Constraints on $s\to d\gamma$ from radiative hyperon and kaon decays

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The quark-level process $b\to s\gamma$ has been used extensively to place constraints on new interactions. These same interactions can also be constrained from the enhancement they induce in the quark-level $s\to d\gamma$ transition, to the extent that the short distance contributions can be separated from the long distance contributions. We parametrize what is known about the long distance amplitudes and subtract it from the data in radiative hyperon and kaon decays to constrain new interactions. These constraints complement existing ones from other rare processes although in most cases they are weaker.

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I. INTRODUCTION

The decay mode $b\to s\gamma$ has been used to place constraints on physics beyond the standard model [1]. The mode is particularly useful, constraining new interactions which remove the chirality suppression that occurs in the standard model. In this case the amplitude is enhanced by factors of a heavy mass scale relative to the b quark mass. The same type of new physics enhances the $s\to d\gamma$ transition by a factor of a heavy mass scale relative to the strange quark mass. In models in which the enhancement is as large as one can expect on dimensional grounds, that is $m_b/m_s\sim 30$, it is possible to place interesting constraints on the new physics from $s\to d\gamma$ even though the decay modes involved are dominated by long distance physics [2].

After introducing our notation for the effective interaction responsible for the $s\to d\gamma$ transition, we study the physical radiative hyperon decay amplitudes and the radiative kaon decays of the form $K\to\pi\pi\gamma$. In both cases we expect the amplitudes to be dominated by long distance physics. We describe this long distance physics guided by chiral perturbation theory and subtract it from the physical amplitudes in order to constrain the new, short-distance, interactions.

In the last two sections we illustrate two types of models in which the short distance transition can be significantly enhanced with respect to the standard model.[1] We are not interested here in the specific details of the models, and, for this reason, we only consider the effective low energy operators that the new models may generate. We illustrate the effects of left-right symmetric models [4,5] and of generalized supersymmetric theories [6].

II. SHORT DISTANCE $s\to d\gamma$ IN THE STANDARD MODEL

The low energy effective Hamiltonian responsible for the $s\to d\gamma(g)$ transition can be written as

$$H_{eff}=\sqrt{2}G_F\frac{V_{id}^*V_{is}}{16\pi^2}\bar{d}[(g_s c_{11}^i T^a G_a^{\mu\nu}+ec_{12}^i F^{\mu\nu})\sigma_{\mu\nu}(m_sP_R+m_dP_L)]s+\text{H.c.},\tag{1}$$

where $G_a^{\mu\nu}$ and $F^{\mu\nu}$ are the gluon and photon field strength tensors respectively, and $P_{L,R}\equiv(1\mp\gamma_5)/2$. In the standard model (SM), the coefficients c_{11}^i and c_{12}^i are given at the one loop level without QCD corrections by [9]

$$c_{11}^i=\frac{x_i(2+5x_i-x_i^2)}{4(1-x_i)^3}+\frac{3x_i^2}{2(1-x_i)^4}\ln x_i$$

$$c_{12}^i=\frac{x_i(7-5x_i-8x_i^2)}{12(1-x_i)^3}+\frac{x_i^2(2-3x_i)}{2(1-x_i)^4}\ln x_i,\tag{2}$$

where $x_i=m_i^2/m_W^2$. This contribution to c_{12} from charm and up quarks is negligibly small. Nevertheless, QCD corrections enhance the charm contribution considerably as first discussed in Ref. [7]; we find $c_{12}^c\approx 0.13$.

We assume in this paper that any contribution beyond the SM is due to heavy degrees of freedom which are integrated out at the W mass scale and obtain the coefficients at a hadronic mass scale $\mu\sim 1$ GeV using the expressions [8,9]

$$c_{12}(\mu)=c_{12}^{SM}(\mu)+c_{12}^{new}(\mu),$$

$$c_{12}^{new}(\mu)=\eta^{16/(33-2n_f)}c_{12}^{new}(m_W)+\frac{8}{3}(\eta^{14/(33-2n_f)}-\eta^{16/(33-2n_f)})c_{11}^{new}(m_W),\tag{3}$$

where $\eta=\alpha_s(m_W)/\alpha_s(\mu)$, and n_f is the number of active quarks.

