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STATUS OF CP VIOLATION IN HYPERON DECAYS

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ABSTRACT

We update the estimate of CP violating observables in hyperon decays. We consider the effect of a heavy top quark, as well as recent bounds on the CKM angles. We present an estimate of the uncertainties associated with the evaluation of the hadronic matrix elements by comparing different models. We find that $A(\Lambda^0_-)$ is likely to occur at the 10^{-5} level within the standard model, although it could reach 10^{-4} .

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Of the current efforts to search for physics beyond the standard model, a very important one is the identification of systems where the observation of CP violation is feasible in the near future. After more than 20 years of experimental efforts, the only evidence for CP violation remains the non-zero value of the parameter ϵ in K^0 decays. The observation of direct CP violation in $K \rightarrow \pi\pi$ remains elusive.[§] Recent results suggest that ϵ' could be $(1.0 \pm 0.5) \times 10^{-3}$ or much smaller if the top quark mass is around 220 GeV^[3] within the standard model. This gives renewed importance to the study of different systems sensitive to direct CP violation. In particular, one would like to know if the presence of a very heavy top quark significantly affects other observables as it does ϵ' .

Interest in CP violation in hyperon decays was revived a few years ago, and some signals were estimated to be of order 10^{-4} within the standard model.^[4] It was also proposed that a low-energy $p\bar{p}$ machine would be an ideal place to search for these signals.^[7] Experiments have only been able to place bounds at the few $\times 10^{-2}$ level so far, but there are studies considering dedicated searches in a machine like SUPERLEAR, or a τ -charm factory. The purpose of this letter is to present an update of the theoretical estimate for some of these signals.

We will concentrate on $\Lambda \rightarrow p\pi^-$ and $\Xi^- \rightarrow \Lambda\pi^-$ decays. The first one is the most likely one to be studied, and the second one allows one to study different types of signals. They would both be of interest in the reactions: $p\bar{p} \rightarrow \Lambda\bar{\Lambda}$ or $p\bar{p} \rightarrow \Xi\bar{\Xi}$ for a machine like SUPERLEAR; and in $J/\Psi \rightarrow \Lambda\bar{\Lambda}$ or $J/\Psi \rightarrow \Xi\bar{\Xi}$ for a τ -charm factory.

For a hyperon non-leptonic decay $B_i(k, w_i) \rightarrow B_f(p, w_f)\pi$, where k, p denote the momenta and w_i, w_f the polarization vectors of the initial and final hyperons, one writes an invariant transition matrix element as:

$$M = G_F m_\pi^2 \bar{u}_f (A - B\gamma_5) u_i. \quad (1)$$

[§] There are at present two different results. The NA31 collaboration has published^[1] $\epsilon' = (3.3 \pm 1.1) \times 10^{-3}$, and it now reports $\epsilon' = (2.3 \pm 0.7) \times 10^{-3}$. The FNAL group has published^[2] $\epsilon' = (-0.4 \pm 1.4) \times 10^{-3}$, and now reports $\epsilon' = (6.0 \pm 5.8 \pm 3.2 \pm 1.8) \times 10^{-4}$.

This leads to a differential decay rate in the rest frame of the initial hyperon that is proportional to:

$$R \sim 1 + \gamma \bar{w}_f \cdot \bar{w}_i + (1 - \gamma) \bar{w}_f \cdot \bar{p} \bar{w}_i \cdot \bar{p} + \alpha (\bar{w}_f \cdot \bar{p} + \bar{w}_i \cdot \bar{p}) + \beta \bar{p} \cdot \bar{w}_f \times \bar{w}_i \quad (2)$$

Defining $S = A$ and $P = |\bar{p}|B/(E_f + m_f)$, the usual S, P -wave amplitudes, one finds:

$$\alpha = \frac{2\text{Re}S^*P}{|S|^2 + |P|^2}, \quad \beta = \frac{2\text{Im}S^*P}{|S|^2 + |P|^2}, \quad \gamma = \frac{|S|^2 - |P|^2}{|S|^2 + |P|^2}. \quad (3)$$

