[9]Generalization to more than one vector mesons is straightforward.

[10]We neglect the small effect of CP violation.

[11]Total cross sections are given, according to (2) and (4), by

$$\sigma(\nu_\mu e)=\frac{G_w^{\,2}}{16(4\pi)}E_{\rm tot}^{\,2}(g_V+g_A)^2(1+\tfrac{1}{3}b)$$
$$=\frac{G_w^{\,2}}{16(4\pi)}E_{\rm tot}^{\,2}[(g_V+g_A)^2+\tfrac{1}{3}(g_V-g_A)^2],$$
$$\sigma(\bar\nu_\mu e)=\frac{G_w^{\,2}}{16(4\pi)}E_{\rm tot}^{\,2}(g_V-g_A)^2\left(1+\frac{1}{3b}\right)$$
$$=\frac{G_w^{\,2}}{16(4\pi)}E_{\rm tot}^{\,2}[(g_V-g_A)^2+\tfrac{1}{3}(g_V+g_A)^2].$$

[12]G. 't Hooft, Phys. Lett. 37B, 195 (1970).

[13]We have neglected contributions from inelastic $\nu_{\mu L}e$ scattering with more than one neutrinos in the final state (e.g., $\nu_\mu e\to\nu_\mu e\nu\bar\nu$), since these cross sections are presumably smaller.

[14]The contamination of the incoming neutrino beam by right-handed neutrinos is (in principle) independent of the $\mathcal{S}$, $\mathcal{P}$, $\mathcal{T}$ terms discussed here, since the neutrino beam comes (mainly) from charged-current processes.

[15]A model which requires the existence of right-handed neutrinos is that of R. Mohapatra and J. Pati, Phys. Rev. D (to be published). See also J. C. Pati and A. Salam, Phys. Rev. D 10, 275 (1974).

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[17]The qualitative nature of the distribution can be easily seen by the argument of helicity conservation. For inclusive processes ($\nu p\to\nu+$ anything), the helicity arguments fail to predict even the qualitative nature of the distribution; thus detection of possible right-handed neutrinos in such processes is difficult.

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Gell-Mann–Low equation and on-mass-shell amplitudes*

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We show how, in certain theories, the Gell-Mann–Low renormalization-group equation can be applied to physical on-mass-shell amplitudes. These theories are characterized by a selection rule softening the zero-mass singularities in external mass channels. Physical amplitudes asymptotically reflect the anomalous dimensions of the theory. In particular, this analysis provides a field-theoretical argument for a connection between the asymptotic behavior of elastic amplitudes at fixed angle and the electromagnetic form factors. This connection is similar to that proposed by Wu and Yang.

The large-momentum behavior of Green's functions in renormalizable field theories has been extensively studied using the renormalization-group equations of Gell-Mann and Low[1] and the related Callan-Symanzik[2] equations.[3] However, most of these treatments involve amplitudes which are unphysical in the sense that the invariant masses of external particles become large and far off the "mass shell." In some cases this restriction can be avoided through use of light-cone expansions; in particular, the structure functions of inelastic electron-proton scattering have been studied in this manner.[4] In another approach, several authors have argued that in certain theories the mass insertion term of the Callan-Symanzik equation for the 3-point function can be neglected even if some external lines are kept on the mass shell.[5]

In this paper we argue that, under reasonable assumptions, the renormalization-group equations are directly applicable to the study of some on-shell amplitudes. This occurs in certain theories where a symmetry softens infrared singularities of an amplitude in external mass variables as the physical mass goes to zero. For example, in ϕ^4 theory the $\phi\to-\phi$ symmetry removes the singularity corresponding to two-particle intermediate states in any channel coupled to an odd number of external lines.[6] We will argue that the infrared singularities are also softened in theories of scalar or pseudoscalar mesons interacting with fermions, as well as in the artifical ϕ^3 theory in six dimensions.

