Non-Reggeization of the Vector Meson*

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Nucleon-antinucleon scattering is treated in a field theory of interacting vector mesons and spin-$\frac{1}{2}$ or spin-0 "nucleons." Some involved cancellations in the asymptotic sixth-order amplitudes are demonstrated and disagreements in earlier papers are resolved. It is found that the vector-meson singularity in the partial-wave amplitudes is unaffected and that therefore the vector meson does not lie on a Regge trajectory to this order.

INTRODUCTION

It was demonstrated by Gell-Mann and others in a series of papers[1–4] that the effect of radiative corrections on the Compton scattering of vector mesons and nucleons is to absorb the Kronecker delta singularity in the partial-wave amplitudes which corresponds to the nucleon, and to place the nucleon on a Regge trajectory. There has been some interest in treating the vector-meson channel in the same field theory[5–7] to see whether it too lies on such a trajectory. The correct solution of this problem is important since it affects the Chew-Frautschi-Mandelstam postulate[7,8] that the absence of Kronecker delta terms in all channels is a criterion for a "bootstrapped" theory.

The purpose of this paper is to resolve the varying results that appear in the literature.[5–7,9] We will show that the vector meson remains un-Reggeized to sixth order in nucleon-antinucleon scattering (i.e., does not lie on a Regge trajectory[1–4]), both for spin-$\frac{1}{2}$ and spin-0 "nucleons."

SPINLESS NUCLEONS

We begin by treating spinless nucleons coupled to massive "photons" (vector mesons), as in Refs. 5, 6, and 9. The photon has $C=-$; since $CP=+$ for the nucleon-antinucleon state, only amplitudes of odd parity can contribute to the photon trajectory. Thus, if the scattering takes place in the s channel, we expect the scattering amplitude to have the asymptotic form

$$f \xrightarrow[t\to\infty]{} t^{\alpha(s)}-u^{\alpha(s)} \tag{1}$$

apart from multiplicative factors.

The scattering amplitude in second order is asymptotically

$$f \xrightarrow[t\to\infty]{} (g^2/8\pi W)(s-\lambda^2)^{-1}(u-t), \tag{2}$$

(λ=photon mass, $W^2=s$) which might turn out to be the lowest order term in perturbation theory for a trajectory

$$\alpha(s)=1+\gamma(s), \tag{3}$$

where

$$\lim_{g^2\to 0}\gamma(s)=0. \tag{4}$$

Explicitly, if the amplitude were dominated by such a trajectory, we would expect the leading terms in each order of perturbation theory to have the form

$$t^{\alpha(s)}-u^{\alpha(s)}\approx t\{1+\gamma(s)\ln t+\tfrac{1}{2}[\gamma(s)\ln t]^2+\cdots\} \\ -u\{1+\gamma(s)\ln u+\tfrac{1}{2}[\gamma(s)\ln u]^2+\cdots\}. \tag{5}$$

The photon would lie on the trajectory if $\gamma(\lambda^2)=0$. Since the lowest order term $t-u$ is provided by the Born approximation, a necessary condition for the vector meson to be Reggeized is that the form $t\ln t-u\ln u$ dominate in some higher order of perturbation theory.

It is well known that the fourth-order diagrams do not provide such an asymptotic form. The only diagrams giving contributions as large as $t\ln t$ are those of Fig. 1. However, their total contribution has the form

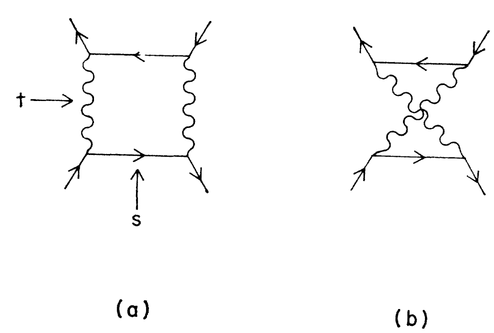

FIG. 1. Fourth-order terms dominant at large t.

