Nonstandard Semileptonic Hyperon Decays

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We investigate the discovery potential of semileptonic hyperon decays in terms of searches of new physics at teraelectronvolt scales. These decays are controlled by a small $SU(3)$-flavor breaking parameter that allows for systematic expansions and accurate predictions in terms of a reduced dependence on hadronic form factors. We find that muonic modes are very sensitive to nonstandard scalar and tensor contributions and demonstrate that these could provide a powerful synergy with direct searches of new physics at the LHC.

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Introduction.—The meson and baryon semileptonic decays have played a crucial role in the discovery of the $V-A$ structure [1] and quark-flavor mixing [2] of the (charged current) electroweak interactions in the standard model (SM). From a modern perspective, high-precision measurements of these decays provide a benchmark to test the SM and complement the direct searches of new physics (NP) at teraelectronvolt (TeV) energies.

For example, the accurate determination of the elements V_{ud} and V_{us} of the Cabibbo-Kobayashi-Maskawa (CKM) matrix can be used to test its unitarity, constraining NP with characteristic scales as high as $\Lambda \sim 10$ TeV [3]. Furthermore, one can test the $V-A$ structure of the charged currents in $d \to u$ transitions using neutron and nuclear β decays [3–9] and pion decays [10,11]. Current limits for the associated NP scale are also at the TeV level, and important improvements are expected from future experiments [12]. Searches of nonstandard $d \to u$ transitions can also be done using LHC data, through, e.g., the collision of d and u partons in the $pp \to e^{\pm} + \mathrm{MET} + X$ channel (where MET stands for missing transverse energy) [12]. This leads to an interesting synergy between low- and high-energy NP searches in these flavor-changing processes.

A similar comprehensive analysis of exotic effects in $s \to u$ transitions has not been done yet. The (semi)leptonic kaon decays are optimal laboratories for this study due to the intense program of high-precision measurements and accurate calculations of the relevant form factors that has been carried out over the past decades [13]. Indeed, bounds on right-handed [14,15] or scalar and tensor [16] NP interactions at the 10^{-2}–10^{-3} level (relative to the SM) can be obtained [17,18]. Generally speaking, (pseudo) scalar and tensor operators modify the spectrum of the decay and a detailed knowledge of the q^2 dependence of the form factors becomes necessary [19].

In this Letter we investigate the physics potential of the semileptonic hyperon decays (SHD) to search for NP. Although the description of these modes may seem involved due to the presence of six nonperturbative matrix elements or form factors, they present interesting features [20–23]. (i) In the isospin limit, there are a total of 8 different channels, each having a differential decay rate with a rich angular distribution that could involve the polarizations of the baryons. (ii) The same form factors in different channels can be connected to each other and with other observables (e.g., electromagnetic form factors) in a model-independent fashion using the approximate $SU(3)$-flavor symmetry of QCD. (iii) The maximal momentum transfer is small compared to the baryon masses and it is parametrically controlled by the breaking of this symmetry. Therefore, a simultaneous $SU(3)$ breaking and "recoil" expansion can be performed that simplifies, systematically, the dependence of the decay rate on the form factors.

On the experimental side there is much room for improvement. Except for the measurements performed by the KTeV and NA-48 Collaborations in the $\Xi^0 \to \Sigma^+$ channel [24–28], most of the SHD data are more than 30 years old [29]. On the other hand, (polarized) hyperons could be produced abundantly in the NA62 experiment at CERN [30,31] or in any other hadron collider like the future $p\bar{p}$ facility PANDA [32] at FAIR/GSI or J-PARC [33].

In the following, we investigate the physics reach of the SHD with a discussion based on the sensitivity of the total decay rates to nonstandard scalar and tensor interactions. We show that the bounds from SHD are competitive with those derived from the LHC data on the $pp \to e^{\pm} +$ $\mathrm{MET} + X$ channel and leave the interplay with kaon decays for future work (see Refs. [14–17] for the current status).