This is the standard result^[6] for the parameters describing the angular distribution of this decay.[¶]

To study the behavior of this differential decay rate under a CP transformation, it is convenient to write the amplitudes S, P as:

$$S = \sum_i S_i e^{i(\phi_i^S + \delta_i^S)} \\ P = \sum_i P_i e^{i(\phi_i^P + \delta_i^P)} \quad (4)$$

to explicitly separate the strong rescattering phases δ_i (which are the same in the corresponding anti-hyperon decay) and the weak CP violating phases ϕ_i (which change sign in the anti-hyperon decay). We use the index i to represent the different isospin amplitudes.

By comparing the hyperon and anti-hyperon decays one then finds three CP odd quantities:

$$\Delta = \frac{\Gamma - \bar{\Gamma}}{\Gamma + \bar{\Gamma}} \\ A = \frac{\Gamma_\alpha + \bar{\Gamma}_\alpha}{\Gamma_\alpha - \bar{\Gamma}_\alpha} \simeq \frac{\alpha + \bar{\alpha}}{\alpha - \bar{\alpha}} \\ B = \frac{\Gamma_\beta + \bar{\Gamma}_\beta}{\Gamma_\beta - \bar{\Gamma}_\beta} \simeq \frac{\beta + \bar{\beta}}{\beta - \bar{\beta}} \quad (5)$$

[¶] Recall that only two of them are independent, and that usually α and ϕ are taken to be the independent ones, where ϕ is defined by $\beta = (1 - \alpha^2)^{1/2} \sin \phi$.

In the above expressions we used Γ for the decay rate, and we used a bar to distinguish the quantities associated with the anti-hyperon decay. General expressions for these quantities can be found in the literature.^[6]

With the knowledge that the strong and weak phases are small, and that the $\Delta I = 3/2$ amplitudes are much smaller than the corresponding $\Delta I = 1/2$ amplitudes one can write approximate expressions to leading order in all the small parameters. Using the notation of Ref. 9 : $S_{ij} = S_{2\Delta I, 2I}$, $P_{ij} = P_{2\Delta I, 2I}$ and δ_{2I} , $\delta_{2I,1}$ for S, P wave amplitudes^[6] one finds:

$$\begin{aligned} \Delta(\Lambda^0_-) &= \sqrt{2} \frac{S_{33}}{S_{11}} \sin(\delta_3 - \delta_1) \sin(\phi_3^S - \phi_1^S) \\ A(\Lambda^0_-) &= -\tan(\delta_{11} - \delta_1) \sin(\phi_1^P - \phi_1^S) \\ B(\Lambda^0_-) &= \cot(\delta_{11} - \delta_1) \sin(\phi_1^P - \phi_1^S) \\ \Delta(\Xi^-) &= 0 \\ A(\Xi^-) &= -\tan(\delta_{21} - \delta_2) \sin(\phi_{12}^P - \phi_{12}^S) \\ B(\Xi^-) &= \cot(\delta_{21} - \delta_2) \sin(\phi_{12}^P - \phi_{12}^S) \end{aligned} \quad (6)$$

Examining these expressions one can see why one has that $B > A > \Delta$ in general. While the three observables are of course proportional to small CP violating phases, $\sin \phi$, both Δ and A have additional suppression factors. A is also proportional to the strong phases, $\tan \delta$, (which are known experimentally to be small) reflecting the fact that this is a naive-T even correlation. Δ is also proportional to these strong phases (again, because it is a naive-T even observable); and is also suppressed by the factor S_{33}/S_{11} revealing its origin in the interference of a $\Delta I = 1/2$ and a small $\Delta I = 3/2$ amplitudes. This factor does not appear in the expressions for A and B , because these ones originate in the interference of s and p waves of the same isospin. The expressions that appear in Ref. 6 are more complicated than the ones we have written. We have not written corrections to the leading terms for A and B that occur through interference of s and p wave amplitudes of

different isospin, since these have an additional $\Delta I = 3/2$ suppression factor.*

Experimentally, only $A(\Lambda^0_-)$ has been searched for, with the results $A(\Lambda^0_-) = -0.07 \pm 0.09$ from LEAR^[10], and $A(\Lambda^0_-) = 0.01 \pm 0.10$ from J/Ψ decays^[11].