For our phenomenological discussion of physics beyond the standard model we will find it convenient to use instead the effective Lagrangian

[1]These modes also enhance the transitions $d\to d'g$ and were discussed in Ref. [3].

$$\mathcal{L}=\frac{eG_F}{2}c\bar{d}\sigma_{\mu\nu}(1+\gamma_5)sF^{\mu\nu}, \tag{4}$$

in terms of which the standard model reads

$$c_{SM}=\frac{\sqrt{2}m_s}{16\pi^2}\sum_i V_{id}^* V_{is}c_{12}^i. \tag{5}$$

Although we will refer to the coefficient c as if it were unique, new interactions may induce this operator with opposite chirality. The distinction, however, is irrelevant for our purpose. Numerically we will use $\alpha_s(m_Z)=0.119$, the Cabibbo-Kobayashi-Maskawa (CKM) matrix of Ref. [10] in the Wolfenstein parametrization, and $m_s=150$ MeV. This corresponds to $c_{SM}\approx 0.04$.

III. RADIATIVE HYPERON DECAYS

The effective Lagrangian for radiative hyperon decays is usually written in the form

$$\mathcal{L}(B_i\to B_f\gamma)=\frac{eG_F}{2}\bar{B}_f(a+b\gamma_5)\sigma^{\mu\nu}B_iF_{\mu\nu}. \tag{6}$$

Each decay mode is then characterized by the constants a and b which have both real and imaginary (absorptive) parts. The two observables are the decay rate and the asymmetry parameter, which are given in terms of a, b and the photon energy, ω, by

$$\Gamma(B_i\to B_f\gamma(\omega))=\frac{G_F^2e^2}{\pi}(|a|^2+|b|^2)\omega^3$$

$$\frac{d\Gamma}{d\cos\theta}\sim 1+\alpha\cos\theta$$

$$\alpha=\frac{2\mathrm{Re}(ab^\star)}{|a|^2+|b|^2}. \tag{7}$$

The latest numbers found by the Particle Data Group [10] are shown in Table I. Some of these measurements will be improved by the KTEV experiment at Fermilab.

The long distance contributions to these decays have been studied within the context of chiral perturbation theory in Refs. [11,12]. The authors of Ref. [11] find that the imaginary parts of a and b are well known and that the real part of b is also known. They also find that the real part of a cannot be predicted or even estimated reliably. For this reason they treat a as a free parameter and attempt to fit the data. In Table II we summarize these results.

TABLE I. Radiative hyperon decay data. We present the Particle Data Group values [10] for the rate and asymmetry parameters as well as the corresponding value for $\sqrt{(|a|^2+|b|^2)}$.

Mode	$\Gamma\times 10^{15}$ MeV	ω MeV	$\sqrt{(\|a\|^2+\|b\|^2)}$ MeV	α
$\Lambda\to n\gamma$	4.38 ± 0.38	162.22	16.07	-
$\Sigma^+\to p\gamma$	10.13 ± 0.41	224.59	15.00	-0.76 ± 0.08
$\Xi^-\to\Sigma^-\gamma$	0.51 ± 0.09	118.06	8.83	1.0 ± 1.3
$\Xi^0\to\Lambda\gamma$	2.4 ± 0.4	184.13	9.85	0.43 ± 0.44
$\Xi^0\to\Sigma^0\gamma$	7.94 ± 0.91	116.57	35.54	0.20 ± 0.32

TABLE II. Radiative hyperon decay amplitudes as described in the text in units of MeV. An entrance $0^\star$ indicates that a fit is not possible.

Mode	a_I	b_R	b_I	$a_{SD}(SU(6))$	a_R(fit)
$\Lambda\to n\gamma$	-0.68	11.11 ± 1.2	11.21	$\sqrt{3/2}\,c$	2.8 ± 18
$\Sigma^+\to p\gamma$	6.18	-1.21	-0.53	$-1/3\,c$	13.6 ± 9
$\Xi^-\to\Sigma^-\gamma$	-1.55	-7.26	-12.34	$-5/3\,c$	$0^\star$
$\Xi^0\to\Lambda\gamma$	0	-2.47 ± 2.12	0	$-1/\sqrt{6}\,c$	9.5 ± 10
$\Xi^0\to\Sigma^0\gamma$	0	2.52 ± 1.22	0	$5/(3\sqrt{2})\,c$	35 ± 85

To construct Table II we have used the imaginary parts of a and b from Ref. [11], which are reliable, and we have also used their estimate[2] for the real part of b. The column labeled a_R(fit) shows the required value of a_R to reproduce the measured rates when combined with the known values of a_I, b_I and b_R. It is not possible to fit the measured rate for $\Xi^-\to\Sigma^-\gamma$ in this way. In the column labeled $a_{SD}(SU(6))$ we show the short distance contribution from the operator of Eq. (4) using $SU(6)$ wave functions to compute the hadronic matrix elements.

