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Charmed penguin versus BAU

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Since the Standard Model most probably cannot explain the large value of CP asymmetries recently observed in D-meson decays we propose the fourth quark-lepton generation explanation of it. As a byproduct weakly mixed leptons of the fourth generation make it possible to save the baryon number of the Universe from erasure by sphalerons. An impact of the 4th generation on BBN is briefly discussed.

1. Introduction. Recently LHCb collaboration has measured the unexpectedly large CP violating asymmetries in $D \to \pi^+\pi^-$ and $D \to K^+K^-$ decays [1]:

$$\Delta A_{\rm CP}^{\rm LHCb} \equiv A_{\rm CP}(K^+K^-) - A_{\rm CP}(\pi^+\pi^-) = \\ = [-0.82 \pm 0.21(\text{stat.}) \pm 0.11(\text{syst.})]\%, \qquad (1)$$

where

$$A_{\rm CP}(\pi^+\pi^-) = \frac{\Gamma(D^0 \to \pi^+\pi^-) - \Gamma(\bar{D}^0 \to \pi^+\pi^-)}{\Gamma(D^0 \to \pi^+\pi^-) + \Gamma(\bar{D}^0 \to \pi^+\pi^-)} \qquad (2)$$

and $A_{\rm CP}(K^+K^-)$ is defined analogously.

This result was later confirmed by CDF collaboration, which obtained [2]:

$$\Delta A_{\rm CP}^{\rm CDF} = [-0.62 \pm 0.21(\text{stat.}) \pm 0.10(\text{syst.})]\%. \qquad (3)$$

The most important question concerning experimental results (1) and (3) is whether in the Standard Model the CP-violation (CPV) in these decays can be as large as 0.5–1%.

In the Standard Model the CPV in $D(\bar{D}) \to \pi^+\pi^-$ decays originates from the interference of the tree and penguin diagrams shown in Fig. 1. For $D(\bar{D}) \to K^+K^-$ decays d-quarks in these diagrams should be substituted by s-quarks.

It is convenient to present the penguin diagram contribution to $D \to \pi^+\pi^-$ decay amplitude in the following form [3]:

$$V_{cd}V_{ud}^* f(m_d) + V_{cs}V_{us}^* f(m_s) + V_{cb}V_{ub}^* f(m_b) = \\ = V_{cd}V_{ud}^*[f(m_d) - f(m_s)] + V_{cb}V_{ub}^*[f(m_b) - f(m_s)], \quad (4)$$

attributing the first term to the tree amplitude and considering the second term only as the penguin amplitude.

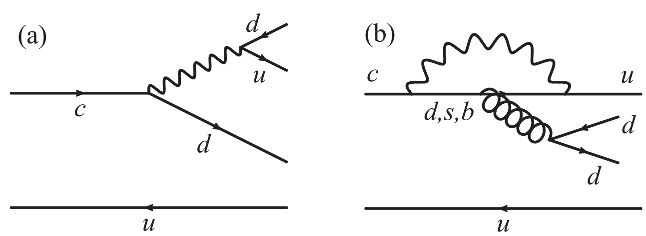

Fig. 1. Quark diagrams describing $D \longrightarrow \pi^+\pi^-$ decay in the Standard Model. A wavy line denotes W-boson, a curly line – gluon

In the case of $D \to K^+K^-$ decay the following presentation is useful [3]:

$$V_{cd}V_{ud}^* f(m_d) + V_{cs}V_{us}^* f(m_s) + V_{cb}V_{ub}^* f(m_b) = \\ = V_{cs}V_{us}^*[f(m_s) - f(m_d)] + V_{cb}V_{ub}^*[f(m_b) - f(m_d)], \ (5)$$

where the first term is attributed to the tree amplitude while the second one is the penguin amplitude.

Denoting the absolute values of $D \to \pi^+\pi^-$ decay amplitudes by T and P we get:

$$A_{\pi^+\pi^-} = T\left[1 + \frac{P}{T}e^{i(\delta-\gamma)}\right], \\ \bar{A}_{\pi^+\pi^-} = T\left[1 + \frac{P}{T}e^{i(\delta+\gamma)}\right], \qquad (6)$$

where δ stands for the difference of the strong interaction phases of the tree and the penguin amplitudes, while $\gamma \approx 70^0$ is the phase of V_{ub} (the product $V_{cd}V_{ud}^*$ as well as V_{cb} are practically real in the standard parametrization of the CKM matrix).

From Eq. (6) for the CPV asymmetry we obtain:

$$A_{\rm CP}(\pi^+\pi^-) = 2\frac{P}{T}\sin\delta\sin\gamma, \tag{7}$$

where in the denominator of (2) we neglect the terms of the order of P/T and $(P/T)^2$ which is a very good approximation because $P/T \sim |V_{cb}V_{ub}^*|/V_{cd} \ll 1$. Here $\sin\gamma$ is close to unity and we use this value in what follows.