To calculate the CP odd observables, we need to know all the quantities that appear in Eq. (6). The magnitude and strong scattering phases of the different amplitudes, S_i , P_i ^[12] δ_i ^[13] are extracted from experiment. We thus need to calculate the CP violating weak phases ϕ_i , and relate them to the CP violating phase in the CKM matrix, since we will only study the minimal standard model. To do this, we consider the relevant $|\Delta S| = 1$ effective weak Hamiltonian:

$$\mathcal{H}_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{ud}^* V_{us} \sum_i C_i(\mu) Q_i(\mu) \quad (7)$$

where the sum is over all the $Q_i(\mu)$ four-quark operators, and the $C_i(\mu)$ are the Wilson coefficients. We follow the notation of Ref. 5, and using $\alpha(m_W) = \frac{1}{128}$ and $m_W = 80$ GeV, we reproduce their results. We will only be interested in the CP-violating imaginary part of the coefficients $C_i(\mu)$, since we will be using the experimental numbers for the CP conserving amplitudes. Following Ref. 5 we write $C_i(\mu) = z_i(\mu) + \tau y_i(\mu)$; where $\tau = -V_{td}^* V_{ts}/V_{ud}^* V_{us}$ and the V_{ij} are the usual CKM matrix elements. We have:

$$\text{Im } M = \langle \pi B_f | \frac{G_F}{\sqrt{2}} V_{ud}^* V_{us} \sum_i \text{Im } \tau y_i(\mu) Q_i(\mu) | B_i \rangle \quad (8)$$

The quantities $y_i(\mu)$ were calculated by taking $y_i(m_W)$ as an initial value and then using the renormalization group equations to reach the scale μ . We have done this following Ref. 5, consistently including electroweak penguins.

* They would be important if, for example, the phase difference $\phi_1^S - \phi_1^P$ turned out to be significantly smaller than $\phi_1^P - \phi_3^S$. We have checked numerically that this doesn't occur in the models we consider.

To restrict the values of the V_{ij} CKM angles that appear in τ we used several experimental constraints. 1) The B-meson lifetime, which leads to $|V_{cb}| = 0.044 \pm 0.009$; 2) The new measurement of $R = \frac{\Gamma(B \rightarrow \tau \nu)}{\Gamma(B \rightarrow e \nu)}$.^[14] Using semileptonic B-decays and taking into account the different models this gives: $|V_{cb}/V_{cb}| = 0.09 \pm 0.04$. 3) $\text{Im}\tau$ is restricted by the parameters ϵ of CP violation in the $K^0 - \bar{K}^0$ system^[14]: $|\epsilon| = (2.27 \pm 0.02) \times 10^{-3}$, and the observed mixing in the $B_d - \bar{B}_d$ system^[14]: $\chi_d = 0.66 \pm 0.11$.

In our analysis, the form factor B_K entering the calculation of ϵ is allowed to vary between $1/3 \leq B_K \leq 1$, and the factor $\sqrt{B_B F_B}$ entering the calculation of χ_d is allowed to vary between $1.3 F_\pi \leq |\sqrt{B_B F_B}| \leq 2.1 F_\pi$. We have performed our analysis by using a program in which the top quark mass is kept as a free parameter. A given set of numbers for $|V_{cb}|$, $|V_{ub}/V_{cb}|$, B_K and $\sqrt{B_B F_B}$ within the allowed ranges is used, such that we fit ϵ and χ_d to the experimental values with a $\chi^2 \leq 9$ (which is approximately at the 99% c.l.). There are two allowed ranges^[14] for the CP violating phase δ , corresponding to it being in the first or second quadrant*. It turns out that for the larger values of m_t , only the first quadrant for δ is allowed.