In most cases the new physics contribution to a_R will be identical to the contribution to b_R, so we can bound the coefficient c by requiring that $a_{SD}(SU(6))$ be less than $a_R(\text{fit})/\sqrt{2}$. The best bound is obtained from the mode $\Lambda\to n\gamma$ and it is

$$|c(\mu)|\lesssim 12\ \text{MeV}. \tag{8}$$

We have not used the asymmetry parameters because the only one that is well measured is not understood [11].

IV. RADIATIVE KAON DECAYS

In this section we look at decays of the form $K\to\pi\pi\gamma$. We start with the decay $K_L\to\pi^+\pi^-\gamma$ in which the "direct emission" has been measured (bremsstrahlung is subtracted from the full amplitude). Assuming CP conservation we can write this direct emission amplitude in the form

$$\mathcal{M}=ig_8\frac{2\sqrt{2}eG_F\lambda f_\pi^2}{M_K^3}\xi_M(z,\nu)\epsilon_{\mu\nu\alpha\beta}p^{+\alpha}p^{-\beta}k^\nu\epsilon^\mu. \tag{9}$$

We use the notation of Ref. [13]: $\lambda\approx 0.22$ is the sine of the Cabibbo angle and $g_8\approx 5.1$. ξ_M is a form factor that depends on the photon energy in the kaon rest frame, $z=E_\gamma/M_K$, and on a pion energy difference, $\nu=(E_{\pi^+}-E_{\pi^-})/M_K$. At

[2]Except that we allow a larger uncertainty by doubling the maximum value used for the unknown counterterm that occurs in chiral perturbation theory [11].

leading order in chiral perturbation theory, it is just a constant which can be fit to the data to obtain [13]

$$\xi_M \equiv \frac{M_K^3}{8\pi^2 f_\pi^3} F_M e^{i(\delta_1^1-\delta_0^0)}$$

$$|F_M| = 0.94 \pm 0.06. \tag{10}$$

Unfortunately, the analysis of this decay mode is more complicated than this. There is experimental evidence for significant variation of the form factor with z. In particular, if we introduce only a slope term into the form factor,

$$\xi_M \equiv \frac{M_K^3}{8\pi^2 f_\pi^3} F_M (1 + c_M E_\gamma / M_K) e^{i(\delta_1^1-\delta_0^0)}. \tag{11}$$

E731 has extracted a value of $c_M = -1.7 \pm 0.5$ [14] which changes the overall constant to $F_M = 1.49 \pm 0.04$. A careful analysis in Ref. [15] parametrizes the long-distance contributions to F_M in terms of one constant (their k_F), finding that a typical range is $0.3 < F_{M,LD} < 0.9$. This analysis, however, might change in view of the latest results from E799 [16].

In this case it is not possible to separate long distance and short distance contributions to F_M. To be conservative, therefore, we demand that any new physics contribution to F_M be at most equal to the measured F_M. As we have seen, there is a large uncertainty at present in the extraction of F_M which can range from 0.3 to 1.5. For definiteness we use

$$|F_M|_{\rm new} < 1. \tag{12}$$

The decay $K^+ \to \pi^+ \pi^0 \gamma$ also has long distance contributions that are not known precisely and it is not as well measured as $K_L \to \pi^+ \pi^- \gamma$. The decay $K_L \to \pi^0 \pi^0 \gamma$ has not been seen, and in chiral perturbation theory, starts at order p^6 [17], making a separation of long and short distance contributions even harder. For these reasons we do not consider additional constraints from these modes.