Let us present an argument demonstrating that δ can also be close to 90^0. The tree diagram gives dominant contribution to the $D \to \pi\pi$ decay rates. The corresponding to it 4-fermion Hamiltonian has parts with isospin 1/2 and 3/2. That is why the produced π-meson may have isospin zero or two. So three decay probabilities, $D^+ \to \pi^+\pi^0$, $D^0 \to \pi^+\pi^-$, and $D^0 \to \pi^0\pi^0$, depend on the absolute values of the decay amplitudes A_0 and A_2 and their strong phases difference $\delta_0 - \delta_2$. From the experimentally measured branching ratios [4]:

$$\mathrm{Br}(D^+ \to \pi^+\pi^0) = [12.6 \pm 0.9]\cdot 10^{-4},$$
$$\mathrm{Br}(D^0 \to \pi^0\pi^0) = [8.0 \pm 0.8]\cdot 10^{-4},$$
$$\mathrm{Br}(D^0 \to \pi^+\pi^-) = [13.97 \pm 0.26]\cdot 10^{-4} \tag{8}$$

we find for the phase difference of the amplitudes with $I = 0$ and $I = 2$:

$$|\delta_0 - \delta_2| = 86^0 \pm 4^0. \tag{9}$$

In Eq. (7) δ stands for the difference of the strong phases of penguin amplitude which has $I = 1/2$ and produces pions with $I = 0$ and tree amplitude, which has parts with $I = 1/2$ and $I = 3/2$ and produces pions with $I = 0$ and $I = 2$, that is why $\delta \neq \delta_0 - \delta_2$. Nevertheless Eq. (9) demonstrates that δ can be large, and so we substitute $\sin\delta = 1$ into Eq. (7).

In the limit of U-spin ($d \leftrightarrow s$ interchange) symmetry the tree amplitude of $D(\bar D) \to K^+K^-$ decay differs by sign from that of $D(\bar D) \to \pi^+\pi^-$ decay, while the penguin amplitudes of these decays are equal, that is why

$$A_{\rm CP}(K^+K^-) = -A_{\rm CP}(\pi^+\pi^-). \tag{10}$$

However since [4]

$$\mathrm{Br}(D^0 \to K^+K^-) = [39.4 \pm 0.7]\cdot 10^{-4}, \tag{11}$$

we obtain from Eq. (8) that $|A_{K^+K^-}/A_{\pi^+\pi^-}| \simeq 1.7$ and U-spin symmetry is heavily broken in D decays. Nevertheless let us suppose that (10) is not badly violated, so finally we get:

$$\Delta A_{\rm CP} = 4\,P/T. \tag{12}$$

Now let us try to understand if in the Standard Model we can obtain

$$P/T = 1.8\cdot 10^{-3}, \tag{13}$$

which is needed to reproduce the average value of the LHCb and CDF results.

2. $D \to \pi\pi$: charmed penguin. Though the four-fermion quark Hamiltonian responsible for these decays is known, strong interactions does not allow to make an exact calculation of the decay amplitudes. What can be done is an estimate of the decay amplitudes with the help of factorization. Let us start from the tree diagram shown in Fig. 1a which dominates in the decay amplitude:

$$T = \frac{G_{\rm F}}{\sqrt{2}} V_{cd}\langle\pi^+\pi^-|\bar d\gamma_\alpha(1+\gamma_5)c\bar u\gamma_\alpha(1+\gamma_5)d|D^0\rangle \times$$
$$\times\left\{\frac{2}{3}[\alpha_s(m_c)/\alpha_s(M_W)]^{-2/b} + \right.$$
$$\left. + \frac{1}{3}[\alpha_s(m_c)/\alpha_s(M_W)]^{4/b}\right\}, \tag{14}$$

where the last factor originates from the summation of the gluon exchanges in the leading logarithmic approximation. Substituting into it $b = 11 - 2/3N_f = 23/3$, $\alpha_s(M_W) = 0.12$, $\alpha_s(m_c) = 0.3$ we find that the factor in the curly brackets is close to one, $\{...\} = 1.1$. Factorizing the decay amplitude we obtain:

$$T = 1.1\frac{G_{\rm F}}{\sqrt{2}} V_{cd}\langle\pi^+|\bar u\gamma_\alpha(1+\gamma_5)d|0\rangle \times$$
$$\times \langle\pi^-|\bar d\gamma_\alpha(1+\gamma_5)c|D^0\rangle =$$
$$= 1.1\frac{G_{\rm F}}{\sqrt{2}} V_{cd} f_\pi k_{1\alpha} \times$$
$$\times [f_+^\pi(0)(p+k_2)_\alpha + f_-^\pi(0)(p-k_2)_\alpha] =$$
$$= 1.1\frac{G_{\rm F}}{\sqrt{2}} V_{cd} f_\pi f_+^\pi(0) m_D^2, \tag{15}$$

where k_1 and k_2 are the momenta of the produced π-mesons, p is the D-meson momentum and we neglect m_π^2 in comparison with m_D^2.

The value of the $D^0 \to \pi^+e^+\nu$ transition formfactor at $q^2 = 0$ can be found in Ref. [4]:

$$f_+^\pi(0)|V_{cd}| = 0.152 \pm 0.005, \quad f_+^\pi(0) = 0.66, \tag{16}$$

and for the decay width we obtain:

$$\Gamma^{\rm theor}_{D\to\pi^+\pi^-} = \frac{G_{\rm F}^2}{2}\frac{[1.1 V_{cd} f_+^\pi(0) f_\pi m_D^2]^2}{16\pi m_D} = 6.2\cdot 10^9\ {\rm s}^{-1}, \tag{17}$$

where $f_\pi = 130$ MeV was used.

