fact it is *not* a unified gauge theory. This is because its tangent group is a product group $H = O(3,1) \otimes O(N)$. In analogy to the framework of grand unified theories, supergravity resembles the nonunified theory of $SU(3)^c \otimes SU(2) \otimes U(1)$ rather than unified theory such as SU(5) or O(10). In usual gauge theories, a product invariance group is the symptom that a spontaneous breaking has already occurred. As mentioned above, the vacuum state and tangent symmetry of supergravity can indeed be thought of as arising from the spontaneous breaking of the truly unified Riemannian space with tangent group $\mathcal{K} = OSp(3,1|4N)$.

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[(a)] On sabbatical leave from Northeastern University, Boston, Mass. 02115.

[1] P. Nath and R. Arnowitt, Phys. Lett. 65B, 73 (1976).

[2] The formulation of Ref. 1 used a supermetric rather than a supervielbein formulation. However, these two approaches are equivalent. See, e.g., R. Arnowitt and P. Nath, Phys. Lett. 78B, 581 (1978).

[3] D. Z. Freedman, P. Van Nieuwenhuizen, and S. Ferrara, Phys. Rev. D 13, 3214 (1976); S. Deser and B. Zumino, Phys. Lett. 62B, 35 (1976).

[4] J. Wess and B. Zumino, Phys. Lett. 66B, 361 (1977); R. Grimm, J. Wess, and B. Zumino, Phys. Lett. 73B, 15 (1978), and 74B, 51 (1978).

[5] L. Brink, M. Gell-Mann, P. Ramond, and J. H. Schwarz, Phys. Lett. 74B, 336 (1978), and 76B, 417 (1978), and Nucl. Phys. B145, 93 (1978).

[6] Arnowitt and Nath, Ref. 2.

[7] V. Ogievetsky and E. Sokatchev, Nucl. Phys. B124, (1977).

[8] W. Siegel, Nucl. Phys. B142, 301 (1978); W. Siegel and S. J. Gates, Nucl. Phys. B147, 77 (1979).

[9] J. G. Taylor, Phys. Lett. 78B, 577 (1978), and 79B, 399 (1978); Y. Ne'eman and T. Regge, Riv. Nuovo Cimento Ser. III 1, No. 5, 1 (1978); F. Mansouri, in *Proceedings of the Integrative Conference on Group Theory and Mathematical Physics, Austin, Texas, 11–16 September 1978*, edited by W. Deigelböck, A. Böhm, and E. Takasugi (Springer-Verlag, New York 1979).

[10] Proceedings of the Supergravity Workshop, Stony Brook, New York, 27 and 28 September, edited by P. Van Nieuwenhuizen and D. Z. Freedman (North-Holland, Amsterdam, to be published).

[11] P. Nath and R. Arnowitt, Northeastern University Report No. NUB 2405 (to be published).

[12] Global indices are denoted by Greek letters, $\Lambda \equiv (\mu, \alpha i)$, $\alpha = 1, \ldots, 4$, $i = 1, \ldots, N$, and local indices by Latin letters, $A \equiv (m, ai)$, $a = 1, \ldots, 4$. Middle lowercase letters are vectorial, early letter are (Majorana) spinorial, and $i, j, \ldots$ denote internal symmetry labels. η_{mn} is the Lorentz metric (with signature +2) and $\eta_{ab} = -(C^{-1})_{ab}$ is the Fermi metric (C = charge conjugation matrix). Comma, e.g., $\varphi_{,\Lambda}$ denotes right derivative.

[13] R. Arnowitt and P. Nath, unpublished; P. G. O. Freund, J. Math. Phys. (N.Y.) 17, 424 (1976).

[14] P. Nath and R. Arnowitt, Phys. Lett. 56B, 171 (1975).

[15] We conjecture that actually $\Delta_{mn\mu}(x)$ [and the $\Delta_{ij\mu}$ of Eq. (17)] vanishes for $N \geq 4$ also, since there are no local supersymmetry tensors independent of derivatives of the fields.