SM effective field theory.—In the SM, and at energies much lower than the electroweak symmetry-breaking scale,

$v = (\sqrt{2}G_F)^{-1/2} \simeq 246$ GeV, all charged-current weak processes involving up and strange quarks are described by the Fermi $(V-A)\times(V-A)$ four-fermion interaction. Beyond the SM, the most general effective Lagrangian is [3]

$$\mathcal{L}_{\rm eff} = -\frac{G_F V_{us}}{\sqrt{2}}(1+\epsilon_L+\epsilon_R) \times \sum_{\ell=e,\mu}\{\bar{\ell}\gamma_\mu(1-\gamma_5)\nu_\ell \cdot \bar{u}[\gamma^\mu-(1-2\epsilon_R)\gamma^\mu\gamma_5]s + \bar{\ell}(1-\gamma_5)\nu_\ell\cdot\bar{u}[\epsilon_S-\epsilon_P\gamma_5]s + \epsilon_T\bar{\ell}\sigma_{\mu\nu}(1-\gamma_5)\nu_\ell\cdot\bar{u}\sigma^{\mu\nu}(1-\gamma_5)s\} + \text{H.c.}, \quad (1)$$

neglecting $\mathcal{O}(\epsilon^2)$ terms and derivative interactions, and where we use $\sigma^{\mu\nu} = [\gamma^\mu,\gamma^\nu]/2$. This Lagrangian has been constructed using only the SM fields relevant at low scales, $\mu \sim 1$ GeV, and demanding the operators to be color and electromagnetic singlets. Furthermore, we have restricted our attention to nonstandard interations that conserve lepton flavor and are lepton universal. Finally, we assume that the Wilson coefficients (WC) ϵ_i are real, since we focus on CP-even observables.

In light of the null results in direct searches of NP at colliders, we assume that its typical scale, Λ, is much larger than v. In such a case, NP can be parametrized using an effective (nonrenormalizable) Lagrangian, $\mathcal{L}_{\rm eff} = \mathcal{L}_{\rm SM} + (1/\Lambda^2)\sum_i \alpha_i \mathcal{O}_i^{(6)} + \cdots$, where the $\mathcal{O}_i^{(6)}$ are now operators built with *all* the SM fields and subject to the structures of its full (unbroken) gauge symmetry group [34,35]. The WC ϵ_i in Eq. (1) are generated by the high-energy WC α_i, which in turn can be obtained by matching to a particular NP model at $\mu = \Lambda$, and by running down to $\mu \sim 1$ GeV using the renormalization group equations, with the heavier fermions and weak bosons integrated out in the process [36–40].

This framework, usually referred to as the SM effective field theory (SMEFT), allows for a bottom-up investigation of NP, describing the implications of collider searches for low-energy experiments and vice versa. Needless to say, this interplay would become crucial in shaping the NP if a discrepancy with the SM is to be found. Examples of top-down applications, with correlated effects at high and low energies, can be found in scenarios with leptoquarks [41] or extra scalar fields [6,19].

Semileptonic hyperon decays.—Neglecting electromagnetic corrections, the amplitude for a particular SHD $B_1(p_1)\to B_2(p_2)\ell^-(p_\ell)\bar{\nu}_\ell(p_\nu)$ factorizes into the leptonic and baryonic matrix elements. For the (axial)vector hadronic currents we have the parametrization in terms of the standard form factors [22,42]:

$$\langle B_2(p_2)|\bar{u}\gamma_\mu s|B_1(p_1)\rangle = \bar{u}_2(p_2)\left[f_1(q^2)\gamma_\mu + \frac{f_2(q^2)}{M_1}\sigma_{\mu\nu}q^\nu + \frac{f_3(q^2)}{M_1}q_\mu\right]u_1(p_1), \quad (2)$$

$$\langle B_2(p_2)|\bar{u}\gamma_\mu\gamma_5 s|B_1(p_1)\rangle = \bar{u}_2(p_2)\left[g_1(q^2)\gamma_\mu + \frac{g_2(q^2)}{M_1}\sigma_{\mu\nu}q^\nu + \frac{g_3(q^2)}{M_1}q_\mu\right]\gamma_5 u_1(p_1); \quad (3)$$

whereas the nonstandard (pseudo)scalar and tensor interactions introduce new form factors [42]:

$$\langle B_2(p_2)|\bar{u}s|B_1(p_1)\rangle = f_S(q^2)\bar{u}_2(p_2)u_1(p_1), \quad (4)$$

$$\langle B_2(p_2)|\bar{u}\gamma_5 s|B_1(p_1)\rangle = g_P(q^2)\bar{u}_2(p_2)\gamma_5 u_1(p_1), \quad (5)$$

$$\langle B_2(p_2)|\bar{u}\sigma_{\mu\nu}s|B_1(p_1)\rangle \simeq f_T(q^2)\bar{u}_2(p_2)\sigma_{\mu\nu}u_1(p_1). \quad (6)$$