To place a bound on new physics we need to estimate the contribution of the short distance operator of Eq. (1) to F_M. This requires a calculation of the matrix element $\langle \pi^+ \pi^- | \bar{d} \sigma_{\mu\nu}(1,\gamma_5) s | K \rangle$. Here we use naive dimensional analysis [18] to match the operator into a meson operator of the form

$$\mathcal{O}_M = g \frac{e G_F \lambda}{\sqrt{2}} \epsilon^{\mu\nu\alpha\beta} \partial_\mu \pi^+ \partial_\nu \pi^- F_{\alpha\beta} K. \tag{13}$$

Following Weinberg [18] we obtain the order of magnitude estimate [19]

$$g = \frac{4\pi}{\sqrt{2}\lambda} \frac{c}{\Lambda^2}. \tag{14}$$

There are many possible chiral Lagrangians that have the same transformation properties as the short distance operator, Eq. (1), that give rise to a meson operator like the one in Eq. (13). For example, guided by dimensional analysis, we can write down at $\mathcal{O}(p^4)$ the following effective Lagrangian:

$$\mathcal{L} = g \frac{e G_F}{\sqrt{2}} \lambda f_\pi^3 \epsilon^{\mu\nu\alpha\beta} {\rm Tr}[h \partial_\mu \Sigma \partial_\nu \Sigma^\dagger (\Sigma F_{R\alpha\beta} + F_{L\alpha\beta} \Sigma)] \tag{15}$$

where the matrix h is a 3×3 matrix with $h_{23}=1$ and all other elements zero to accomplish the $\Delta s = 1$ transition. Notice that Eq. (15) has the same transformation properties as the short distance operator and is not suppressed by light-quark masses as is appropriate for the new physics interactions of interest, Eq. (4). Within the standard model, the short distance operator is suppressed by light-quark masses and this results in a different effective chiral Lagrangian involving the usual chiral symmetry breaking factor. However, the standard model short distance coefficient is small enough to completely neglect its contribution. We have also included a Levi-Cività tensor to select the contribution of new physics to the magnetic transition. For the radiative decays in question we replace $F_{R\mu\nu} = F_{L\mu\nu} = eQF_{\mu\nu}$ where Q is the quark charge matrix and $F_{\mu\nu}$ the photon field strength tensor.

After using the chiral Lagrangian, Eq. (15), to calculate the amplitude for $K_L \to \pi^+ \pi^- \gamma$, we find

$$|F_M|_{\rm new} = \frac{\pi}{2\sqrt{2}\lambda} \frac{c}{f_\pi} \tag{16}$$

which, from Eq. (12), implies

$$|c(\mu)| \lesssim 18 \text{ MeV}. \tag{17}$$

This is comparable to the bound obtained from radiative hyperon decay, and can be improved with an improved determination of F_M.

It is easy to check that Eq. (15) gives a similar contribution to the decay $K^+ \to \pi^+ \pi^0 \gamma$, but as argued before, this does not place additional constraints. Similarly, Eq. (15) does not contribute to $K_L \to \pi^0 \pi^0 \gamma$ in accordance with the fact that this amplitude starts at order p^6.

V. RIGHT HANDED W COUPLINGS

Left-right mixing is the obvious case in which the light-quark mass suppression of the operator, Eq. (1), can be removed, and turned into an enhancement from a heavy top-quark mass. Left-right symmetric models have been studied in the context of $b \to s\gamma$ in detail [4]. For our purpose it suffices to look at the effective Lagrangian at the W mass scale that results after integrating out a heavy right handed W. This can be done easily following the formalism of Peccei and Zhang [20]. In unitary gauge the new coupling of interest is

$$\mathcal{L}_{\rm eff} = \frac{g}{\sqrt{2}} V_{UD} \kappa_R^{UD} \bar{U}_R \gamma^\mu D_R W_\mu^+ + \text{H.c.} \tag{18}$$

In writing Eq. (18) we have assumed CP conservation and ignored modifications to the left-handed W couplings which do not lead to enhanced effects. In general κ_R^{UD} will be different for each UDW coupling.

