α_{Λ_b} in unparticle physics model is equal to 0.98.

In the vector operator case the situation is drastically different from that above-mentioned scalar operator case. In other words, in the region $(E_\Lambda)_{\min} \leqslant E_\Lambda \leqslant 1.70$ GeV, α_{Λ_b} gets negative values and at $E_\Lambda = 1.70$ GeV, α_{Λ_b} becomes zero. Starting from $E_\Lambda = 1.70$ GeV on, α_{Λ_b} increases with the increasing values of E_Λ. Similar situation occurs for the SM case as well (see [11]). More essential than that, similar to the scalar operator case, α_{Λ_b} is insensitive to the values of the parameters C_P and C_A. Note that in the SM $\langle \alpha_{\Lambda_b} \rangle = -0.33$, while this model predicts $\langle \alpha_{\Lambda_b} \rangle = 0.86$. Therefore the study of α_{Λ_b} on E_Λ, as well as measuring the average value of $\langle \alpha_{\Lambda_b} \rangle$ can play significant role for establishing unparticle physics.

Let us pay our attention to the study of the Λ baryon polarization. Dependence of the longitudinal polarization of Λ baryon on the Λ baryon energy in presence of the scalar and vector unparticle operators are presented in Figs. 9 and 10, respectively. From these figures we see that:

In the presence of scalar and vector operators, longitudinal polarization of Λ exhibits practically the same behavior, namely, up to $E_\Lambda = 1.25$ GeV, $\mathcal{P}_{\Lambda,L}$ increases and from that point on it remains constant for all kinematical region, and its value is independent of the parameters C_S, C_P, C_V and C_A.

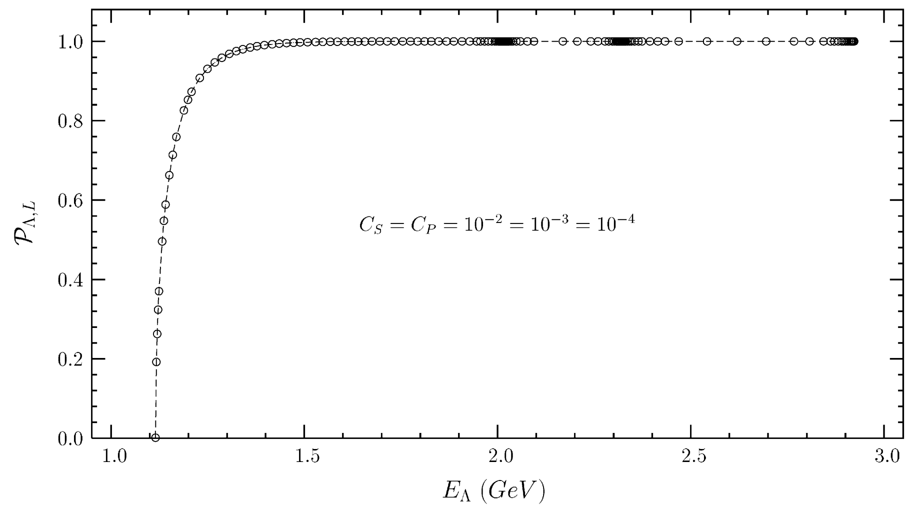

Fig. 9. The dependence of the distribution of longitudinal polarization $\mathcal{P}_{\Lambda,L}$ of Λ baryon on E_Λ at fixed values of C_S and C_P for the scalar operator.