In Eqs. (2)–(6), $u_{1,2}$ are the parent and daughter baryon spinor amplitudes, $M_{1,2}$ their respective masses, and $q = p_1 - p_2$ is the momentum transfer, with $m_\ell^2 \le q^2 \le (M_1 - M_2)^2$. Furthermore, in Eq. (6) we have neglected other contributions to the matrix element of the tensor current since they are kinematically suppressed $\sim\mathcal{O}(q/M_1)$ [42].

A crucial aspect in the study of the SHD is the approximate $SU(3)$-flavor symmetry of QCD. It controls the phase space of the decay and allows for a systematic expansion of the observables in the *generic* symmetry-breaking parameter, $\delta = (M_1 - M_2)/M_1$ [21]. Furthermore, with currents J_b that transform as an octet under $SU(3)$ with flavor index b, one has $\langle B_a|J_b|B_c\rangle = F_J(q^2)f_{bac} + D_J(q^2)d_{bac}$, where $F_J(q^2)$ and $D_J(q^2)$ are reduced matrix elements, f_{abc} are the $SU(3)$ structure constants, and d_{abc} the so-called d coefficients (see, e.g., Ref. [43]). The $\sim\mathcal{O}(\delta)$ symmetry-breaking corrections can be calculated using model independent methods [44–51]. In addition, the form factors can be expanded around $q^2 = 0$ in powers of $q^2/M_X^2 \sim \delta^2$, where $M_X \sim 1$ GeV is an hadronic scale related to the mass of the resonances coupling to the currents [52,53].

Let us illustrate this with the total decay rate for the electronic mode in the SM which, expanded up to next-to-leading order (NLO) in δ and neglecting m_e, is [21]

$$\Gamma_{e,\rm SM} \simeq \frac{G_F^2|V_{us}f_1(0)|^2\Delta^5}{60\pi^3}\left[\left(1-\frac{3}{2}\delta\right) + 3\left(1-\frac{3}{2}\delta\right)\frac{g_1(0)^2}{f_1(0)^2} - 4\delta\frac{g_2(0)}{f_1(0)}\frac{g_1(0)}{f_1(0)}\right], \quad (7)$$

with $\Delta = M_1 - M_2$. This expression contains a minimal dependence on the form factors. No information on their q^2 dependence is required and, moreover, the last term can be neglected because the weak-electric charge, $g_2(0)$, is itself $\mathcal{O}(\delta)$ [42]. Thus, besides G_F and V_{us}, and up to a theoretical accuracy of $\mathcal{O}(\delta^2) \sim 1\%$–$5\%$, the total decay rate of the electronic mode in the SM only depends on hyperon vector and axial charges, $f_1(0)$ and $g_1(0)$. Equation (7) makes manifest that $f_1(0)$ is essential for extracting V_{us} from the rates, while the ratio $g_1(0)/f_1(0)$ can be obtained measuring the angular distribution of the final lepton [21,22]. Neglected electromagnetic corrections are of a few percent [21,54], well within the accuracy achieved at NLO in the $SU(3)$ expansion.

Beyond the SM, we generally have two types of effects. On one hand, (axial)vector modifications to the SM, described by the WC $\epsilon_{L,R}$, can be arranged [cf. Eq. (1)] into a change of the normalization of the rate according to the replacement $V_{us} \to \tilde{V}_{us} = (1 + \epsilon_L + \epsilon_R)V_{us}$ and of the axial coupling to the leptonic current by the factor $(1 - 2\epsilon_R)$. The former combination involves a modification of V_{us} which has been tightly constrained by testing CKM unitarity [3]. The latter could be determined in SHD from the measured $g_1(0) \to \tilde{g}_1(0) = (1 - 2\epsilon_R)g_1(0)$ only if $g_1(0)$ was known accurately from QCD (for recent progress in the lattice, see Refs. [55,56]).