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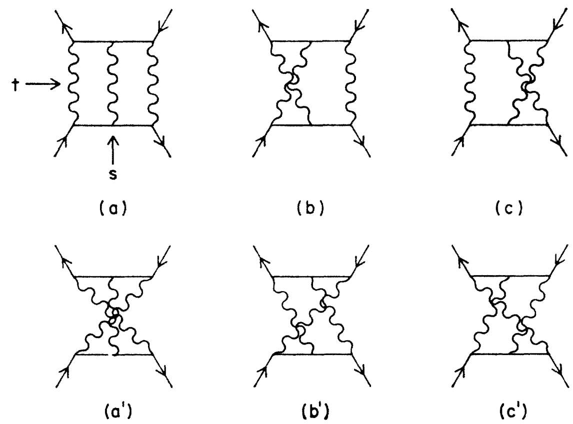

FIG. 2. Sixth-order terms dominant at large t.

$t\ln t+u\ln u$, which has the wrong parity (it should be obvious anyway that the two-photon intermediate state has $C=+$ and could not contribute to a trajectory with $C=-$). The lowest order in which the form $t\ln t-u\ln u$ might appear is therefore sixth order.

The only diagrams in sixth order which contribute terms of order $t\ln t$ are those depicted in Fig. 2, as pointed out by Freund and Oehme.[5] We will show that although each of these diagrams contributes a leading term proportional to $t\ln t$, a number of cancellations occur which make the sum of lower asymptotic order.

The four-dimensional integrals for each of these sixth-order diagrams contain seven propagator factors in their denominators. We will combine these in a particular way that simplifies subsequent work. Using the usual Feynman parameterization technique, we first combine the propagators arising from the initial nucleon line into a single factor. Next we combine the two factors from the final nucleon line into a single term. Finally, we combine the five factors which now remain by yet another parameterization.

For spinless nucleons, the only relevant term in the numerator in the large-t limit turns out to be $(-8t^3)$ and we have simplified accordingly in what follows.

For diagram (a) of Fig. 2, the scattering amplitude thus obtained is effectively (all masses have been set equal to unity in the rest of this section)

$$
\begin{aligned}
f_a=&(1/8\pi W)(-g^6/16\pi^4)t^3\\
&\times\int_0^1 dx_1dx_2dy_1dy_2dz_1\cdots dz_5\delta(x_1+x_2-1)\\
&\times\delta(y_1+y_2-1)\delta(z_1+\cdots+z_5-1)z_4z_5CD_a^{-3}. \quad (6)
\end{aligned}
$$

The method developed by Polkinghorne[10] for extracting the asymptotic limit for planar graphs may be applied to this amplitude. Only the region where z_4 and z_5 are small contributes, so that the form of the amplitude for large t will be unchanged if we simplify C and D_a to the forms

$$C\to C'=z_1z_2+z_2z_3+z_3z_1, \quad (7)$$

$$D_a\to D_a'=d+z_4z_5(z_2+z_3x_1y_1+z_1x_2y_2)t, \quad (8)$$

where

$$d=z_1z_2z_3s-C'. \quad (9)$$

For large t, we find

$$f_a \underset{t\to\infty}{\longrightarrow} g(s)t\ln t, \quad (10)$$

where

$$
\begin{aligned}
g(s)=&(1/8\pi W)(-g^6/32\pi^4)\\
&\times\int_0^1 dx_1dx_2dy_1dy_2dz_1dz_2dz_3\\
&\times\delta(x_1+x_2-1)\delta(y_1+y_2-1)\delta(z_1+z_2+z_3-1)\\
&\times C'd^{-1}(z_2+z_3x_1y_1+z_1x_2y_2)^{-2}\\
=&(1/8\pi W)(-g^6/32\pi^4)\int_0^1 dz_1dz_2dz_3\delta(z_1+z_2+z_3-1)\\
&\times d^{-1}\ln[(z_1+z_2)(z_2+z_3)/z_2^2]. \quad (11)
\end{aligned}
$$

This result was obtained by Freund and Oehme.[5]

The asymptotic form of the amplitude for diagram (a′) of Fig. 2 is obtained by replacing t by u and affixing an over-all negative sign [this follows from the crossing symmetry which relates diagrams 2(a) and 2(a′), and can be seen by inspecting the four-dimensional integrals for the two amplitudes]:

$$f_{a'} \underset{t\to\infty}{\longrightarrow} -g(s)u\ln u. \quad (12)$$

There seems to be no disagreement in the literature about these amplitudes for diagrams (a) and (a′) of Fig. 2.