Nonspectator Quark Interactions and the Λ_c^+ Lifetime

V. Barger, J. P. Leveille, and P. M. Stevenson
Physics Department, University of Wisconsin, Madison, Wisconsin 53706
(Received 21 November 1979)

It is found that nonspectator interactions between c and d quarks in the Λ_c^+ give a substantial enhancement of its decay rate. The lifetime relation $\tau(\Lambda_c^+)/\tau(D^+) = [1 + \tau(D^+)/(3\times10^{-13}\ \mathrm{s})]^{-1}$ is derived, based on the assumption that the D^+ decays purely via charm quark decay. The semileptonic branching fractions are related by $B_l(\Lambda_c^+) = B_l(D^+) \times \tau(\Lambda_c^+)/\tau(D^+)$.

Inclusive weak decays of a hadron with a new flavor are usually assumed to proceed through the heavy quark, with the light quarks acting as spectators.[1,2] The lifetimes of charm states D^+ ($c\bar{d}$), F^+ ($c\bar{s}$), D^0 ($c\bar{u}$), and Λ_c^+ (cud) are thereby expected to be equal, and of order a few times 10^{-13} s. Recent lifetime measurements[3-6] give D^+ and F^+ values in qualitative accord with this, but may indicate D^0 and Λ_c^+ lifetimes that are significantly shorter.[3,4] The discrepancy suggests that interactions involving light quarks in the decaying hadron may play a role.[7] The subject of this Letter is the Λ_c^+ lifetime. We find that the nonspectator transition $cd \to su$ of Fig. 1(a) is at least comparable to the spectator process of Fig. 1(b), leading to a Λ_c^+ lifetime that is

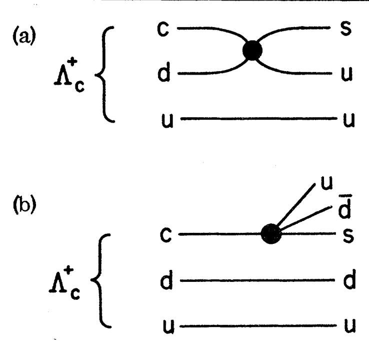

FIG. 1. (a) The nonspectator and (b) the spectator contributions to $\Lambda_c{}^+$ decay. The black circle represents the effective weak interactions [Eq. (1)].

shorter than the D^+ lifetime by a factor 1.5 to 3.

Before specializing to the $\Lambda_c{}^+$ case, we would like to comment briefly on the D^0 and F^+ cases. The D^0 and F^+ decay rates receive contributions from nonspectator processes $c\bar{u}\to s\bar{d}$ and $c\bar{s}\to u\bar{d}$, respectively. These contributions are small because of helicity suppression and are usually neglected. However, it has been suggested[8] that the emission of a single hard gluon by the initial quarks ($c\bar{u}\to c\bar{u}g\to s\bar{d}g$) can allow the nonspectator contribution to be very significant. This can occur for both D^0 and F^+, but is larger in the D^0 case because of different strong-enhancement and mass factors.[9] Thus nonspectator interactions can potentially explain the observed pattern of charmed-meson lifetimes, although it is not clear whether higher-order gluon corrections will preserve the above mechanism. In the $\Lambda_c{}^+$ case there is no helicity suppression of the process $cd\to su$, and there seems to be no clear reason for its neglect in previous inclusive calculations.

Nonspectator interactions have been considered previously in the case of Λ and Σ hyperon decays,[10,11] where they were found to be an important part of the explanation of the large nonleptonic decay rate and the $\Delta I=\frac{1}{2}$ rule. In these cases soft-pion techniques, or pole-dominance assumptions, have been used to estimate the diagrams. The much larger phase space involved in $\Lambda_c{}^+$ decays makes this approach inappropriate here. Nonspectator contributions were included in the analysis of two-body and quasi two-body decays of charmed baryons by Körner, Kramer, and Willrodt.[12] Arguing that the multihadron decay channels are resonance dominated, they estimated the nonleptonic $\Lambda_c{}^+$ decay rate by summing over all two-body channels. However, such an exclusive calculation of the inclusive rate requires detailed, untested assumptions about the symmetry of couplings and about wave-function overlaps.

Our approach is to calculate directly the nonspectator contribution to the $\Lambda_c{}^+$ rate in the free-quark model, including short-distance enhancement. We use a nonrelativistic approximation for the quarks in the $\Lambda_c{}^+$. The nonspectator contribution is proportional to the square of the modulus of the wave function for two quarks at the origin, $|\psi(0)|^2\equiv\langle\psi|\delta^3(\vec{r}_1-\vec{r}_2)|\psi\rangle$, and we estimate this from the $\Sigma_c{}^+$-$\Lambda_c{}^+$ mass difference. To obtain the inclusive decay rate we integrate over the phase space of the final-state quarks, on the assumption that the $\Lambda_c{}^+$ mass is large enough for this to represent (in the usual parton-model fashion) the sum of all hadronic final states.