On the other hand, the WC $\epsilon_{S,P,T}$ introduce new structures in the energy and angular distributions. Restricting ourselves to $\mathcal{O}(v^2/\Lambda^2)$ (or linear in the WC), they appear from the interference of the NP terms with the SM, and the contributions of the (pseudo)scalar and tensor operators are suppressed by $m_\ell/\sqrt{q^2}$. Therefore, while the electronic channels can be analyzed specifically to measure and study the normalization of the rates $|\tilde{V}_{us}f_1(0)|$ and the relevant form factors, the muonic modes could use the information thus obtained to constrain the (pseudo)scalar and tensor operators. Besides that, it is important to note that the pseudoscalar quark bilinear receives a kinematical $\mathcal{O}(q/M_1)$ suppression that largely neutralizes the sensitivity of SHD to ϵ_P (see, however, Ref. [9]). For this reason, we center our discussion below on the study of ϵ_S and ϵ_T.

We expand the contributions in the SM up to $\mathcal{O}(\delta)$, but we keep only the leading terms in the NP terms. This implies a relative $\mathcal{O}(\delta^2)$ error in the SM predictions, which we fix to 5% in all channels for definiteness, and an uncertainty $\mathcal{O}(\delta) \sim 10\%$–$20\%$ in the sensitivity to NP that will not affect the conclusions of our analysis.

Bounds on scalar and tensor operators.—Let us now introduce the ratio

$$R^{\mu e} = \frac{\Gamma(B_1 \to B_2 \mu^- \bar{\nu}_\mu)}{\Gamma(B_1 \to B_2 e^- \bar{\nu}_e)}. \tag{8}$$

This observable is not only sensitive to lepton-universality violation but is also linearly sensitive to ϵ_S and ϵ_T. In addition, one expects the dependence on the form factors in the SM to simplify in the ratio. In fact, working at NLO we obtain

$$R^{\mu e}_{\rm SM} = \sqrt{1 - \frac{m_\mu^2}{\Delta^2}}\left(1 - \frac{9}{2}\frac{m_\mu^2}{\Delta^2} - 4\frac{m_\mu^4}{\Delta^4}\right) + \frac{15}{2}\frac{m_\mu^4}{\Delta^4}\,\mathrm{arctan\,h}\left(\sqrt{1 - \frac{m_\mu^2}{\Delta^2}}\right). \tag{9}$$

This is a remarkable result: up to a relative theoretical accuracy of $\mathcal{O}(\delta^2)$, $R^{\mu e}$ in the SM does not depend on any form factor. In Table I we compare the experimental ratios

TABLE I. Comparison between the predictions of $R^{\mu e}$ in the SM at NLO and experimental measurements for different SHD.

	$\Lambda \to p$	$\Sigma^- \to n$	$\Xi^0 \to \Sigma^+$	$\Xi^- \to \Lambda$
Experiment	0.189(41)	0.442(39)	0.0092(14)	0.6(5)
SM NLO	0.153(8)	0.444(22)	0.0084(4)	0.275(14)

to the results predicted in the SM. As discussed above, the main reason for the large experimental errors is that most of the data in the muonic channel are very old and scarce. At this level of precision, which generously covers the theoretical accuracy attained by Eq. (9), we observe that the experimental data on $R^{\mu e}$ agree with the SM.

One can now use this consistency of the data with the SM to set bounds on the WC of the scalar and tensor operators, which generate the following nonstandard contribution:

$$R^{\mu e}_{\rm NP} \simeq \frac{\left(\epsilon_S \frac{f_S(0)}{f_1(0)} + 12\epsilon_T \frac{g_1(0)}{f_1(0)}\frac{f_T(0)}{f_1(0)}\right)}{(1 - \frac{3}{2}\delta)\left(1 + 3\frac{g_1(0)^2}{f_1(0)^2}\right)}\Pi(\Delta, m_\mu), \tag{10}$$

where $\Pi(\Delta, m_\mu)$ is a phase-space integral:

$$\Pi(\Delta, m_\mu) = \frac{5}{2}\frac{m_\mu}{\Delta}\left[\left(2 + 13\frac{m_\mu^2}{\Delta^2}\right)\sqrt{1 - \frac{m_\mu^2}{\Delta^2}} - 3\left(4\frac{m_\mu^2}{\Delta^2} + \frac{m_\mu^4}{\Delta^4}\right)\mathrm{arctan\,h}\left(\sqrt{1 - \frac{m_\mu^2}{\Delta^2}}\right)\right]. \tag{11}$$