In all of these theories, our arguments imply that the anomalous powers of the renormalization group solutions are measured by asymptotic on-shell amplitudes. As a consequence, various on-shell amplitudes can be related. In particular, we obtain a prediction similar to that of Wu and Yang[7]; i.e., differential cross sections at fixed angles are related in asymptotic behavior to electromagnetic form factors.

Although the analysis only applies in a certain class of theories, this class is remarkably wide. The only renormalizable theories in four dimensions for which we do not draw conclusions either involve vector mesons, i.e., quantum electrodynamics, or involve scalar mesons with a ϕ^3 coupling.

Our analysis treats the scattering particles as elementary in the sense that corresponding to each particle is an elementary field in an underlying local Lagrangian. This approach is complementary to treating hadrons as bound states of more elementary constituents, such as quarks. These two approaches may not be exclusive; indeed, we obtain some results similar to those of a composite picture.

To establish notation and conventions, we review the derivation of the renormalization-group equations. To be specific, consider the theory obtained from the Lagrangian density

$$\begin{aligned}\mathcal{L}(x) &= \tfrac{1}{2}\partial_\mu\phi(x)\partial^\mu\phi(x) - \tfrac{1}{2}m_0{}^2\phi^2(x)\\ &\quad - (g_0/4!)\phi^4(x)\,, \end{aligned} \tag{1}$$

where $\phi(x)$ is a Hermitian field, and the parameters g_0 and m_0 are the bare coupling constant and mass. The renormalized vertex functions[8] Γ_n can be expressed as functions of the physical mass m^2, the renormalized coupling constant g, a renormalization point μ^2, and the n external momenta p_i entering the vertex. The Γ_n are normalized so that

$$\Gamma_2(p^2; m^2, g, \mu^2)|_{p^2=m^2} = 0\,, \tag{2}$$

$$\Gamma_2(p^2; m^2, g, \mu^2)|_{p^2=\mu^2} = i(\mu^2 - m^2)\,, \tag{3}$$

$$\Gamma_4(s, t, p_i{}^2; m^2, g, \mu^2)|_{s=t=(4/3)\mu^2,\ p_i{}^2=\mu^2} = -ig\,, \tag{4}$$

where

$$s = (p_1 + p_2)^2, \quad t = (p_1 - p_3)^2\,.$$

Although we express the theory in terms of the three parameters m^2, g, and μ^2, the theory has an underlying dependence on only two parameters, i.e., $m_0{}^2$ and g_0. A change in μ^2 can be compensated by a change in coupling constant and field normalization. This is manifested in the renormalization-group equation[3]

$$\left[\mu^2\frac{\partial}{\partial\mu^2} + \beta\left(g, \frac{m^2}{\mu^2}\right)\frac{\partial}{\partial g} + \frac{n}{2}\gamma\left(g, \frac{m^2}{\mu^2}\right)\right]\Gamma_n(p_i; m^2, g, \mu^2) = 0\,. \tag{5}$$

Here β and γ are dimensionless functions determined by the theory. The integrated form of Eq. (5) is

$$\begin{aligned}\Gamma_n(p_i; m^2, g, \mu^2) &= \Gamma_n(p_i; m^2, \hat{g}(\lambda), \lambda^2\mu^2)\\ &\quad\times\exp\left[-\frac{n}{2}\int_1^{\lambda^2}\frac{d\lambda'^2}{\lambda'^2}\gamma\left(\hat{g}(\lambda'), \frac{m^2}{\lambda'^2\mu^2}\right)\right]\,, \end{aligned}\tag{6}$$

where λ is an arbitrary real number and $\hat{g}(\lambda)$ is determined by the differential equation

$$\lambda^2\frac{d\hat{g}(\lambda)}{d\lambda^2} = \beta\left(\hat{g}(\lambda), \frac{m^2}{\lambda^2\mu^2}\right)\,, \tag{7}$$

with the constraint $\hat{g}(1) = g$. To obtain information on asymptotic behavior from these equations one uses dimensional analysis, which implies

$$\Gamma_n(p_i; m^2, g, \lambda^2\mu^2) = \lambda^{4-n}\Gamma_n\left(\frac{p_i}{\lambda}; \frac{m^2}{\lambda^2}, g, \mu^2\right)\,. \tag{8}$$