We work with the conventional form of the effective weak-interaction Lagrangian[13] which, for the processes $cd\to su$ and $c\to su\bar{d}$, etc., can be written as

$$\mathcal{L}_{\rm eff}=(G_F/\sqrt{2})U_{ud}U_{cs}{}^*\tfrac{1}{2}(f_+\hat{O}^+ + f_-\hat{O}^-), \tag{1}$$

where U is the Kobayashi-Maskawa[14] mixing matrix, and the operators are $\hat{O}^\pm=[(\bar{u}d)(\bar{s}c)\pm(\bar{s}d)(\bar{u}c)]$, where $(\bar{q}Q)$ denotes a color singlet V-A current, $\bar{q}^\alpha\gamma_\mu(1-\gamma_5)Q_\alpha$. The coefficients f_+ and f_- (both equal to unity in the absence of strong interactions) are the usual short-distance enhancement factors[13]

$$f_-=[\alpha_s(m_c{}^2)/\alpha_s(m_W{}^2)]^\gamma\simeq 2.09,$$
$$f_+=(f_-)^{-1/2}\simeq 0.69, \tag{2}$$

where $\gamma=12/(33-2F)$, with F the effective number of flavors.

The calculation of the nonspectator contribution [Fig. 1(a)] is like a calculation of free-quark scattering, $cd\to su$ (with the c and d initially at rest), except that we must correctly fold in the color, spin, and spatial parts of the $\Lambda_c{}^+$ wave function. Because of the color antisymmetry of the wave function, the operator $\hat{O}^+$ does not contribute, since it is symmetric under interchange of color indices between the c and d quark field operators.[12,15] (This is because the exchange of color indices, followed by a Fierz rearrangement, turns the first term of $\hat{O}^+$ into the second, and vice versa.) This is an important fact because it means that the short-distance factor is $f_-{}^2$, which gives a substantial enhancement (~4 in the rate) to this contribution.

The spin-isospin wave function of the $\Lambda_c{}^+$ has the u,d quarks in an $I=0$, and hence a $J=0$, state. Consequently, the spin component of the c quark is just that of the Λ_c^+ itself. Also the spin component of the d quark is opposite to that of the spectator u quark. Hence the four spin states of the incoming c,d quarks are each distinct and noninterfering, so that the spin projection reduces to the usual sum-and-average procedure applied as if to free-quark scattering.

The spatial wave function enters in the following way.[16] To calculate the decay rate one evaluates the usual expression for a q-q scattering cross section, replacing the flux factor $1/|\vec{v}_1-\vec{v}_2|$ by $|\psi(0)|^2$. This calculation is straightforward and leads to

$$\Gamma_{cd\to su}=f_-^{\,2}|\psi(0)|^2(G_F^{\,2}/2\pi)|U_{ud}|^2|U_{cs}|^2[\Delta^2/(m_c+m_d)^2](\Delta^4-4m_s^{\,2}m_u^{\,2})^{1/2}, \tag{3}$$

where $\Delta^2\equiv(m_c+m_d)^2-m_s^{\,2}-m_u^{\,2}$.

We can estimate $|\psi(0)|^2$ from the Σ_c^+-Λ_c^+ mass splitting using the mass formula of De Rújula, Georgi, and Glashow (DGG)[17] where this mass difference is determined by the quantum chromodynamics analog of the Fermi-Breit "hyperfine" interaction. This gives (cf. Ref. 11)

$$\begin{aligned}|\psi(0)|^2&\equiv\langle\psi|\delta^3(\vec{r}_1-\vec{r}_2)|\psi\rangle\\&=\frac{9}{16\pi}\frac{1}{\alpha_s}\frac{m_u^{\,2}m_c}{(m_c-m_u)}(\Sigma_c^+-\Lambda_c^+).\end{aligned}\tag{4}$$