It is particularly convenient to express the dependence on the WC in "units" of the SM ratio:

$$\frac{R^{\mu e}}{R^{\mu e}_{\rm SM}} = 1 + r_S\epsilon_S + r_T\epsilon_T, \tag{12}$$

where $r_{S,T}$ are dimensionless numbers encapsulating the net sensitivity to the WC.

The values of the form factors that we use to calculate $r_{S,T}$ are given in Table II. The ratio $g_1(0)/f_1(0)$ is measured from the angular distribution of the electronic channels [29]. The scalar form factor can be obtained, up to electromagnetic corrections, using the conservation of vector current in QCD, $f_S(0)/f_1(0) = \Delta/(m_s - m_u)$ [9]. For the tensor form factors we need to use model calculations [57], whose errors are difficult to quantify. Nevertheless, it is interesting to note

TABLE II. SHD data for $g_1(0)/f_1(0)$ and theoretical determinations of $f_{S,T}(0)/f_1(0)$ at $\mu = 2$ GeV used in this work. The corresponding $r_{S,T}$ are shown in the last two lines.

	$\Lambda \to p$	$\Sigma^- \to n$	$\Xi^0 \to \Sigma^+$	$\Xi^- \to \Lambda$
$g_1(0)/f_1(0)$	0.718(15)	−0.340(17)	1.210(50)	0.250(50)
$f_S(0)/f_1(0)$	1.90(10)	2.80(14)	1.36(7)	2.25(11)
$f_T(0)/f_1(0)$	0.72	−0.28	1.22	0.22
r_S	1.60	4.1	0.56	3.7
r_T	5.2	1.7	7.2	1.1

that the tensor form factor for the neutron β decay is predicted to be 1.22, which is in the ballpark of the values obtained in the lattice [6,56,58–60]. This situation should be easily improved by future lattice calculations of the hyperon decay tensor charges.

The sensitivities to $\epsilon_{S,T}$ exhibited by the SHD (last two lines of Table II) are strongly channel dependent. In Fig. 1, we show 90% confidence level contours in the (ϵ_S, ϵ_T) plane using a χ^2 that includes the experimental measurements of $R^{\mu e}$ and where we propagate the experimental and theoretical uncertainties of the SM predictions in quadratures. For $r_{S,T}$ we use the values in Table II. As we can see, even though the experimental data on $R^{\mu e}$ are not precise, the strong sensitivity of SHD to NP leads to stringent bounds in $\epsilon_{S,T}$, namely,

$$\epsilon_S = 0.003(40), \qquad \epsilon_T = 0.017(34), \tag{13}$$

at 90% C.L. Accounting for the running of $\epsilon_{S,T}$ on the renormalization scale μ [61], and assuming natural values for the WC at $\mu = \Lambda$, these bounds translate into $\Lambda \sim v(V_{us}\epsilon_{S,T})^{-1/2} \sim 2$–4 TeV [12].

Limits from LHC data.—As discussed above, the SMEFT allows us to interpret model-independently high- and low-energies searches of NP. In particular, the cross section $\sigma(pp \to e + \text{MET} + X)$ with transverse mass higher than $\bar{m}_T$ is modified by nonstandard $\bar{u}s \to e\bar{\nu}$ partonic interactions as follows:

$$\sigma(m_T > \bar{m}_T) = \sigma_W + \sigma_S|\epsilon_S|^2 + \sigma_T|\epsilon_T|^2, \tag{14}$$

where $\sigma_W(\bar{m}_T)$ represents the SM contribution and $\sigma_{S,T}(\bar{m}_T)$ are new functions, whose explicit form can be found in Ref. [12] (up to trivial flavor indices changes). Thus, comparing the observed events above $\bar{m}_T$ with the SM expectation, we can set bounds on $\epsilon_{S,T}$. In particular, one (three) event is found with a transverse mass above $\bar{m}_T = 1.5$ TeV (1.2 TeV) in the 20 fb^{-1} (5 fb^{-1}) data set recorded at $\sqrt{s} = 8$ TeV (7 TeV) by the CMS Collaboration [62,63], in good agreement with the SM background of 2.02 ± 0.26 (2.8 ± 1.0) events. Using Eq. (14), this agreement translates in the 90% C.L. limits on $\epsilon_{S,T}$ shown in Fig. 1. We use the MSTW2008 PDF sets evaluated at $Q^2 = 1\ \text{TeV}^2$ [64] to calculate $\sigma_{S,T}$. Further details can be found in Ref. [12].