Conventional applications of the renormalization group combine Eqs. (8) and (6) to give

$$\begin{aligned}\Gamma_n(\lambda p_i; m^2, g, \mu^2) &= \lambda^{4-n}\Gamma_n\left(p_i; \frac{m^2}{\lambda^2}, \hat{g}(\lambda), \mu^2\right)\\ &\quad\times\exp\left[-\frac{n}{2}\int_1^{\lambda^2}\frac{d\lambda'^2}{\lambda'^2}\gamma\left(\hat{g}(\lambda'), \frac{m^2}{\lambda'^2\mu^2}\right)\right]\,. \end{aligned}\tag{9}$$

We emphasize that this is an exact consequence of the theory. If $\hat{g}(\lambda)$ has a finite limit of the form

$$\lim_{\lambda\to\infty}\hat{g}(\lambda) = g_\infty\,, \tag{10}$$

then $\beta(x, m^2/\lambda^2\mu^2)$ must have a zero at a value of x which approaches g_∞ as λ goes to infinity. This connection between zeros of β and limiting values of $\hat{g}(\lambda)$ has been extensively discussed before.[3] For convenience of discussion, we assume the existence of an appropriate zero and the behavior expressed in Eq. (10). Consequently, by taking λ large we have[9]

$$\begin{aligned}\Gamma_n(\lambda p_i; m^2, g, \mu^2) &\underset{\lambda\to\infty}{\sim} \lambda^{4-n-n\gamma(g_\infty, 0)}\\ &\quad\times\Gamma_n\left(p_i; \frac{m^2}{\lambda^2}, \hat{g}(\lambda), \mu^2\right)\,. \end{aligned}\tag{11}$$

If $\Gamma_n(p_i\,;m^2/\lambda^2, \bar{g}(\lambda), \mu^2)$ approaches a finite nonvanishing quantity as λ goes to infinity, then we know the asymptotic behavior of the Green's functions

$$\Gamma_n(\lambda p_i\,;m^2,g,\mu^2) \underset{\lambda\to\infty}{\sim} \lambda^{4-n-n\gamma(g_\infty,0)}\ . \tag{12}$$

It has been argued that the massless version of the theory normalized at μ^2 exists.[3] Thus Eq. (12) should be valid if the p_i are selected to avoid singularities of the massless theory. We shall henceforth consider this conventional application of the renormalization group equation as valid.

We are now ready to discuss asymptotic behavior when some external masses are held fixed. For definiteness we study the four-point function $\Gamma_4(s,t,p_i{}^2;m^2,g,\mu^2)$, where s and t were defined previously. We write the renormalization-group equation in the form

$$\Gamma_4(\lambda^2 s,\lambda^2 t,p_i{}^2;m^2,g,\mu^2)=\Gamma_4\left(s,t,\frac{p_i{}^2}{\lambda^2}\,;\frac{m^2}{\lambda^2},\bar{g}(\lambda),\mu^2\right)\exp\left[-4\int_1^{\lambda^2}\frac{d\lambda'^2}{\lambda'^2}\gamma\left(g(\lambda'),\frac{m^2}{\lambda'^2\mu^2}\right)\right]$$
$$\underset{\lambda\to\infty}{\sim}\lambda^{-4\gamma(g_\infty,0)}\Gamma_4\left(s,t,\frac{p_i{}^2}{\lambda^2}\,;\frac{m^2}{\lambda^2},g_\infty,\mu^2\right). \tag{13}$$