Using quark masses from DGG ($m_u=m_d=336$ MeV, $m_s=540$ MeV, and $m_c=1650$ MeV) and using[18] $\Sigma_c^+-\Lambda_c^+=170$ MeV, with $\alpha_s\simeq0.58$ {obtained from

$$\alpha_s(m_c^{\,2})=12\pi/[(33-2F)\ln(m_c^{\,2}/\Lambda^2)],$$

with a scale factor $\Lambda=0.5$ GeV and $F=3$ flavors}, we estimate

$$|\psi(0)|_{\Lambda_c}{}^2\simeq4.3\times10^{-3}/\alpha_s\simeq7.4\times10^{-3}\ \mathrm{GeV}^3. \tag{5}$$

Using this estimate in Eq. (3), with quark masses as above, and including the Cabibbo suppressed process $cd\to du$, we obtain

$$\Gamma_{\mathrm{nonsp}}\simeq21\times10^{-13}\ \mathrm{GeV}=0.32\times10^{13}\ \mathrm{s}^{-1}. \tag{6}$$

The estimate of $|\psi(0)|^2$ is the major source of uncertainty. Apart from this, the result is fairly independent of the light-quark masses (but scales roughly as the square of the assumed charm-quark mass).

The spectator contribution to the decay rate is well known,[1,2] but is quite sensitive to the u,d,s quark mass effects in the phase space, as well as depending on the fifth power of the c-quark mass. To avoid these uncertainties, we fix the spectator contribution in terms of the D^+ lifetime. This implicitly includes all gluon corrections to the free-quark decay for both semileptonic and nonleptonic modes. We can then write

$$\begin{aligned}\Gamma(\Lambda_c^+)&=\Gamma_{\mathrm{spec}}(\Lambda_c^+)+\Gamma_{\mathrm{nonsp}}(\Lambda_c^+)\\&=\Gamma(D^+)+\Gamma_{\mathrm{nonsp}}(\Lambda_c^+).\end{aligned}\tag{7}$$

This assumes that the conventional dogma[1,2] (charmed hadron decay = charmed quark decay) is at least correct for the D^+ (where the only possible nonspectator interactions are Cabibbo suppressed). Thus we obtain the lifetime relation

$$\begin{aligned}\tau(\Lambda_c^+)&=\frac{\tau(D^+)}{[1+\Gamma_{\mathrm{nonsp}}\tau(D^+)]}\\&=\frac{\tau(D^+)}{[1+\tau(D^+)/(3\times10^{-13}\ \mathrm{s})]}.\end{aligned}\tag{8}$$

Numerically this gives, in the range of possible values of $\tau(D^+)$,

$\tau(D^+)$ (s)	$\Gamma_{\mathrm{nonsp}}/\Gamma_{\mathrm{spec}}=\Gamma_{\mathrm{nonsp}}\tau(D^+)$	$\tau(\Lambda_c^+)$ (s)
2×10^{-13}	0.6	1.2×10^{-13}
4×10^{-13}	1.3	1.8×10^{-13}
6×10^{-13}	1.9	2.1×10^{-13}

The semileptonic branching fractions of Λ_c^+ and D^+ are related by the same factor:

$$B_l(\Lambda_c^+)=\frac{B_l(D^+)}{[1+\Gamma_{\mathrm{nonsp}}\tau(D^+)]}=B_l(D^+)\frac{\tau(\Lambda_c^+)}{\tau(D^+)}. \tag{9}$$

If $B_l(D^+)\simeq(23\pm6)\%$, as suggested by Kirkby,[4] then we expect $B_l(\Lambda_c^+)\simeq[(8-14)\pm3]\%$, depending on the D^+ lifetime.

We have not considered explicit gluon corrections to the nonspectator process.[19] There are difficulties of principle involved: One must decide *a priori* which of the virtual corrections should be absorbed into the wave function and which have already been included in the effective Lagrangian, Eq. (1). Thus a proper calculation of gluon radiative corrections would require a better understanding of the bound-state problem than we have at present. In any case one would want to improve upon the nonrelativistic approximation and the estimate of $|\psi(0)|^2$ (which again presupposes a better understanding of the bound-state problem) before one tried to calculate radiative corrections.