Once we have fixed the short distance Wilson coefficients and the CKM angles appearing in Eq. (8), we are left with the outstanding problem of computing the matrix elements of the four-quark operators. We will use a few models for our estimates, and present results for two of them. Another approach consists of trying to write an effective Lagrangian for the ground state octet of baryons and the pseudoscalar meson octet. This approach will not be very useful to us, mainly because there seem to be a large number of operators that one needs to consider.^{††}

Our approach to the hadronic matrix elements will consist of using the experimental values^[17] for the different amplitudes assuming them to be real, and then to calculate the weak phases by using the imaginary part of the amplitude that is predicted by different models in the literature.

* Here δ is the CP violating phase in the Maiani parameterization of the CKM matrix.

† Recall that within the effective Lagrangian approach each operator contains an unknown coupling constant that must be fit from experiment.

Let us first consider the weak phases in vacuum saturation.^[18] Starting with the effective weak hamiltonian of Eq. (7) we compute the amplitude for the decay $B_i \rightarrow B_f \pi^-$ by factorizing the four-quark operators. In this form, our result applies both for $\Lambda \rightarrow \pi^-$ and for $\Xi^- \rightarrow \Lambda \pi^-$.

Our result in vacuum saturation[†] is given by:

$$M(B_i \rightarrow B_f \pi) = \frac{G_F}{\sqrt{2}} V_{ud}^* V_{us} \left((M_1^S + M_3^S) V + (M_1^P + M_3^P) A \right)$$

with

$$\begin{aligned} M_1^S &= \left[\frac{(C_1 - 2C_2)}{3} - \frac{C_7}{2} + \xi \left(\frac{-2C_1 + C_2}{3} - C_3 - \frac{C_8}{2} \right) \right] \\ &\quad - \frac{2m_\pi^2}{(m_u + m_d)(m_s - m_u)} \left[C_6 + \frac{C_8}{2} + \xi(C_5 + \frac{C_7}{2}) \right] \\ M_1^P &= \left[\frac{(C_1 - 2C_2)}{3} - \frac{C_7}{2} + \xi \left(\frac{-2C_1 + C_2}{3} - C_3 - \frac{C_8}{2} \right) \right] \\ &\quad + \frac{2m_\pi^2}{(m_u + m_d)(m_s + m_u)} \left[C_6 + \frac{C_8}{2} + \xi(C_5 + \frac{C_7}{2}) \right] \\ M_3^S &= \left[\frac{-(C_1 + C_2)}{3} (1 + \xi) + \frac{C_7}{2} + \xi \frac{C_8}{2} \right] - \frac{m_\pi^2}{(m_u + m_d)(m_s - m_u)} (\xi C_7 + C_8) \\ M_3^P &= \left[\frac{-(C_1 + C_2)}{3} (1 + \xi) + \frac{C_7}{2} + \xi \frac{C_8}{2} \right] + \frac{m_\pi^2}{(m_u + m_d)(m_s + m_u)} (\xi C_7 + C_8) \end{aligned} \quad (9)$$

We have used the fact that in vacuum saturation, tadpoles cancel to leading order.[‡] We have also parameterized the matrix elements of scalar and pseudoscalar densities in terms of those of vector and axial vector currents and quark masses as:

$$\begin{aligned} \langle \pi^- | \bar{d} \gamma_\mu u | 0 \rangle \langle B_f | \bar{u} s | B_i \rangle &= \frac{m_\pi^2}{(m_u + m_d)(m_s - m_u)} V \\ \langle \pi^- | \bar{d} \gamma_5 u | 0 \rangle \langle B_f | \bar{u} \gamma_5 s | B_i \rangle &= - \frac{m_\pi^2}{(m_u + m_d)(m_s + m_u)} A \end{aligned} \quad (10)$$