From the branching ratio of the $D^0 \to \pi^+\pi^-$ decay (8) and D^0-meson mean life, $\tau_{D^0} = 0.41 \cdot 10^{-12}$ s, we find:

$$\Gamma^{\text{exp}}_{D\to\pi^+\pi^-} = 3.4 \cdot 10^9\,\text{s}^{-1}. \tag{18}$$

So the naive factorization overestimates the decay amplitude by the factor $\sqrt{6.2/3.4} \approx 1.4$.

Calculating the $D \to K^+K^-$ decay probability we obtain:

$$\Gamma^{\text{theor}}_{D\to K^+K^-} = \left[\frac{f_K}{f_\pi}\frac{f_+^K(0)}{f_+^\pi(0)}\right]^2 \Gamma^{\text{theor}}_{D\to\pi^+\pi^-} = 12.2 \cdot 10^9\,\text{s}^{-1}, \tag{19}$$

where we substituted $f_K/f_\pi = 1.27$ and $f_+^K(0) = 0.73$ taken from Ref. [4].

From Eq. (11) it follows:

$$\Gamma^{\text{exp}}_{D\to K^+K^-} = 9.6 \cdot 10^9\,\text{s}^{-1}, \tag{20}$$

so the factorization overestimates the decay amplitude by the factor $\sqrt{12.2/9.6} = 1.1$.

We see that in the case of the tree diagrams the accuracy of the factorization approximation is very good.

Let us make a brief remark on the $D \to K^0\bar{K}^0$ decay. At the tree level it proceeds through the diagram with W-boson exchange in t-channel, so it should be suppressed. Even more, $c\bar{u} \to d\bar{d}$ and $c\bar{u} \to s\bar{s}$ amplitudes interfere destructively and in the U-spin symmetry limit their sum is zero [5]. According to experimental data [4]:

$$\text{Br}(D^0{\to}K^0\bar{K}^0) = 4\text{Br}(D^0{\to}2K_S^0) = (6.8 \pm 1.2) \cdot 10^{-4}, \tag{21}$$

which is approximately 6 times smaller than $\text{Br}(D \to K^+K^-)$. It means that the decay amplitude is smaller than that to charged kaons by factor 2.5. This unexpectedly small suppression may indicate that large distance effects like $D^0 \to K^{*+}K^{*-} \to K^0\bar{K}^0$ rescattering can be important.

The four-fermion QCD penguin amplitude which describes $D \to \pi^+\pi^-$ decay looks like:

$$\begin{aligned} H(P) &= \frac{G_{\text{F}}}{\sqrt{2}} V_{cb}V_{ub}^* \frac{\alpha_s(m_c)}{12\pi} \times \\ &\times \ln\left(\frac{m_b}{m_c}\right)^2 [\bar{u}\gamma_\alpha(1+\gamma_5)\boldsymbol{\lambda}c)](\bar{d}\gamma_\alpha\boldsymbol{\lambda}d) = \\ &= \frac{G_{\text{F}}}{\sqrt{2}} V_{cb}V_{ub}^* \frac{\alpha_s(m_c)}{12\pi} \times \\ &\times \ln\left(\frac{m_b}{m_c}\right)^2 \{[\bar{u}\gamma_\alpha(1+\gamma_5)d][\bar{d}\gamma_\alpha(1+\gamma_5)c] - \\ &- 2\bar{u}(1-\gamma_5)d\bar{d}(1+\gamma_5)c\} \cdot \frac{8}{9}, \end{aligned} \tag{22}$$

where $\boldsymbol{\lambda}$ are the Gell-Mann SU(3) matrices and we use the Fierz identities:

$$\begin{gathered} \boldsymbol{\lambda}_{ab}\boldsymbol{\lambda}_{cd} = -2/3\delta_{ab}\delta_{cd} + 2\delta_{ad}\delta_{bc}, \\ \bar{\psi}\gamma_\alpha(1{+}\gamma_5)\varphi\bar{\chi}\gamma_\alpha(1{+}\gamma_5)\eta = \bar{\psi}\gamma_\alpha(1{+}\gamma_5)\eta\bar{\chi}\gamma_\alpha(1{+}\gamma_5)\varphi, \\ \bar{\psi}\gamma_\alpha(1{+}\gamma_5)\varphi\bar{\chi}\gamma_\alpha(1{-}\gamma_5)\eta = -2\bar{\psi}(1{-}\gamma_5)\eta\bar{\chi}(1{+}\gamma_5)\varphi. \end{gathered}$$

Also the identity $\langle\pi^+|\bar{u}_a O d_b|0\rangle = 1/3\delta_{ab}\langle\pi^+|\bar{u}Od|0\rangle$, where $O \equiv \gamma_\alpha\gamma_5$ or γ_5, was used.