The amplitudes for the nonplanar diagrams 2(b) and 2(c) are more complicated, because the coefficient of t in the denominator can vanish within the region of integration as well as on its boundary. The regions of z_4 and z_5 small are still important, but we will have to be more careful about extracting the large-t form. We have

$$
\begin{aligned}
f_b=f_c=&(1/8\pi W)(-g^6/16\pi^4)t^3\\
&\times\int_0^1 dx_1dx_2dy_1dy_2dz_1\cdots dz_5\\
&\times\delta(x_1+x_2-1)\delta(y_1+y_2-1)\\
&\times\delta(z_1+\cdots+z_5-1)z_4z_5CD_b^{-3}, \quad (13)
\end{aligned}
$$

[10] J. C. Polkinghorne, J. Math. Phys. 4, 503 (1963).

where asymptotically we may write

$$C \to C', \tag{14}$$

$$D_b \to D_b' = d + z_4 z_5 [(z_2 + x_2 z_1) - (z_2 + x_1 z_3) y_1 t]. \tag{15}$$

We may not directly apply Polkinghorne's method for planar graphs to find the asymptotic form of this integral. If this is done, an expression is obtained containing a double pole on the path of integration [this occurs for the term f_2 in Eq. (14) of Ref. 5]. This double pole arises from two poles having pinched the contour of integration in the large-t limit; the correct path of integration actually passes between these poles. It is necessary to delay taking the large-t limit until the contour of integration has been pulled across one of these poles. There arises an extra "pinch" contribution from the residue of the integrand at the pole, in addition to the ordinary term one obtains by simply avoiding the double pole while doing the integral.

The method of obtaining the pinch contribution is discussed in Polkinghorne's later paper[11]; it involves doing a number of the integrations exactly before extracting the large-t form. We perform the x and y integrations exactly in Eq. (13) (using the still valid replacement $C \to C'$, $D_b \to D_b'$) to obtain

$$\begin{aligned} f_b = f_c \approx\ & (1/8\pi W)(-g^6/16\pi^4) t^3 \\ & \times \int_0^1 dz_1 \cdots dz_5 \delta(z_1 + \cdots + z_5 - 1) \\ & \times C'(z_3/2t)(z_3 d + z_4 z_5 t C')^{-2} \\ & \times \{R + \ln[(d + z_4 z_5 t z_2)(d + z_4 z_5 t z_1) / \\ & \qquad (d - z_4 z_5 t z_3)(d + z_4 z_5 t (z_1 + z_2))]\}. \end{aligned} \tag{16}$$

R consists of rational terms which are not affected when the contour is pulled across one of the pinching poles.

If $t \to +\infty + i\epsilon$, the logarithmic term has the form

$$\ln[(+\infty + i\epsilon)(+\infty + i\epsilon)/(-\infty - i\epsilon)(+\infty + i\epsilon)]. \tag{17}$$

Below threshold, d is negative, and in analytically continuing a term of the form $\ln[d + \alpha(+\infty + i\epsilon)]$ from positive to negative α, one passes counterclockwise around the origin. Thus we have

$$\ln(-\infty - i\epsilon) = \ln(+\infty + i\epsilon) + i\pi \tag{18}$$

in this problem, and the correct contribution from the logarithmic term of Eq. (17) in the limit $t \to +\infty + i\epsilon$ is just $-i\pi$.

This quantity represents the additional contribution obtained in pulling the contour of integration across one of the pinching poles, as Polkinghorne showed.[11] The extra contribution to the amplitudes thus obtained is

$$\begin{aligned} \delta f_b = \delta f_c =\ & (1/8\pi W)(-g^6/16\pi^4) t^3 \\ & \times \int_0^1 dz_1 \cdots dz_5 \delta(z_1 + \cdots + z_5 - 1) \\ & \times C'(z_3/2t)(z_3 d + z_4 z_5 t C')^{-2}(-i\pi) \\ & \underset{t\to\infty}{\longrightarrow} (1/8\pi W)(-g^6/16\pi^4) t \ln t \\ & \times \int_0^1 dz_1 dz_2 dz_3 \delta(z_1 + z_2 + z_3 - 1)(-i\pi/2d). \end{aligned} \tag{19}$$

To this must be added the result obtained from the rational terms R in Eq. (16), which may be correctly obtained by simply avoiding the double pole referred to previously in the large-t limit. The net result for diagrams 2b and 2c is

$$f_b + f_c \underset{t\to+\infty+i\epsilon}{\longrightarrow} [g'(s) - 2\pi i h(s)] t \ln t. \tag{20}$$