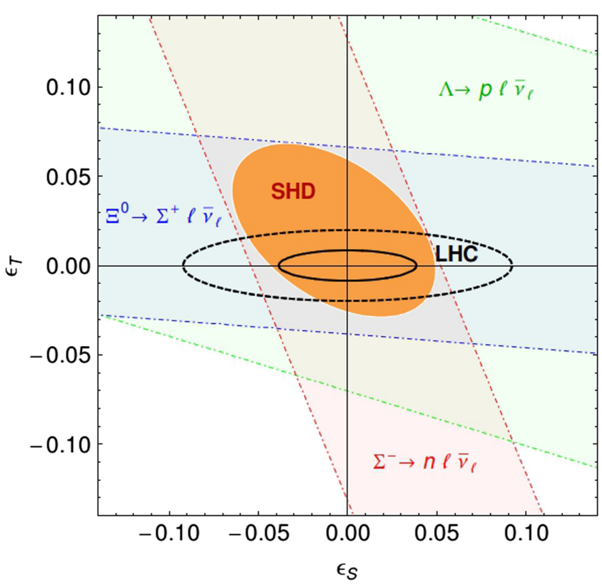

FIG. 1 (color online). 90% C.L. constraints on $\epsilon_{S,T}$ at $\mu = 2$ GeV from the measurements of $R^{\mu e}$ in different channels (dot-dashed lines) and combined (filled ellipse). LHC bounds obtained from CMS data at $\sqrt{s} = 8$ TeV (7 TeV) are represented by the black solid (dashed) ellipse.

Figure 1 illustrates the interesting competition that future SHD measurements could have with LHC searches of NP affecting $s \to u$ transitions. It is important to note that the dependence of the cross section [Eq. (14)] on the WC is quadratic, whereas in SHD is linear. In addition to reducing the sensitivity of the future collider searches of NP in this channel, one might also need to consider possible cancellations with linear effects from dimension-8 operators in the SMEFT.

Conclusions and outlook.—In summary, the most important features of SHD in relation to searches of NP at TeV scales are as follows. (i) The SHD are controlled by a small $SU(3)$-breaking parameter, allowing for systematic expansions that lead to accurate expressions in terms of a reduced dependence on form factors, cf. Eq. (9). (ii) The interference of the (pseudo)scalar and tensor NP operators with the SM in the rate is chirally suppressed. Therefore, electronic modes are well suited to measure normalization factors $|\tilde{V}_{us}f_1(0)|$, NP-modified $\tilde{g}_1(0)$, and other form factors. (iii) The muonic modes, on the other hand, show a strong linear sensitivity to scalar and tensor contributions that depend on the different combinations of form factors in each channel. This allows us to constrain them using SHD alone, with a precision that is competitive with the LHC data, cf. Fig. 1 and Eq. (13).

Our hope is that the present study triggers a program of high-precision measurements of different observables in the SHD. Hyperons can be produced in great numbers in current [30,33] and future facilities [32]. One may also wonder if better measurements could be extracted from the analysis of the data collected in past experiments like HyperCP [65], KTeV, and NA48. Any development on the experimental side will directly improve the bounds on NP obtained in this work with an observable as simple as $R^{\mu e}$, and using data with ~10%–20% relative errors.

Future improvements on the experimental precision will need to be accompanied by similar efforts on the theory side. In particular, the inclusion of $\mathcal{O}(\delta^2)$ terms in the SM predictions would improve the accuracy to the 1% level or below. Besides that, further nonperturbative calculations of the tensor form factors would improve the assessment of the sensitivity to ϵ_T. Finally, it will be important to perform this comprehensive analysis of the

SHD in complementarity with the kaon decays. Work along these lines is in progress.

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