The nonspectator interaction discussed in this Letter may also be important for the decays of baryons containing still heavier quarks (b,t), but uncertainties in the wave functions and masses make reliable estimates impossible at this stage. We suspect, however, that this contribution may be less significant for heavier quarks. If $|\psi(0)|^2$ varies slowly with heavy-quark mass M_Q, then the nonspectator contribution scales as $M_Q^{\,2}$, whereas $\Gamma_{\text{spec}} \propto M_Q^{\,5}$. Then the conventional spectator description of heavy-baryon weak decays may be a better approximation than it seems to be for charmed baryons.

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[1]See, for example, M. K. Gaillard, in *Proceedings of the SLAC Summer Institute on Particle Physics, 1978*, edited by M. C. Zipf (Stanford Univ. Press, Stanford, Cal., 1978); J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. B100, 313 (1975); L. Maiani and T. Walsh, in *Proceedings of the International Symposium on Lepton and Photon Interactions at High Energies, Hamburg, 1977*, edited by F. Gutbrod (DESY, Hamburg, Germany, 1977); B. W. Lee, M. K. Gaillard, and J. Rosner, Rev. Mod. Phys. 47, 277 (1975).

[2]M. Suzuki, Nucl. Phys. B145, 420 (1978); N. Cabibbo and L. Maiani, Phys. Lett. 79B, 109 (1978).

[3]J. D. Prentice, in Proceedings of the Topical Workshop on the Production of New Particles in Super High Energy Collisions, University of Wisconsin, Madison, October, 1979 (to be published).

[4]J. Kirkby, in Proceedings of the International Symposium on Lepton and Photon Interaction at High Energies, Fermilab, Batavia, Illinois, 1979 (to be published).

[5]H. C. Ballagh *et al.*, University of Hawaii Report No. UH-411-351-79, 1979 (unpublished).

[6]C. Angelini *et al.*, Phys. Lett. 84B, 154 (1979).

[7]See, for example, S. P. Rosen, Phys. Rev. Lett. 44, 4 (1979).

[8]M. Bander, D. Silverman, and A. Soni, Phys. Rev. Lett. 44, 7 (1979); M. Suzuki, unpublished.

[9]The ratio of nonspectator contributions in $D°$ and F^+ decays, based on the model of Ref. 8, is

$$\frac{(f_+ + f_-)^2 |\psi(0)|_D^{\,2}\, m_D^{\,4}\, m_u^{\,-2}}{(f_+ - f_-)^2 |\psi(0)|_F^{\,2}\, m_F^{\,4}\, m_s^{\,-2}} .$$

[10]C. Schmid, Phys. Lett. 66B, 353 (1977); A. Le Yaouanc, L. Oliver, O. Pène, and J.-C. Raynal, Phys. Lett. 72B, 53 (1977).

[11]Riazuddin and Fayyazuddin, International Center for Theoretical Physics, Trieste, Report No. IC/79/24, 1979 (to be published).

[12]J. G. Körner, G. Kramer, and J. Willrodt, Phys. Lett. 78B, 492 (1978), and 81B, 419 (E) (1979), and Z. Phys. C 2, 117 (1979).

[13]B. W. Lee and M. K. Gaillard, Phys. Rev. Lett. 33, 108 (1974); G. Altarelli and L. Maiani, Phys. Lett. 52B, 351 (1974).

[14]M. Kobayashi and K. Maskawa, Prog. Theor. Phys. 49, 652 (1973).

[15]A. I. Vainshtein, V. I. Zakharov, and M. A. Shifman, Zh. Eksp. Teor. Fiz. 72, 1275 (1977) [Sov. Phys. JETP 45, 670 (1977)]; J. C. Pati and C. H. Woo, Phys. Rev. D 3, 2920 (1971).

[16]See, for example, J. M. Jauch and F. Rohrlich, *The Theory of Photons and Electrons* (Addison-Wesley, Reading, Mass., 1955); R. Van Royen and V. F. Weisskopf, Nuovo Cimento 50A, 617 (1967).

[17]A. De Rújula, H. Georgi, and S. L. Glashow, Phys. Rev. D 12, 147 (1975).

[18]C. Bricman *et al.*, Phys. Lett. 75B i, 1 (1978).

[19]The "penguin" contributions, considered for strange hadron decays, are not expected to be important for charm decays; see, for example, A. I. Vainshtein, V. I. Zakharov, and M. A. Shifman, Pis'ma Zh. Eksp. Teor. Fiz. 22, 123 (1975) [JETP Lett. 22, 55 (1975)].