In this expression, $h(s)$ is the pinch contribution found above:

$$h(s) = (1/8\pi W)(-g^6/32\pi^4) \times \int_0^1 dz_1 dz_2 dz_3 \delta(z_1 + z_2 + z_3 - 1) d^{-1}. \tag{21}$$

The function $g'(s)$ is the same form as obtained in Ref. 5, with the prescription that the double pole is to be displaced from the path of integration when actually doing the integrals:

$$\begin{aligned} g'(s) =\ & (1/8\pi W)(-g^6/32\pi^4) \\ & \times \int_0^1 dx_1 dx_2 dy_1 dy_2 dz_1 dz_2 dz_3 \\ & \times \delta(x_1 + x_2 - 1)\delta(y_1 + y_2 - 1)\delta(z_1 + z_2 + z_3 - 1) \\ & \times 2C' d^{-1}(z_1 x_2 + z_2 y_2 - z_3 x_1 y_1)^{-2} \\ =\ & (1/8\pi W)(-g^6/32\pi^4) \int_0^1 dz_1 dz_2 dz_3 \delta(z_1 + z_2 + z_3 - 1) \\ & \times d^{-1} 2 \ln[z_1 z_2 / (-z_3)(z_1 + z_2)]. \end{aligned} \tag{22}$$

However, a simple change of variables shows that $g'(s) = -g(s)$. Thus we have simply

$$f_b + f_c \underset{t\to+\infty+i\epsilon}{\longrightarrow} [-g(s) - 2\pi i h(s)] t \ln t, \tag{23}$$

and the leading nonpinch terms cancel[12] for diagrams 2(a), 2(b), and 2(c).

[11] J. C. Polkinghorne, J. Math. Phys. 4, 1396 (1963).

[12] The cancellation of the terms involving $g(s)$ between Eqs. (10) and (23) has been verified by Freund (private communication), and it is this cancellation to which he refers in Ref. 9. However, he did not treat the pinch contribution.

In addition, the pinch contribution from diagrams 2(b) and 2(c) is cancelled by the pinch terms from diagrams 2(b′) and 2(c′), as we shall now demonstrate.[13]

The calculation of the amplitudes for diagrams 2(b′) and 2(c′) requires great care. It is true that the *functional* form of the amplitudes for diagrams 2(b′) and 2(c′) is correctly obtained by replacing t by u in Eq. (13) and affixing an over-all negative sign [this occurred for diagrams 2(a) and 2(a′), and again may be seen by simply examining the form of the amplitudes]. However, it is *not* true that the correct *asymptotic* forms for diagrams 2(b′) and 2(c′) may be obtained from Eq. (23) by the same prescription.

The reason is as follows. When we take the limit $t\to+\infty+i\epsilon$ in the logarithmic term of Eq. (16) to determine the pinch contribution, we must take the *same* asymptotic limit in both the crossed and uncrossed diagrams; since $t\to+\infty+i\epsilon$ for diagrams 2(b) and 2(c), we must take $u\to-\infty-i\epsilon$ for diagrams 2(b′) and 2(c′). That is, the analogy to Eq. (16) for the crossed diagrams is

$$\begin{aligned} f_{b'}=f_{c'}\approx &-(1/8\pi W)(-g^6/16\pi^4)u^3 \\ &\times\int_0^1 dz_1\cdots dz_5\delta(z_1+\cdots+z_5-1) \\ &\times C'(z_3/2u)(z_3d+z_4z_5uC')^{-2} \\ &\times\{R+\ln[(d+z_4z_5uz_2)(d+z_4z_5uz_1)/ \\ &\qquad (d-z_4z_5uz_3)(d+z_4z_5u(z_1+z_2))]\}. \end{aligned} \tag{24}$$

However, when $t\to+\infty+i\epsilon$, $u\to-\infty-i\epsilon$ and the logarithmic term becomes

$$\ln[(-\infty-i\epsilon)(-\infty-i\epsilon)/(+\infty+i\epsilon)(-\infty-i\epsilon)]=+i\pi. \tag{25}$$

This introduces an extra negative sign into the pinch contribution, so that while the pinch contribution and the "ordinary" contribution have the same signs for diagrams 2(b) and 2(c) [as in Eq. (23)], they have *opposite* signs for diagrams 2(b′) and 2(c′):

$$f_{b'}+f_{c'}\xrightarrow[t\to+\infty+i\epsilon]{}-[-g(s)+2\pi ih(s)]u\ln u. \tag{26}$$

In other words, because u and t approach different limits, we asymptotically reach different sheets of the same function when we study the large-t forms of diagrams 2(b) and 2(b′).