Calculating the matrix element in the factorization approximation with the help of the equations of motion for quark fields we find:

$$\begin{aligned} P &= \frac{G_{\text{F}}}{\sqrt{2}}|V_{cb}V_{ub}^*|\frac{\alpha_s(m_c)}{12\pi}\ln\left(\frac{m_b}{m_c}\right)^2 \times \\ &\times \frac{8}{9} f_\pi f_+^\pi(0) m_D^2 \left[1 + \frac{2m_\pi^2}{m_c(m_u+m_d)}\right]. \end{aligned} \tag{23}$$

Dividing it by the experimental value of the tree amplitude and using Eq. (15) we obtain:

$$\begin{aligned} P/T &= \frac{1.4}{1.1} \cdot \frac{8}{9}\frac{|V_{cb}V_{ub}^*|}{|V_{cd}|}\frac{\alpha_s(m_c)}{12\pi} \times \\ &\times \ln\left(\frac{m_b}{m_c}\right)^2 \left[1 + \frac{2m_\pi^2}{m_c(m_u+m_d)}\right]. \end{aligned} \tag{24}$$

Substituting $|V_{cd}| = 0.23$, $|V_{ub}| = 3.9 \cdot 10^{-3}$, $V_{cb} = 41{\cdot}10^{-3}$, $\alpha_s(m_c) = 0.3$, $m_b = 4.5$ GeV, $m_c = 1.3$ GeV, $m_u + m_d = 6$ MeV we come to:

$$P/T \approx 9 \cdot 10^{-5}. \tag{25}$$

Comparing it with Eq. (13) we see that in order to fit the experimental data on ΔA_{CP} the penguin amplitude should be enhanced by the factor 20 in comparison with what factorization gives. Concerning the tree amplitudes, we have found in this section that factorization result differs from the experimental value by the factor 1.4 in the case of $D \to \pi^+\pi^-$ decay and by 1.1 in the case of $D \to K^+K^-$ decay. In the next two sections we will study how accurate is the factorization approximation to the penguin amplitudes in B- and K-meson decays.

3. $B \to \pi K$: beautiful penguin. $B_u \to \pi^+K^0$ decay is described by the penguin amplitude shown in Fig. 2.

The Hamiltonian responsible for this decay looks like:

$$\hat{H} = \frac{G_{\text{F}}}{\sqrt{2}} V_{tb}V_{ts}^*(c_3O_3 + c_4O_4 + c_5O_5 + c_6O_6), \tag{26}$$

$V_{tb}V_{ts}^*$ is substituted for $V_{cb}V_{cs}^* + V_{ub}V_{us}^*$ (the contribution of a loop with the virtual t-quark is negligible) and

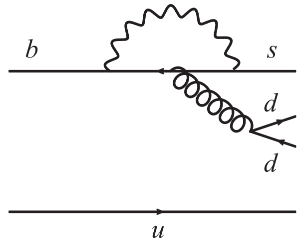

Fig. 2. $B_u \to \pi^+ K^0$ decay proceeds through the penguin amplitude only

$$
\begin{aligned}
O_3 &= \bar{s}\gamma_\alpha(1+\gamma_5)b\bar{d}\gamma_\alpha(1+\gamma_5)d,\\
O_4 &= \bar{s}_a\gamma_\alpha(1+\gamma_5)b_c\bar{d}_c\gamma_\alpha(1+\gamma_5)d_a,\\
O_5 &= \bar{s}\gamma_\alpha(1+\gamma_5)b\bar{d}\gamma_\alpha(1-\gamma_5)d,\\
O_6 &= \bar{s}_a\gamma_\alpha(1+\gamma_5)b_c\bar{d}_c\gamma_\alpha(1-\gamma_5)d_a, \qquad (27)
\end{aligned}
$$

where $a, c = 1, 2, 3$ are the color indexes.

Using the Fierz identities as well as $\langle K^0|\bar{s}_a O d_b|0\rangle = \frac{1}{3}\delta_{ab}\langle K^0|\bar{s}Od|0\rangle$ identity we obtain:

$$
\hat{H} = \frac{G_F}{\sqrt{2}} V_{tb}V_{ts}^*[a_4\bar{s}\gamma_\alpha(1+\gamma_5)d\bar{d}\gamma_\alpha(1+\gamma_5)b - - 2a_6\bar{s}(1-\gamma_5)d\bar{d}(1+\gamma_5)b], \qquad (28)
$$

where $a_4 = \frac{1}{3}c_3 + c_4$, $a_6 = \frac{1}{3}c_5 + c_6$. Calculating the matrix element in the factorization approximation we obtain:

$$
M = \frac{G_F}{\sqrt{2}} V_{tb}V_{ts}^* f_K f_+(0) m_B^2 \left(a_4 + a_6\frac{2m_K^2}{m_b m_s}\right), \qquad (29)
$$

where in the leading logarithmic approximation the following approximate equation is valid:

$$
a_4 = a_6 = -\frac{\alpha_s(m_b)}{12\pi}\ln\left(\frac{M_W}{m_b}\right)^2 \approx -0.03, \qquad (30)
$$

while at NLO approximation from Table 1 of [6] we obtain: $a_4 = -0.031$, $a_6 = -0.042$. Substituting $m_s = 100\,\mathrm{MeV}$, $m_b = 4.5\,\mathrm{GeV}$ we find:

$$
\Gamma^{\mathrm{fact}}_{(B_u\to\pi^+K^0)} = \frac{G_F^2|V_{ts}|^2}{32\pi} f_K^2 f_+^2(0) m_B^3 (0.076)^2 = = 4.1\cdot 10^6\,\mathrm{s}^{-1}, \qquad (31)
$$

where $V_{ts} = 39\cdot 10^{-3}$ and $f_+(0) = 0.25$ from [4] was used. The experimental result is:

$$
\Gamma^{\mathrm{exp}}_{(B_u\to\pi^+K^0)} = 14\cdot 10^6\,\mathrm{s}^{-1}. \qquad (32)
$$

So, the factorization result is enhanced by the factor

$$
P/P_{\mathrm{fact}} = \sqrt{14/4.1} = 1.8. \qquad (33)
$$

The numerical value of the penguin amplitude is important in the calculation of CP asymmetries in $B \to \pi K$ and $B \to \pi\pi$ decays [7].