This interaction has been considered before in the study of $b \to s\gamma$ [5], and also in the study of CP violation in B decays [21]. A trivial generalization of those results leads to

$$c_{11}^i = -\frac{m_i}{m_s}\kappa_R^{is}F_{GR}(x_i),$$

$$c_{12}^i = -\frac{m_i}{m_s}\kappa_R^{is}F_{AR}(x_i), \tag{19}$$

where $i=u,c,t$, $x_i=(m_i/m_W)^2$ and

$$F_{GR}(x_i)=\frac{6x_i}{(1-x_i)^3}\log(x_i)+\frac{3(1+x_i)}{(1-x_i)^2}+1$$

$$F_{AR}(x_i)=\frac{x_i(2-3x_i)}{(1-x_i)^3}\log(x_i)+\frac{5x_i^2-31x_i+20}{6(1-x_i)^2}. \tag{20}$$

In this case, the intermediate charm and top quark states lead to a value of c much larger than in the standard model. If we assume that the V_{UD} matrix elements of Eq. (18) are the same as the CKM elements in the SM, and use the approximate Wolfenstein parametrization for the CKM angles with $\rho=0$, we find, at the W mass scale (in units of MeV),

$$c(m_W)\approx -10\kappa_R^{cs}-2\kappa_R^{ts}. \tag{21}$$

These numbers are reduced by about a factor of 2 at $\mu=1$ GeV when we use Eq. (3). We find $c(\mu)\approx -6\kappa_R^{cs}-1.4\kappa_R^{ts}$.

If we assume that κ_R^{cs}, κ_R^{ts} are of the same order of magnitude, then the charm-quark contribution is dominant. Our strongest constraint from the radiative hyperon decays, $|c(\mu)|<12$ MeV, implies that

$$|\kappa_R^{cs}|\lesssim 2$$

$$|\kappa_R^{ts}|\lesssim 8.5 \tag{22}$$

which are not as restrictive as $b\to s\gamma$. In fact, if we use the result of Ref. [22] and require that the predicted branching ratio to be within the 95% C.L. allowed range [23] $2\times10^{-4}<\mathrm{BR}(b\to s\gamma)<4.5\times10^{-4}$, we find

$$|\kappa_R^{cs}|\lesssim 0.95$$

$$|\kappa_R^{ts}|\lesssim 0.02. \tag{23}$$

Similarly, using the results of Ref. [21] and requiring that new contributions to K-$\bar{K}$ mixing be less than the standard model charm-quark contribution we obtain

$$|\kappa_R^{cs}|\lesssim 0.2$$

$$|\kappa_R^{ts}|\lesssim 2.2. \tag{24}$$

The constraints we obtain from radiative hyperon decay are not quite as good. Nevertheless, additional assumptions are implicit in Eqs. (23), (24). For example they have been obtained by taking only one non-zero coupling at a time and by ignoring possible counterterms in the effective theory at one-loop. For this reason we regard our constraints as being complementary to existing ones.

VI. SUPERSYMMETRIC MODELS

Another class of models in which the coefficient c is naturally large is the general supersymmetric (SUSY) extension of the standard model. In this class of models one can generate the operator at one loop via intermediate squarks and gluinos. The enhancement is due both to the strong coupling constant and to the removal of the chirality suppression that results in a gluino mass replacing the light-quark mass in Eq. (1) [24].

In the interaction basis, the down-quark-squark-gluino vertex is given by

$$\mathcal{L}=-\sqrt{2}g_s(\overline{d_L^i}T^a\tilde{g}_a\cdot\tilde{D}_L^i-\overline{d_R^i}T^a\tilde{g}_a\cdot\tilde{D}_R^i) \tag{25}$$

where i is the generation index. Soft SUSY breaking squark masses will in general induce mixing between different generations of squarks. The interaction eigenstates $\tilde{D}_{L,R}^i$ are then different from the mass eigenstates $\tilde{D}^k$ inducing flavor changing neutral currents. Here we stay away from specific models and follow Ref. [6] to write the contributions to the $s\to d\gamma$ transition in the mass insertion approximation. The result found in Ref. [6], after introducing an overall normalization to match our definition of $c_{11,12}$ in Eq. (1), is

$$c_{11}=\frac{\sqrt{2}\alpha_s\pi}{G_F m_{\tilde{q}}^2 V_{td}^* V_{ts}}\left[\delta_{12}^{LL}\left(\frac{1}{3}M_3(m_{\tilde{g}}^2/m_{\tilde{q}}^2)+3M_4(m_{\tilde{g}}^2/m_{\tilde{q}}^2)\right)\right.$$
$$\left.+\delta_{12}^{LR}\frac{m_{\tilde{g}}}{m_s}\left(\frac{1}{3}M_1(m_{\tilde{g}}^2/m_{\tilde{q}}^2)+3M_2(m_{\tilde{g}}^2/m_{\tilde{q}}^2)\right)\right]$$