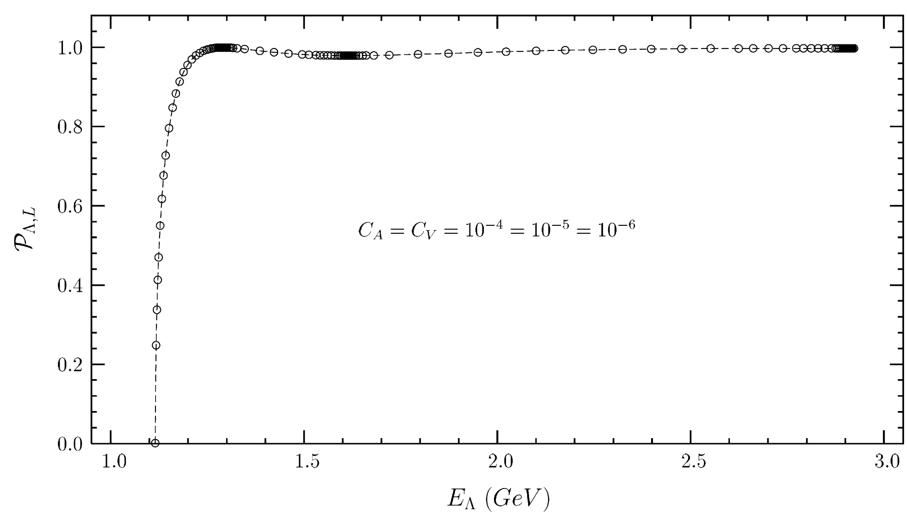

Fig. 10. The same as in Fig. 9, but for the vector operator.

We also calculate the averaged value of the longitudinal polarization of the Λ baryon in the scalar and vector operator cases, and obtain that $\langle \mathcal{P}_{\Lambda,L}\rangle \approx 0.98$ and $\langle \mathcal{P}_{\Lambda,L}\rangle \approx 0.99$, respectively, while in the SM $\langle P_{\Lambda,L}\rangle \approx -0.3$.

Transversal polarization of Λ decreases with increasing values of E_Λ and from $E_\Lambda = 1.75$ GeV on it approaches to zero in the presence of the scalar operator, being insensitive to the parameters C_S and C_P (see Fig. 11). This behavior is very different in the vector operator case. While the sign of $\mathcal{P}_{\Lambda,T}$ is negative up to $E_\Lambda = 1.25$ GeV, it starts increasing from this point on, and attains at constant value 10% after $E_\Lambda = 1.5$ GeV. Similar to the scalar operator case, $\mathcal{P}_{\Lambda,T}$ is insensitive to the numerical values of C_V and C_A (see Fig. 12).

Our calculations leads to the result that, the averaged value of the transversal polarization is $\langle \mathcal{P}_{\Lambda,T}\rangle = 7\%$ in the vector and $\langle \mathcal{P}_{\Lambda,T}\rangle = 4.6\%$ in the scalar operator case, respectively, and $\langle \mathcal{P}_{\Lambda,T}\rangle = 5.4\%$ in the SM case.

Note that, all these results on the polarization effects of the Λ and Λ_b baryons are practically independent of the value of the parameter $d_{\mathcal{U}}$, under the assumption of the equality of the coupling constants C_S and C_P, as well as C_V and C_A. From these results we can deduce that the polarization effects can play quite an essential role in establishing unparticle physics.

Few words on the assumption about equality of the coupling constants $C_S = C_P$ and $C_V = C_A$ are in order. In order to analyze the sensitivity of the branching ratio to this assumption, we consider the deviation of the branching ratio from the case when the above mentioned coupling constants are equal to each other. In Figs. 13 and 14, we depict the dependence of the ratio

$$\epsilon = \frac{\mathcal{B}(\Lambda_b \to \Lambda + U)_{\alpha\neq 1}}{\mathcal{B}(\Lambda_b \to \Lambda + U)_{\alpha=1}}, \tag{25}$$

on α, where $\alpha = C_P/C_S(C_A/C_V)$ for scalar (vector) unparticles. We observe from these figures that ϵ changes from 0.3 up to 10, when α changes between 0 and 5, for both type of couplings. In other words, the deviation of the branching ratio from its value