4. $K \to \pi\pi$: strange penguin. $s \to d$ penguin transition changes the isospin by 1/2 in this way explaining the famous $\Delta I = 1/2$ rule in $K \to \pi\pi$ decays. The calculation of the $K_S \to \pi^+\pi^-$ decay amplitude generated by the penguin transition using the factorization underestimates the amplitude by the factor 2–3 according to Refs. [6, 8].

In view of the results for B and K decays we can cautiously assume that for $D \to \pi^+\pi^-$ decay the factorization calculation underestimates the penguin amplitude at most by factor 5 leading to:

$$
\left(\Delta A_{\mathrm{CP}}^{\mathrm{theor}}\right)_{\mathrm{SM}} \lesssim 0.2\%. \qquad (34)
$$

Thus the following alternative emerges: either the experimental results are wrong or New Physics is found. Of course we cannot determine what kind of new particles and interactions are responsible for large CPV asymmetry in $D \to \pi^+\pi^-$ (K^+K^-) decays. However, in the next section we will propose the straightforward generalization of the Standard Model in which large CPV in D decays can be explained.

5. The fourth generation: enhancement of CPV in D decays. As it was stated in paper [9] the introduction of the fourth quark-lepton generation may easily remove Standard Model upper bound (34) matching the experimental results [1, 2]. In the case of the fourth generation the additional term with the intermediate b' quark should be added to the expression for the penguin amplitude. In this way expression (4) is substituted by:

$$
\begin{aligned}
&V_{cd}V_{ud}^* f(m_d) + V_{cs}V_{us}^* f(m_s) + V_{cb}V_{ub}^* f(m_b) + \\
&+ V_{cb'}V_{ub'} f(m_{b'}) = V_{cd}V_{ud}^*[f(m_d) - f(m_s)] + \\
&+ V_{cb}V_{ub}^*[f(m_b) - f(m_s)] + \\
&+ V_{cb'}V_{ub'}[f(m_{b'}) - f(m_s)], \qquad (35)
\end{aligned}
$$

where the unitarity of 4×4 quark mixing matrix is used. According to the experimental constraints from the direct searches of the fourth generation quarks b' should weigh several hundreds GeV, that is why $f(m_{b'})$ is small and can be neglected just as it is done with t-quark contribution to $b \to s$ penguin, see the remark after Eq. (26). In order to enhance SM contribution to the penguin amplitude we should suppose that the term $V_{cb'}V_{ub'}f(m_s)$ dominates.

Then the enhancement of A_{CP} in the case of the fourth generation is equal to:

$$
\frac{P_4}{P_{\mathrm{SM}}} = \frac{\ln(m_W/m_c)}{\ln(m_b/m_c)}\frac{|V_{cb'}V_{ub'}^*|}{|V_{cb}V_{ub}|}\frac{\sin(\arg V_{cb'}V_{ub'}^*)}{\sin\gamma} \approx \approx 3.3\frac{3\cdot 10^{-4}}{1.5\cdot 10^{-4}} \approx 6, \qquad (36)
$$

where in the last equality we use the allowed values of the product $|V_{cb'}V^*_{ub'}|\sin(\arg V_{cb'}V_{ub'})$ taken from Fig. 1 of paper [10] [1)]. So we see that the enhancement necessary to describe the experimental data on ΔA_{CP} can be achieved in the case of the fourth generation.

6. Saving baryon number by long-lived fourth generation neutrino. If weakly mixed particles exist, then the sphaleron processes can create the baryon asymmetry of the universe [11]. As it is noted in Ref. [12], the long-lived fourth generation particles save baryon asymmetry generated in the early universe from erasure by the sphaleron transitions. The sphaleron transitions conserve $B-L$, thus, if in the early universe $B_0 = L_0 \neq 0$ is generated, then the final baryon and lepton asymmetries being proportional to $B-L$ are completely erased. If the fourth generation particles weakly mix with three quark-lepton generations of the Standard Model, then two additional quantities are conserved: B_4-L_4 and $L-3L_4$, where B_4 and L_4 are the densities of baryons and leptons of the fourth generation, while B and L are the densities of baryons and leptons of three light generations. In Ref. [12] initial asymmetries $B_0 = L_0 = 3\Delta$ and $B_4^0 = L_4^0 = 0$ were chosen and since $L - 3L_4 = 3\Delta \neq 0$, the total baryonic number density, $B + B_4$, being proportional to a linear superposition of conserved quantities is nonzero at the sphaleron freeze-out temperature. After the sphaleron freeze-out $B + B_4$ is conserved in comoving volume and is equal to the present day baryon density of the Universe. However, if heavy baryons of the 4th generation do not decay prior to big bang nucleosynthesis (BBN), the light baryon number density at BBN could be different from that determined from the angular fluctuations of CMB. The impact of this effect on the light element abundances is discussed below.