To obtain information on the behavior of the left-hand side of this equation as λ becomes large, we need to know how the Green's functions behave as both the physical mass and external invariant masses go to zero. In the massless theory, singularities are expected as any invariant upon which the Green's functions depend is taken to zero. Consequently, we must investigate the theory at these "infrared" singular points. We will argue that under certain circumstances, as for external masses in ϕ^4 theory, these singularities are sufficiently mild that the Green's functions will remain finite at the infrared points. With this argument, Eq. (13) yields

$$\Gamma_4(\lambda^2 s,\lambda^2 t,p_i{}^2;m^2,g,\mu^2)\underset{\lambda\to\infty}{\sim}\lambda^{-4\gamma(g_\infty,0)}\ , \tag{14}$$

and, in general, for $n\geq 4$ we have

$$\Gamma_n(\lambda^2 q_j{}^2,p_i{}^2;m^2,g,\mu^2)\underset{\lambda\to\infty}{\sim}\lambda^{4-n-n\gamma(g_\infty,0)}\ , \tag{15}$$

where the $p_i{}^2$ are the squared masses of the external legs and the $q_j{}^2$ are the remaining independent Lorentz invariants upon which Γ_n depends.

Our argument is based on phase space. Suppressing all dependences other than on the physical mass and one external mass, we study the behavior of an amplitude $\Gamma(p^2,m^2)$ on the variable p^2 by dispersing in this variable

$$\Gamma(p^2,m^2)=\frac{1}{\pi}\int_{\sigma_0}^{\infty}\frac{A(\sigma,m^2)}{\sigma-p^2-i\epsilon}\,d\sigma\ , \tag{16}$$

where σ_0 is the threshold of the lowest intermediate states. Possible subtractions are irrelevant to our discussion. The absorptive part $A(\sigma,m^2)$ is obtained from a sum over intermediate states

$$A(\sigma,m^2)=\sum_n T^*_{pn}T_{fn}(2\pi)^4\delta^4(p_n-p)\ , \tag{17}$$

where T_{pn} is the amplitude for the external leg p to produce the intermediate state n and T_{fn} is the amplitude for the remaining external lines to produce the same intermediate state.[10]

Temporarily assume that both T_{pn} and T_{fn} are finite in the infrared limit $p^2, m^2\to 0$. The threshold behavior of an n particle intermediate state in $A(p^2,m^2)$ is then given by n-particle phase space:

$$\begin{aligned}
\mathcal{P}_n(p,m^2)&=\int(2\pi)^4\delta^4(\Sigma q_i-p)\\
&\quad\times\prod_{i=1}^n\left(\frac{d^4q_i}{(2\pi)^4}\,2\pi\delta(q_i{}^2-m^2)\theta(q_{i_0})\right)\ ,\\
\mathcal{P}_2(p,m^2)&=\frac{1}{4\pi}\,\frac{(p^2-4m^2)^{1/2}}{(p^2)^{1/2}}\,\theta(p^2)\\
&\underset{m^2\to 0}{\longrightarrow}\frac{1}{4\pi}\,\theta(p^2)\theta(p_0)\ ,\\
\mathcal{P}_3(p,m^2)&\underset{m^2\to 0}{\longrightarrow}\frac{1}{2^8\pi^3}\,p^2\theta(p^2)\theta(p_0)\ ,\\
\mathcal{P}_n(p,m^2)&\underset{m^2\to 0}{\sim}(p^2)^{n-2}\theta(p^2)\theta(p_0)\ .
\end{aligned} \tag{18}$$

With the exception of $\mathcal{P}_2$, any of these threshold behaviors inserted into Eq. (16) will give a finite result for $\Gamma(p^2/\lambda^2, m^2/\lambda^2)$ as λ goes to infinity. The crucial point is that because of the $\phi\to-\phi$ symmetry of ϕ^4 theory, the troublesome two-particle intermediate state does not occur in external mass channels.