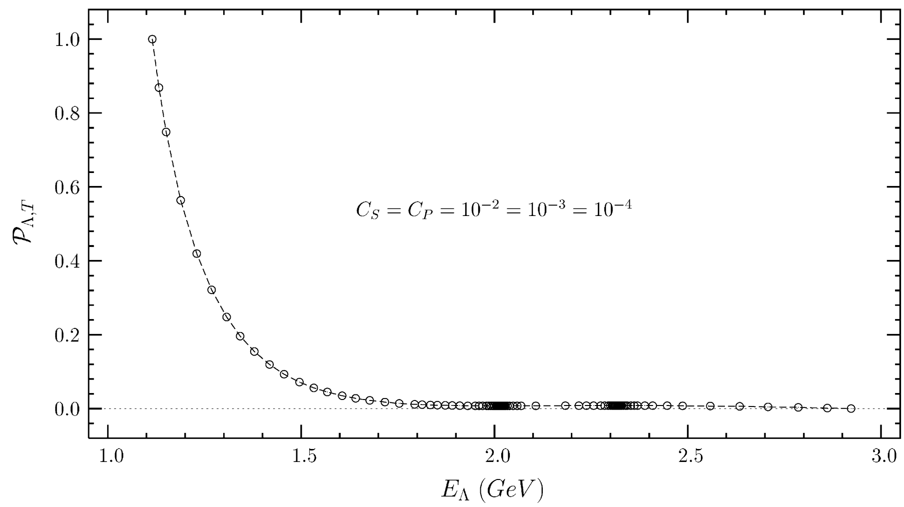

Fig. 11. The same as in Fig. 9, but for the transversal polarization $\mathcal{P}_{\Lambda,T}$ of Λ baryon.

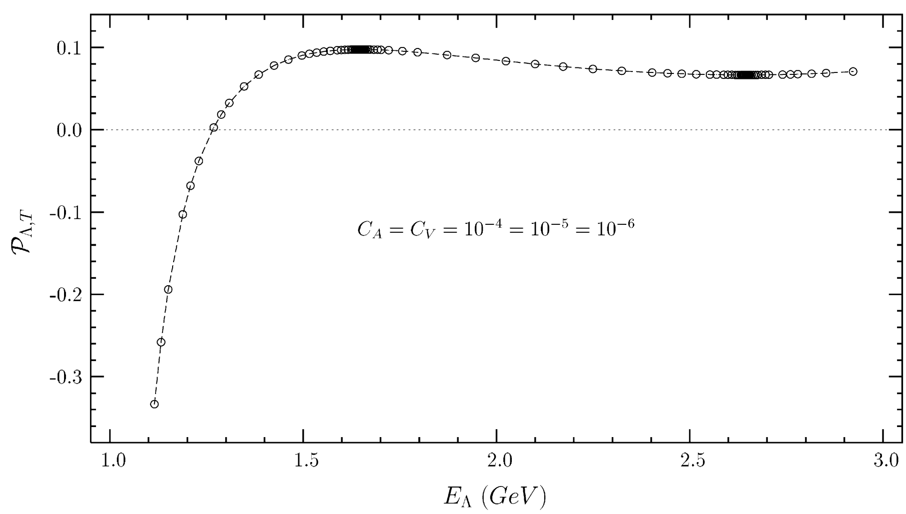

Fig. 12. The same as in Fig. 10, but for the vector operator.

for the case $\alpha = 1$ is not drastic at all. For this reason, the comments which we have made under the assumption of equality of the coupling constants, practically, remain valid.

At the end of this section, let us briefly discuss the prospects in measuring the $\Lambda_b \to \Lambda +$ *missing energy* channel at LHC. Note that, about 10^{10}–10^{11} b-quarks will be produced at LHC, at the total luminosity $L = 2 \times 10^{32}$ cm^{-2} s^{-1}. Taking into account that the fragmentation of b-quark to Λ_b is about 0.1 (this result is announced by the OPAL Collaboration), about 10^9–10^{10} Λ_b baryons are expected to be produced per year. Assuming 1000 events are enough for detecting the corresponding decay channel, $\Lambda_b \to \Lambda +$ *missing energy* channel with a branching ratio up to 10^{-7} seems to be quite measurable at LHC.

4. Conclusion

In the present work we have studied the possible manifestation of unparticles on the missing energy signatures of the rare Λ_b decays. The branching ratio, Λ and Λ_b baryon polarizations are studied in unparticle physics in the presence of scalar and vector operators. We obtain that the energy distribution of the Λ baryon, as well as the value of the branching ratio, can discriminate the scale dimension $d_\mathcal{U}$, especially at $d_\mathcal{U} < 1.7$ (1.85) when $\Lambda_\mathcal{U} = 2$ TeV (1 TeV) for the scalar, and at $d_\mathcal{U} < 2$ (2.04) when $\Lambda_\mathcal{U} = 2$ TeV (1 TeV) for the vector operators under the assumption that $C_S = C_P$, and $C_V = C_A$, respectively.