† In vacuum saturation one should take $\xi = 1/3$. We have left it as an arbitrary parameter, since some models advocate using factorization with $\xi = 0$.^[19]

and we have defined:

$$\begin{aligned} V &\equiv (\pi^- | \bar{d}_{\mu} \gamma_5 u | 0) \langle B_f | \bar{u} \gamma^\mu s | B_i \rangle \\ A &\equiv (\pi^- | \bar{d}_{\mu} \gamma_5 u | 0) \langle B_f | \bar{u} \gamma^\mu \gamma_5 s | B_i \rangle \end{aligned} \quad (11)$$

To calculate V and A , we use, for $\Lambda \rightarrow p\pi^-$:

$$\begin{aligned} V &= i\sqrt{2}F_\pi(m_\Lambda - m_p)f_1(m_\pi^2)\bar{\Psi}_p\Psi_\Lambda \\ A &= i\sqrt{2}F_\pi\left(-(m_\Lambda + m_p)g_1(m_\pi^2) + m_\pi^2g_3(m_\pi^2)\right)\bar{\Psi}_p\gamma_5\Psi_\Lambda \\ &= i\sqrt{2}F_\pi(p|q_\mu\bar{u}\gamma^\mu\gamma_5s|\Lambda) \\ &\simeq -i\frac{2F_\pi F_K m_\pi^2}{m_K^2 - m_\pi^2}g_{\Lambda p K^-}\bar{\Psi}_p\gamma_5\Psi_\Lambda \end{aligned} \quad (12)$$

and similarly for $\Xi^- \rightarrow \Lambda\pi^-$.

In the above expressions we have $F_\pi = 93 \text{ MeV}$, $F_K \simeq 1.3F_\pi$ and $f_1 = \sqrt{3/2}$ (for $\Lambda \rightarrow p\pi^-$).[§] The value of g_3 is not known, so we estimated the second matrix element assuming kaon pole dominance. For that we used $g_{\Lambda p K^-} \simeq -g_{\pi NN} \simeq -13.3$.

Although our complete result will be shown in Figs. 1, 2 we give here a simplified expression with the dominant terms for the weak phases in Eq. (6):

$$\begin{aligned} \phi_1^S &\simeq -0.3 \frac{m_\pi^2}{(m_u + m_d)m_s} y_6 \text{Im}\tau \\ \phi_3^S &\simeq \left(3.56(y_1 + y_2) + 8.22\left(\frac{1}{3}y_7 + y_8\right)\frac{m_\pi^2}{(m_u + m_d)m_s}\right) \text{Im}\tau \\ \phi_1^P &\simeq -0.03 \frac{m_\pi^2}{(m_u + m_d)m_s} y_6 \text{Im}\tau \end{aligned} \quad (13)$$

The difference $\phi_3^S - \phi_1^S$ arises from the interference of the two isospin amplitudes, and it is the one that could be very sensitive to the value of the top quark mass,

[§] f_1, g_1, g_3 are the usual form factors appearing in the baryonic matrix elements of the vector and axial vector currents^[6]. We have not included f_3 and g_2 that correspond to "second class currents". The term proportional to f_2 vanishes after contraction with q_μ .