We must still discuss the assumption that T_{pn} and T_{fn} are finite at $p^2=m^2=0$. Two types of problem can arise here. The first comes from further singularities in the p^2 channel. An inductive argument in perturbation theory indicates that these singularities give no trouble: to lowest order the T's remain finite, while singularities in higher orders are related to singularities in lower orders using the above phase-space argument. The second difficulty in discussing T_{pn} and T_{fn} for $p^2=m^2=0$ involves particle exchanges in other than the p^2 channel.[10] Since we are taking the mass to zero, singularities in other channels of T_{fn} approach the boundary of the physical region for the intermediate state n. Generally the integration over the phase space of the intermediate particles will smooth these singularities to give in A at worst a logarithmic factor of the mass.[11] Thus we expect in perturbation theory

$$\Gamma\left(\frac{p^2}{\lambda^2}, \frac{m^2}{\lambda^2}\right)=\text{finite} + o(\lambda^{\epsilon-1}) , \qquad (19)$$

where ϵ is some arbitrary small positive power. To obtain Eq. (15) we must now assume that after summing to all orders in perturbation theory, Eq. (19) is still valid or, at worst, ϵ can still be kept less than unity. This is consistent with considering the renormalization group as a technique for summing the logarithms occurring in perturbation theory, and ignoring terms down from the leading ones by powers of the large invariants. We must emphasize that this assumption is a cornerstone of our argument.

It should now be clear why our technique does not allow the momentum transfer t to be kept fixed in the asymptotic study of $\Gamma_4(s, t, p_i{}^2, m^2)$. In the t channel two-particle intermediate states will occur. Because two-particle phase space for massless particles is constant at threshold, we expect perturbation theory to give powers of $\log t$ as t goes to zero. Thus, we cannot take t to zero on the right-hand side of Eq. (13). In particular, our method does not apply to deeply inelastic electron-proton scattering.

In general, our arguments apply when there is some condition reducing the phase space available to intermediate states in the respective channel. Such is the case in a theory of fermions interacting with spin-zero mesons through either a scalar or pseudoscalar coupling. For example, considering intermediate states containing a single fermion and using the normalization condition for the spinors $\overline{U}(p)U(p)=1$, the $\mathcal{P}_n$'s in Eq. (18) acquire an additional factor of the fermion mass m_f. Using the fact that in the zero-mass limit the intermediate particles are collinear with the initial particle, one can show that T_{pn} of Eq. (17) remains finite for scalar and γ_5 coupling; in contrast, T_{pn} behaves as $m_f{}^{-1}$ in a theory of vector gluons interacting with fermions through a γ_μ coupling. With the scalar and pseudoscalar couplings, the factor m_f in the phase space removes the infrared divergence. An alternative argument notes that the quantities $1\pm\gamma_5$ commute with the interaction; consequently, as the fermion mass goes to zero its helicity must flip in an interaction. Angular momentum conservation then provides a vanishing factor as the masses go to zero.[12] Finally, the technique also applies to the artificial ϕ^3 theory in six-dimensional space-time, because the extra dimensions provide the needed extra phase-space factor.

In all these theories, we have looked at low-order Feynman graphs that contain the most dangerous infrared singularities discussed above. They confirm the previous conclusions.

Our methods can be extended to electromagnetic form factors. Normalizing so that the electric charges of the particles are μ^2-independent, the electromagnetic current does not carry an anomalous dimension. In theories of the type discussed above where external legs can be kept on-shell, a dimensionless electromagnetic form factor displays the anomalous dimensions of the on-shell legs[9]

$$F(t) \underset{t\to\infty}{\sim} t^{-\gamma(g_\infty,0)} . \qquad (20)$$

Combining this with Eq. (14), we obtain

$$\Gamma_4(s, t) \underset{\substack{t\to\infty \\ s/t \text{ fixed}}}{\sim} F(t)^2 f(\cos\theta) , \qquad (21)$$

where $f(\cos\theta)$ is an unknown function of the scattering angle. In terms of the elastic differential cross section,

$$\frac{d\sigma}{dt} \underset{\substack{t\to\infty \\ s/t \text{ fixed}}}{\sim} \frac{1}{t^2} F_1(t)^2 F_2(t)^2 f(\cos\theta) , \qquad (22)$$

where the indices on the F's refer to the respective particles involved in the reaction. If we assume an asymptotic form factor behavior of t^{-2} for baryons and t^{-1} for mesons, then for baryon-baryon elastic scattering at fixed angle $d\sigma/dt \sim t^{-10}$ and for meson-baryon scattering $d\sigma/dt \sim t^{-8}$. These behaviors have been conjectured in composite models of hadrons.[13]