Next we study the polarization effects due to the polarizations of Λ and Λ_b, which can be very useful for establishing unparticle physics, especially at lower values of the scale dimension $d_\mathcal{U}$, since the branching ratios at larger values of $d_\mathcal{U}$ are extremely smaller compared to that of the SM prediction, under the above-mentioned assumption about the coupling constants. Therefore in this region, study of the polarization effects seems to be the unique approach in establishing unparticle physics. We see that, the longitudinal polarization of Λ_b baryon, and longitudinal and transversal polarizations of Λ baryon, are practically independent on

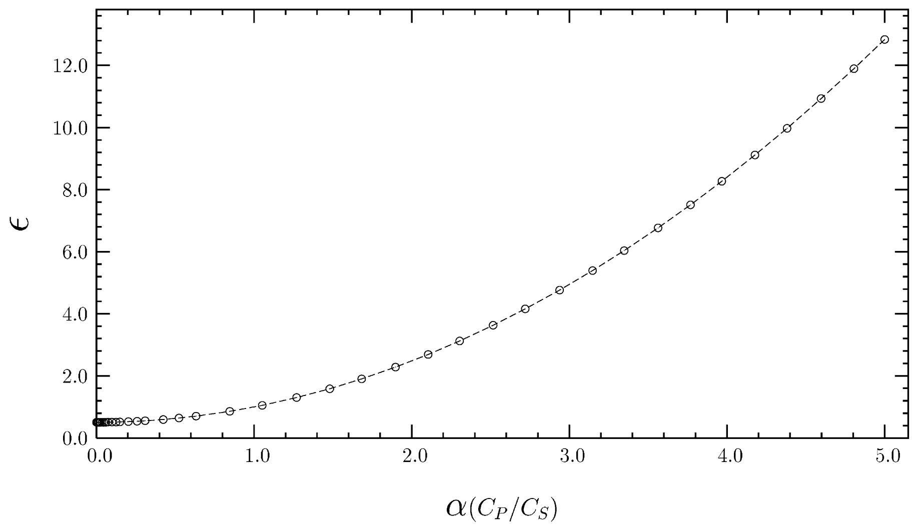

Fig. 13. The dependence of $\epsilon = \mathcal{B}(\Lambda_b \to \Lambda + U)_{\alpha\neq 1}/\mathcal{B}(\Lambda_b \to \Lambda + U)_{\alpha=1}$ on the parameter $\alpha = C_P/C_S$ for the scalar unparticle.

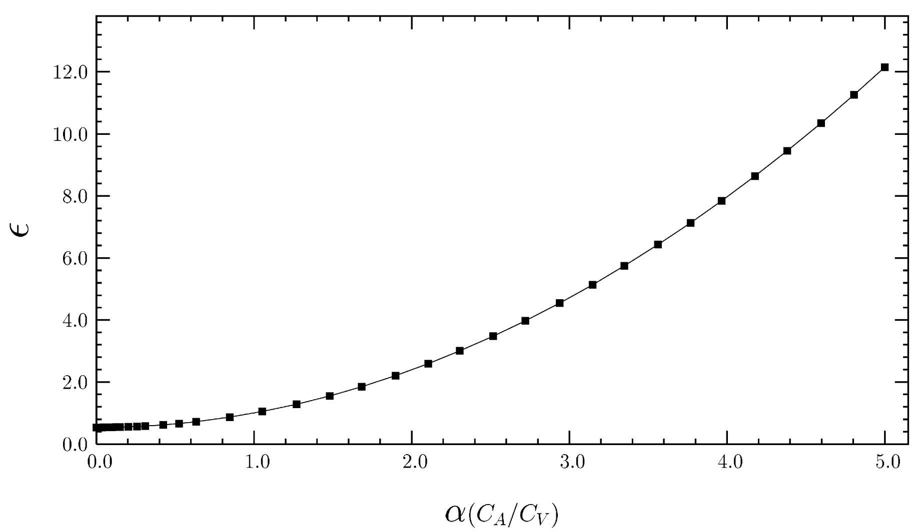

Fig. 14. The same as in Fig. 13, but on the parameter $\alpha = C_A/C_V$ for the vector unparticle.

the value of the scale dimension $d_{\mathcal{U}}$, as well as, on the couplings C_P, C_S of the scalar, and on the couplings C_V, C_A of the vector operators with fermions, if $C_S = C_P$ and $C_V = C_A$.

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