in analogy with ϵ'/ϵ . This is easy to understand, as the effect of a heavy top is to enhance contributions of electro-weak penguin diagrams. These are important for $\Delta I = 3/2$ amplitudes but not so much for $\Delta I = 1/2$ amplitudes. We can in fact check that these two phases (ϕ_1^S, ϕ_3^S) do have the same sign and that ϕ_3^S approaches the value of ϕ_1^S for a top quark mass of around 250 GeV , suggesting that $\Delta(\Lambda_-^0)$ may indeed have the same cancellation that occurred for ϵ'/ϵ . This is not the case for $A(\Lambda_-^0)$. We can see this explicitly in Figs. 1, 2. In those figures we can also see that there is a reduction of the signals by a factor of about two when going from a renormalization scale $\mu = 0.5 \text{ GeV}$ to $\mu = 1.0 \text{ GeV}$. It is difficult to justify using perturbative results at such a low scale as $\mu = 0.5 \text{ GeV}$, however, we have included those values since they are the ones used in most models. This gives us an idea of the uncertainty involved in not knowing where to match the perturbative calculation to the estimates of the matrix elements.

Factorization of the effective weak Hamiltonian with the use of vacuum saturation is a very crude method of estimating the amplitudes, and it fails to reproduce the experimental results. However, it is instructive to look at it in some detail, since it is the simplest model, and it seems to have all the relevant features that go into the estimate of the weak phases. In terms of the different terms that should contribute to the hyperon decay amplitudes, it gives a contribution to the S -waves that is analogous to the commutator term when one uses PCAC. It also contributes to the P -waves through the K -pole diagram.

The different models one finds in the literature, attempt to improve upon these results by including several effects. The first thing they consider are additional diagrams. For the S -waves it is common to include a K^* pole, as well as baryon poles with the $(\frac{1}{2})^-, (70, 1^-)$ states. For the P -waves additional diagrams include the baryon poles with both the ground state octet and with the $(\frac{1}{2})^+, (56, 0^+)$ states.^[21] The second ingredient that goes into the different models, is the way to estimate the weak transitions in all the different diagrams. Apart from simple factorization with vacuum insertion, it is common to find bag model matrix elements^[22] or non-relativistic quark model matrix elements.^[23,24] In Fig. 3 we plot $A(\Lambda_-^0)$ as

a function of m_t using MIT bag model matrix elements as in the calculations of Donoghue *et al.*^[20]. It is not too different from the vacuum saturation results.[†] In vacuum saturation we find that ϕ_1^S is typically much larger than ϕ_1^P , because $-(m_\Lambda + m_p)g_1(m_\pi^2) + m_\pi^2 g_3(m_\pi^2)$ is small. All the signals are dominated by ϕ_1^S . On the other hand, in the MIT bag model^[23] $\phi_1^P > \phi_1^S$.

We refer the reader to the original papers for a detailed description of the models, and will simply comment on some relevant features. We have implemented CP violation into the model of Stech and Xu^[6] by letting $\xi = 0$ in our vacuum saturation result as they advocate. We have thus assumed that there is no additional CP violation in the di-quark contributions to amplitudes. Although this is consistent with their work, it is not necessarily true. It illustrates a further complication: the models have concentrated on trying to reproduce the magnitude of the amplitudes, to within a 20% (approximate) agreement with experiment. Since CP violating imaginary parts are negligible contributions to the magnitude of the amplitudes, it is possible to modify the models in many ways to include CP violation. Doing this, we obtain results that are indistinguishable (within errors) from vacuum saturation. In the model of Bonvin^[21], the amplitudes for $\Lambda \rightarrow p\pi^-$ are dominated by what he calls the "meson topology". That would mean that the phase for the S -wave is mostly given by CP violation in the $K^* - \pi$ transition, and that for the P -wave by CP violation in the $K - \pi$ transition. Since these phases are also unknown, there is further ambiguity in the model. Using different assumptions for these phases we get results that can be up to three times smaller or larger than our vacuum saturation curves.

Finally, we have looked at the results of Tadić and Trampetić^[24] where both: a) the MIT bag model and b) a non-relativistic quark model with a harmonic oscillator potential have been used to estimate the matrix elements. In both cases we find relatively large phases, but a large cancellation between them for the $\Lambda \rightarrow$

[†] A different model^[22] in which the CP violation is assumed to occur completely in the meson sector, finds signals remarkably similar to those found using the MIT bag model for Λ_c^0 decays (although this agreement can only be fortuitous).