In summary, we have considered application of the renormalization group equations to amplitudes on the mass shell. Our arguments apply in field theories possessing some selection rule softening infrared singularities encountered as the physical mass and the external masses are taken to zero. In theories of this type, the asymptotic behavior of on-shell amplitudes with all other invariants large provides a direct measure of the anomalous dimensions in the field theory. From this general result, we obtain a behavior similar to that conjectured by Wu and Yang[7] connecting electromagnetic form factors with fixed-angle elastic scattering.

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[8]We consider the amputated one-particle irreducible parts of connected Green's functions for $n>2$, and Γ_2 is the negative of the inverse propagator. A similar renormalization-group equation applies to the full connected Green's functions.

[9]If $\gamma(g_\infty, 0)$ vanishes, under some circumstances factors of $\log\lambda$ can enter the right-hand side of Eq. (11). For simplicity we do not explicitly exhibit such factors, although our arguments do not depend on their absence.

[10]Note that although Γ is one-particle irreducible, T_{fn} need not be.

[11]An exception to this occurs when T_{fn} is an elastic on-shell amplitude in a theory where a single particle can be exchanged. One consequence of this is that all legs of a three-point function cannot simultaneously be kept on-shell in our argument. Indeed, our conclusions must fail in this case, because when all legs are on-shell, the three-point function has nothing else on which to depend.

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Radiative corrections to the nonleptonic Ξ decays and the $\Delta I = \frac{1}{2}$ selection rule

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The effects of radiative corrections on the $\Delta I = 1/2$ rule predictions for the nonleptonic Ξ decays are studied. The radiative corrections are estimated using standard perturbation theory, assuming that all particles are structureless. Two different Lagrangian models are considered and it is found that the radiative corrections are sensitive to the choice of the Lagrangian. Most of the results have the usual divergence difficulties which are overcome with the aid of a cutoff. The results for the branching ratio $\Gamma(\Xi^- \to \Lambda^0\pi^-)/\Gamma(\Xi^0 \to \Lambda^0\pi^0)$ give a corrected value in disagreement with the latest experimental data by as much as 5.3 standard deviations. The corrections to the asymmetry parameters lead to good agreement between $\Delta I = \frac{1}{2}$ and experiment.

I. INTRODUCTION

One of the more intriguing regularities in the weak, nonleptonic decays of strange particles is that they experimentally obey the isospin selection rule $\Delta I = \frac{1}{2}$ to a surprising accuracy. This selection rule, proposed several years ago,[1] requires that the weak-interaction Lagrangian responsible for these decays transforms as a spinor under the isospin group SU(2). Over the past years this selection rule has been tested for both K-meson and hyperon decays with results suggesting its general validity. However, with recent experiments[2] furnishing very precise measurements of these decays, it has become increasingly important to take into account electromagnetic corrections to the $\Delta I = \frac{1}{2}$ rule if one wishes to test the limits of its validity to the order of one percent. Such corrections have been estimated[3] for K^0, Λ^0, and $\Sigma^\pm$ decays and the results tend to suggest that $\Delta I = \frac{1}{2}$ may be violated as the result of a small admixture (~4%) of $\Delta I = \frac{3}{2}$ in the decay amplitudes.

In this paper we consider the radiative corrections to the $\Delta I = \frac{1}{2}$ rule for nonleptonic Ξ decay. Our calculations are made using standard perturbation theory and treating all particles as point particles, thus neglecting all effects due to strong interactions.[4] This feature of the calcula-