For such a scenario to occur the lifetimes of the fourth generation quarks and leptons should be at least larger than the Universe age at the sphaleron freeze-out: $\tau_4 > M_{\rm Pl}/T^2_{\rm sph} \sim 10^{-10}$ s. For the mixing angles in the case of $b' \to (c,u)W$ decay it gives $\theta < 10^{-8}$ [12], much smaller than what we need to explain the large CPV in D-decays, see Eq. (36).

So in our case quarks of the fourth generation should be much stronger mixed with quarks of three light generations. However, let us suppose that leptons of the fourth generation are weakly mixed with the leptons of three light generations. Let us introduce the total baryon density, $B' \equiv B + B_4$, and take the initial conditions analogous to those in Ref. [12]: $B'_0 = L_0 = 3\Delta$ and $L_4^0 = 0$. We can choose four independent chemical potentials as: μ_{u_L}, μ_W, μ_{N_L} and $\mu \equiv \mu_{\nu_e}+\mu_{\nu_\mu}+\mu_{\nu_\tau}$, which are the chemical potentials for the upper type quarks, W-bosons, $4G$ neutrino and sum over all SM neutrino chemical potentials (see Appendix). In the limit $\mu_i/T \ll 1$ the baryon and lepton densities are linear combinations of these chemical potentials with the coefficients which depend on the ratio of masses of the corresponding particles to the temperature. We will take into account the masses of W-boson, t-quark, t'- and b'-quarks of the fourth generation and the fourth generation leptons N and E, the masses of all the other components of the primeval plasma can be neglected in comparison with $T_{\rm sph}$.

Finally we have four equations for four unknown chemical potentials: two quantities are conserved under the sphaleron transitions; we can choose them as

$$\begin{aligned} B' - L - L_4 &= 0, \\ L - 3L_4 &= 3\Delta. \end{aligned} \tag{37}$$

The third equation is that of the electric neutrality of the primeval plasma, $Q = 0$, and, finally, the sum of the chemical potentials of all the particles which are converted into nothing by sphaleron ($qqql$ of each generation) equals zero. The values of masses of the 4th generation particles we take from paper [13] in which the fit to the electroweak observables for higgs mass $m_H = 125$ GeV was performed and recent LHC bounds on the masses of t'- and b'-quarks were taken into account:

$$\begin{aligned} m_{t'} &= 634 \text{ GeV}, \quad m_{b'} = 600 \text{ GeV}, \\ m_E &= 107.6 \text{ GeV}, \quad m_N = 57.8 \text{ GeV}. \end{aligned} \tag{38}$$

The dashed blue line in Fig. 3 corresponds to the case of the unmixed fourth generation particles considered in [12]. The results for the case of the strongly mixed fourth generation quarks and the unmixed fourth generation leptons are shown by the solid green line. In order that leptons, N, do not decay before the sphaleron freeze-out, which happens at $t_U \sim 10^{-10}$ s, the mixing angles of N with three light neutrinos should be small: $\theta_i < 10^{-5}$ (N decays through four fermion interaction). Assuming similar bound $\theta < 10^{-6}$, the existence of heavy Dirac sequential neutrino with $m_N = (50-100)$ GeV is compatible with the search at LEP II [14].

According to the standard cosmological scenario nonrelativistic matter started to dominate the cosmic energy density at redshift $z \approx 10^4$. If we demand that N should decay before that epoch, its life-time should be sufficiently short, $\tau_N < 10^{13}$ s, from which we obtain

[1)]Let us stress that the logarithmic ($\log(m_W/m_c)$) enhancement originates not from the diagram with the intermediate b' quark but from the term $f(m_s)$.

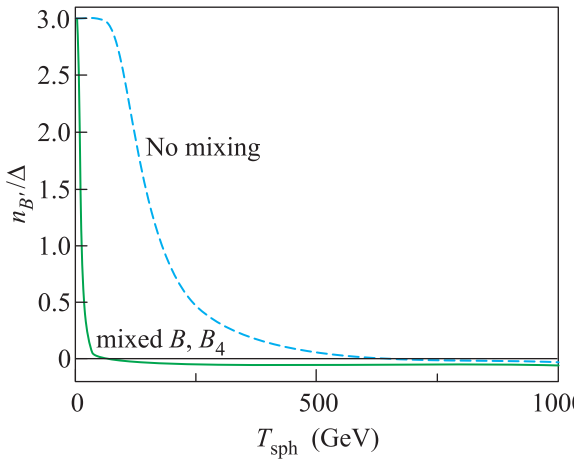

Fig. 3. (Color online) The final baryon asymmetry versus the initial asymmetry $n_{B'}/\Delta$ as a function of sphaleron freeze-out temperature $T_{\rm sph}$ for the unmixed fourth generation is shown by a dashed (blue) line. It is analogous to Fig. 2 from [12] but for $m_N = 57.8$ GeV, $m_E = 107.6$ GeV, $m_{t'} = 634$ GeV, $m_{b'} = 600$ GeV. The final baryon asymmetry for the case of the mixed fourth generation quarks and the unmixed fourth generation leptons is shown by a solid (green) line

the lower bound $\theta > 10^{-16}$. (Let us note that direct searches exclude N as a unique dark matter candidate [15].)