$p\pi^-$ decay; this produces an $A(\Lambda_c^0)$ three to four times smaller than our vacuum saturation result. We found that the tendency of ϕ_1^S and ϕ_1^P to cancel is common to all the models we looked at for $\Lambda \rightarrow p\pi^-$. It reverses itself for $\Xi^- \rightarrow \Lambda\pi^-$, where the two phases tend to add.

In conclusion, if we use the different models to estimate our uncertainty in the matrix elements, we find that the results of Figs. 1, 2, 3 can increase or decrease by factors of two to three. We interpret this to mean that at present, the best we can do, is to give an order of magnitude estimate for the different signals.

For $\Xi^- \rightarrow \Lambda\pi^-$ we have also estimated the CP odd observable $\tilde{B}$ defined in Ref. 7:

$$\tilde{B} \simeq \frac{\pi}{8} \mathcal{P}_{\Xi}^0 \alpha_{\Lambda}^0 \rho_{\Xi}^0 \left(\frac{\mathcal{P}_{\Xi} - \tilde{\mathcal{P}}_{\Xi}}{\mathcal{P}_{\Xi} + \tilde{\mathcal{P}}_{\Xi}} + A(\Lambda_c^0) + B(\Xi^-) \right) \quad (14)$$

where $\mathcal{P}_{\Xi}$ is the polarization of the Ξ , and the superscript 0 indicates the averaged quantity for particle and antiparticle. From Eq. (6) we get $B_{\Xi} \simeq 12.2 A_{\Xi}$. For $m_t = 120 \text{ GeV}$, $\mu = 0.5 \text{ GeV}$ $\tilde{B}_{\Xi}$ is in the range $(0.3 - 3.1) \times 10^{-3}$. Using the experimental value for β_{Ξ} and assuming $\mathcal{P}_{\Xi} \simeq 1$, we find $\tilde{B} \simeq 7.6 \times 10^{-3} B_{\Xi}$. It has been argued that this is not easier to measure than $A(\Xi^-)$ in a machine like SUPERLEAR.^[24] However, a very crude estimate shows us that in a τ -charm factory one could need as few as 10^8 J/ψ decays to see $\tilde{B}(\Xi^-)$; and this could be fewer than those needed to see $A(\Lambda_c^0)$. Clearly, a detailed study of this issue, considering the different efficiencies for detecting the different particles should be done before reaching any conclusion.

We have found that when all uncertainties are taken into account, $A(\Lambda_c^0)$ is likely to be of the order of a few times 10^{-5} within the standard model, although it could be as large as 10^{-4} . As long as one does not have a framework to compute the hadronic matrix elements of the four-quark operators in a systematic way, it is not possible to estimate the errors in the results we have presented. For this same reason, it must be stressed that the efforts in studying hyperon decays should not be geared exclusively towards finding CP violation. In the process of collecting

a large sample of decays, an effort must be made to improve our knowledge of the amplitudes and strong scattering phases. This knowledge will be extremely important in understanding the origin of a CP violating signal when one is seen.

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FIGURE CAPTIONS

- 1) In this figure we show the results of vacuum saturation for $A(\Lambda_c^0)$ as a function of m_t , for a) $\mu = 0.5 \text{ GeV}$; and b) $\mu = 1.0 \text{ GeV}$. We have used the experimental constraints described in the text.
- 2) In this figure we show the results of vacuum saturation for $\Delta(\Lambda_c^0)$ as a function of m_t , for a) $\mu = 0.5 \text{ GeV}$; and b) $\mu = 1.0 \text{ GeV}$. We have used the experimental constraints described in the text.
- 3) Same as figure 1, but with the weak matrix elements calculated using the MIT bag model of Ref. 22.

Fig. 1a

Fig. 1b

Fig. 2a

Fig. 3a

Fig. 2b

Fig. 3b