$$c_{12}=\frac{\sqrt{2}\alpha_s\pi}{G_F m_{\tilde{q}}^2 V_{td}^* V_{ts}}\frac{8}{9}\left[\delta_{12}^{LL}M_3(m_{\tilde{g}}^2/m_{\tilde{q}}^2)\right.$$
$$\left.+\delta_{12}^{LR}\frac{m_{\tilde{g}}}{m_s}M_1(m_{\tilde{g}}^2/m_{\tilde{q}}^2)\right], \tag{26}$$

where the loop functions are given by

$$M_1(x)=\frac{1+4x-5x^2+2x(2+x)\ln x}{2(1-x)^4},$$

$$M_2(x)=-\frac{5-4x-x^2+2(1+2x)\ln x}{2(1-x)^4},$$

$$M_3(x)=\frac{-1+9x+9x^2-17x^3+6x^2(3+x)\ln x}{12(x-1)^5},$$

$$M_4(x)=\frac{-1-9x+9x^2+x^3-6x(1+x)\ln x}{6(x-1)^5}. \tag{27}$$

The parameters δ_{12}' characterize the mixing in the mass insertion approximation [6], $x=m_{\tilde{g}}^2/m_{\tilde{q}}^2$, and $m_{\tilde{g}}$, $m_{\tilde{q}}$ are the

TABLE III. c (in units of MeV) in generalized SUSY models due to δ_{12}^{LR} with $m_{\tilde{q}}=500$ GeV for different values of $x=m_{\tilde{g}}^2/m_{\tilde{q}}^2$.

x	$c(\mu)$	$c(m_W)$
0.3	88 δ_{12}^{LR}	73 δ_{12}^{LR}
1	57 δ_{12}^{LR}	60 δ_{12}^{LR}
4	24 δ_{12}^{LR}	32 δ_{12}^{LR}

average gluino and squark masses, respectively. In general the quark-squark-gluino interactions also generate the $s\rightarrow d\gamma$ transition with different chiralities. This can be obtained by replacing δ_{12}^{LL} and δ_{12}^{LR} by δ_{12}^{RR} and δ_{12}^{RL} in the above expressions.

As anticipated, there is an enhancement factor $m_{\tilde{g}}/m_s$ in the term proportional to δ_{12}^{LR} and this is the term with a potentially large contribution to $s\rightarrow d\gamma$.

In Table III we illustrate some representative values for c both at a scale $\mu=1$ GeV and at the W mass scale. The real part of δ_{12}^{LR} is constrained from $\Delta m_{K_S\text{-}K_L}$ to be typically less than a few times 10^{-2} [6]. This is to be compared with our constraint for $x=0.3$:

$$\delta_{12}^{LR}\lesssim 0.14. \tag{28}$$

The imaginary part is more severely constrained from ϵ'/ϵ [6]. Even though our constraint is not as good as that from $K\text{-}\bar{K}$ mixing the two are complementary in a general model. Our bound would be important, for example, in models that try to avoid the $K\text{-}\bar{K}$ mixing bound through an interplay of δ_{12}^{LR} and δ_{12}^{RL}. Notice that, unlike the case of right-handed W couplings, in the supersymmetric case $b\rightarrow s\gamma$ constrains different parameters of the model.

VII. CONCLUSIONS

We have studied radiative hyperon decays and radiative kaon decays of the form $K\rightarrow\pi\pi\gamma$ with the aim of constraining new interactions. Guided by chiral perturbation theory and dimensional analysis we have parametrized the long distance physics that dominates these modes and in this way we have quantified a bound on possible contributions from short-distance new physics. We found that these processes can place additional constraints on certain models in which the $s\rightarrow d\gamma$ transition is significantly enhanced with respect to the standard model. The constraints that we find are, in general, not competitive with those from other rare processes like $b\rightarrow s\gamma$ or $K\text{-}\bar{K}$ mixing. The constraints are, however, complementary and may become important for sufficiently detailed probes of new models.

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