A stronger bound on τ_N follows from the equilibrium form of the energy spectrum of CMB. According to Ref. [16] a large influx of energy into the usual cosmological cosmic background would be thermalized if it took place before $z \sim 10^7$. Otherwise the observed black body spectrum of CMB would be noticeably distorted. Since the precision of the spectral shape is at the level of 10^{-4}, only a very small distortion is permitted.

The condition that N decays before or at $z \sim 10^7$ demands $\tau_N < 10^6$ s, or $\theta > 10^{-13}$. If N indeed decays before $z \sim 10^7$, the contribution from its decay to the energy density of CMB would be not larger than 1% and the ratio of baryon to photon number densities $\eta_B \equiv n_B/n_\gamma$ at BBN epoch and at CMB recombination would be slightly different but in principle measurable by the light element abundances.

More interesting and pronounced effect appears if heavy quarks of the 4th generation are long-lived. In this case we cannot explain the large value of CPV in D decays but may explain the difference of $\eta_{\rm B}$ at BBN epoch ($\eta_{\rm BBN}$) and at the recombination ($\eta_{\rm rec}$) which is probably requested by the recent data on the light element abundances [17]. If heavy baryons of the 4th generation decays after BBN but before the hydrogen recombination, the number of light baryons in the comoving volume at BBN would be different from that at the recombination. The ratio $\eta_{\rm BBN}/\eta_{\rm rec}$ at these epochs could be either larger or smaller than unity depending upon the value of the baryon asymmetry in the heavy quark sector and the energy influx to CMB from the heavy baryon decays. So in principle both rise or decrease of η_{BBN} is possible[2)].

In the limit $T_{\rm sph} \to 0$ heavy particles of the fourth generation are not produced: $B_4 = L_4 = 0$, $B' = L = 3\Delta$. In the physically interesting opposite limit $T_{\rm sph} \gg m_N$ the value of baryon asymmetry is nonzero since the right-handed neutrinos of three light generations are not produced in the primordial plasma violating symmetry between the leptons of four generations which would occur at $T \gg m_N$. The characteristic time of the right-handed neutrino to thermalize is T/m_ν^2 and for $m_\nu \lesssim 1$ keV (which is valid for three light neutrinos) this time is longer than the Universe age, $t_U = M_{\rm Pl}/T^2$ for $T = T_{\rm sph} \approx 200$ GeV [11].

7. Conclusions. In Introduction we determined what ratio of the penguin to the tree amplitudes of $D \to \pi^+\pi^-$ decay is needed to get the observed CP asymmetry. In Section 2 we found that the factorization describes the tree amplitude with good accuracy; concerning the penguin amplitude it appears to be twenty times smaller than one needs to describe the experimental data on $A_{\rm CP}$. In Section 3 we demonstrated that in the case of $B \to \pi^+ K^0$ decay the factorization underestimates the penguin amplitude by factor 2. In the case of $K_S \to \pi^+\pi^-$ decay the penguin amplitude is enhanced by factor 2–3 in comparison with the factorization result.

Thus if confirmed on larger statistics and future systematics result (1) demands New Physics.

In Section 5 we demonstrated that the fourth quark-lepton generation may enhance the penguin amplitude describing the experimental data. If the leptons of the fourth generation weakly mix with three light generation leptons, then the baryonic charge generated at high scale escapes the erasure by sphalerons and survives till now according to the results presented in Section 6.

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2) Since both the value of $\eta_{\rm BBN}$ and the number of light neutrino species influence nucleosynthesis, the change in the value of $\eta_{\rm BBN}$ can be formulated as an additional (positive or negative) number of light neutrino species [17].

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Appendix. Below we derive equations used in Section 6 to find the dependence of the baryon asymmetry of the Universe on the sphaleron freeze-out temperature. In this Appendix we closely follow paper [12].

Being interested in the values of the asymmetries at sphaleron freeze-out temperature we should assume that the electroweak phase transition already has occured and the neutral Higgs boson condenses. That is why the Higgs boson chemical potential is zero. Sometimes in the literature the baryon density in the electroweak unbroken phase is looked for. In this case the Higgs boson does not condense and its chemical potential is nonzero. To find it an additional equation is needed. It is provided by the condition that the density of charges with which the massless bosons interact should be zero, and in an unbroken phase there are two such charges: the hypercharge and the third projection of a weak isospin. The baryon density in the unbroken phase is analyzed, for example, in book [18] and it differs from its value in a broken phase. Since the right-handed components of quarks and leptons emitting neutral Higgs transform to the left-handed components the chemical potential of both components are equal: $\mu_{u_R} = \mu_{u_L} \equiv \mu_u$, $\mu_{d_R} = \mu_{d_L} \equiv \mu_d$, $\mu_{e_R} = \mu_{e_L} \equiv \mu_e$. The analogous relations are valid for the particles of the second and third families. The right-handed neutrinos of three light generations are not thermalized and should not be taken into account (see the end of Sect. 6). The fourth generation right-handed neutrinos, being heavy, rapidly thermalize: $\mu_{N_R} = \mu_{N_L} \equiv \mu_N$. The chemical potentials of up and down weak isospin components are related by W^- chemical potential: $\mu_d = \mu_W + \mu_u$, $\mu_e = \mu_W + \mu_\nu$, $\mu_E = \mu_W + \mu_N$. Mixing of quarks of four families and leptons of three families equilibrates the chemical potentials of the particles with the identical gauge quantum numbers. As a result four independent chemical potentials remain: μ_u, μ_N, μ_W, and $\mu \equiv \mu_{\nu_1} + \mu_{\nu_2} + \mu_{\nu_3} \equiv 3\mu_\nu$.

The particle number densities depend on their (Fermi or Bose) statistics, temperature, chemical potential, and masses. The chemical potential of an antiparticle is opposite to that of the particle. The asymmetries and, hence, chemical potentials are very small. Expanding the equilibrium integrals for the asymmetry over μ we obtain:

$$n_p = \frac{g_p}{\pi^2} T^3 \frac{\mu}{T} \int\limits_x^\infty y\sqrt{y^2 - x^2}\frac{e^y}{(1 \pm e^y)^2} dy =$$

$$= \begin{cases} \dfrac{g_p T^3}{3}\left(\dfrac{\mu}{T}\right)\alpha_b(x), & \text{if } p \text{ is a boson,} \\ \dfrac{g_p T^3}{6}\left(\dfrac{\mu}{T}\right)\alpha_f(x), & \text{if } p \text{ is a fermion,} \end{cases} \tag{A.1}$$

where g_p is the number of the degrees of freedom of the particle p ($g_q = g_l = 2$, $g_\nu = 1$, $g_N = 2$, $g_W = 3$) and $x = m/T$. Functions $\alpha(x)$ are normalized in such a way that $\alpha_b(0) = \alpha_f(0) = 1$. In what follows we take into account the nonzero masses of the particles of the fourth generation, of t-quark, and of W-boson.

The condition of electroneutrality of the primeval plasma looks as:

$$Q = 3 \cdot \frac{2}{3}[2(\alpha_u + \alpha_c + \alpha_t + \alpha_{t'})\mu_u] -$$
$$- 3 \cdot \frac{1}{3}[2(\alpha_d + \alpha_s + \alpha_b + \alpha_{b'})(\mu_W + \mu_u)] -$$
$$- 2[(\alpha_e + \alpha_\mu + \alpha_\tau)(\mu_W + \mu_\nu)] - 2\alpha_E(\mu_W + \mu_N) -$$
$$- 3 \cdot 2\alpha_W \mu_W = 0, \tag{A.2}$$

$$(1 + 2\alpha_t + 2\alpha_{t'} - \alpha_{b'})\mu_u -$$
$$-(6 + \alpha_{b'} + \alpha_E + 3\alpha_W)\mu_W - \mu - \alpha_E \mu_N = 0. \tag{A.3}$$

Here and below we omit irrelevant factor $T^2/6$.

The sphaleron transition converts $qqql$ combination of each generation into vacuum, which gives:

$$12\mu_u + 8\mu_W + \mu + \mu_N = 0. \tag{A.4}$$

The remaining two equations are two superpositions of B', L, and L_4 conserved under sphaleron transitions thus being equal to their initial values. The expressions for these quantities look like:

$$L_4 = 2\alpha_E\mu_E + 2\alpha_N\mu_N = 2(\alpha_E + \alpha_N)\mu_N + 2\alpha_E\mu_W, \tag{A.5}$$

$$L = 2(\alpha_e + \alpha_\mu + \alpha_\tau)\mu_e + (\alpha_{\nu_e} + \alpha_{\nu_\mu} + \alpha_{\nu_\tau})\frac{\mu}{3} = 3\mu + 6\mu_W, \tag{A.6}$$

$$B' = 2 \cdot 3 \cdot \frac{1}{3}[(\alpha_u + \alpha_c + \alpha_t + \alpha_{t'})\mu_u +$$
$$+ (\alpha_d + \alpha_s + \alpha_b + \alpha_{b'})\mu_d] =$$
$$= 2(2 + \alpha_t + \alpha_{t'})\mu_u + 2(3 + \alpha_{b'})(\mu_u + \mu_W). \tag{A.7}$$

Thus we have four equations which determine the chemical potentials: (A.3), (A.4), and the remaining two:

$$B' - L - L_4 = 2(5 + \alpha_t + \alpha_{t'} + \alpha_{b'})\mu_u + 2(\alpha_{b'} - \alpha_E)\mu_W -$$
$$-3\mu - 2(\alpha_E + \alpha_N)\mu_N = 0, \tag{A.8}$$

$$L - 3L_4 = 6(1 - \alpha_E)\mu_W + 3\mu - 6(\alpha_E + \alpha_N)\mu_N = 3\Delta, \tag{A.9}$$

where we take the initial values analogous to those of Ref. [12]: $B'_0 = L_0 = 3\Delta$ and $L_4^0 = 0$.

When temperature is much larger than the masses of all the particles, all α_i are equal to one we obtain:

$$\frac{B'}{\Delta}\bigg|_{T\gg m_i} = -\frac{11}{179}\,. \qquad \text{(A.10)}$$

If the right-handed neutrinos of three light generations thermalized then the Eq. (A.6) would be substituted by

$$L = 4\mu + 6\mu_W, \qquad \text{(A.11)}$$

and the baryon asymmetry at $T \gg m_i$ would vanish.

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