The total contribution to the asymptotic form of the sixth-order scattering amplitude is the algebraic sum of Eqs. (10), (12), (23), and (26); there are no surviving terms of order $t\ln t$.[14] The terms which remain do not give the asymptotic form of the amplitude because of our previous neglect of terms of order t. Reference 6 did not properly obtain the additional negative sign and so does not exhibit the cancellation of the pinch contributions.

We conclude that no terms of the form $t\ln t-u\ln u$ appear in the scattering amplitude through sixth order; therefore, the Kronecker delta term is not removed and the vector meson remains un-Reggeized to sixth order when interacting with spinless nucleons.

SPIN-$\frac{1}{2}$ NUCLEONS

The analogous problem with spin-$\frac{1}{2}$ nucleons may be reduced to the spinless case. We take nucleon-antinucleon scattering in a theory with interacting nucleons and vector mesons and Reggeize as in Ref. 3. Near a trajectory with $J\approx\alpha$, the partial-wave amplitudes will have the form

$$\begin{aligned} &F_{\mu\nu}{}^{J\pm}\approx\eta_\mu{}^\pm\eta_\nu{}^\pm\alpha(\alpha+1)/(J-\alpha), \\ &F_{\mu0}{}^{J\pm}[J(J+1)]^{-1/2}\approx\eta_\mu{}^\pm\eta_0{}^\pm/(J-\alpha), \\ &F_{0\nu}{}^{J\pm}[J(J+1)]^{-1/2}\approx\eta_0{}^\pm\eta_\nu{}^\pm/(J-\alpha), \\ &F_{00}{}^{J\pm}\approx\eta_0{}^\pm\eta_0{}^\pm/(J-\alpha), \end{aligned} \tag{27}$$

where μ and ν take the values 1 and -1.

From this we find that if the parity-conserving helicity amplitudes are dominated by this pole at large z, they will have the asymptotic form

$$\begin{aligned} &f_{\mu\nu}{}^\pm\xrightarrow[z\to\infty]{}N_\alpha(\pi/\sqrt{2}\sin\pi\alpha)\eta_\mu{}^\pm\eta_\nu{}^\pm(\alpha+1)(-\alpha^2)(-z)^{\alpha-1}, \\ &f_{10}{}^\pm=-f_{01}{}^\pm\xrightarrow[z\to\infty]{}N_\alpha(\pi/\sqrt{2}\sin\pi\alpha) \\ &\qquad\qquad\times\eta_1{}^\pm\eta_0{}^\pm(\alpha+1)\alpha(-z)^{\alpha-1}, \\ &f_{0-1}{}^\pm=-f_{-10}{}^\pm\xrightarrow[z\to\infty]{}N_\alpha(\pi/\sqrt{2}\sin\pi\alpha) \\ &\qquad\qquad\times\eta_0{}^\pm\eta_{-1}{}^\pm(\alpha+1)\alpha(-z)^{\alpha-1}, \\ &f_{00}{}^\pm\xrightarrow[z\to\infty]{}N_\alpha(\pi/\sqrt{2}\sin\pi\alpha)\eta_0{}^\pm\eta_0{}^\pm(\alpha+1)(-z)^\alpha. \end{aligned} \tag{28}$$

As in Ref. 3, we define

$$N_\alpha\equiv\sqrt{2}2^{\alpha+1}\Gamma(\alpha+\tfrac{3}{2})/\pi^{1/2}\Gamma(\alpha+2). \tag{29}$$

We expand as we did in the spinless case, taking

[13] Reference 6 considers the pinch contribution but obtains (-2) times that given in Eq. (19). In that work, the conventional parameterization for the four-dimensional integrals was used. It is essential to explicitly perform one integral to remove the delta function before finding the pinch contribution, but then the remaining integrals have interdependent limits and the pinch only occurs for certain values of the other variables. On the other hand, in the treatment here presented the transition from Eqs. (13), (14), (15) to Eq. (16) could be carried out exactly because the limits on the integrals were simplified by the method of parameterization. Actually, however, the main point we wish to emphasize is not the numerical value of the pinch terms, but the important fact that they cancel between the crossed and uncrossed diagrams. This is shown in the discussion that follows.

[14] This same cancellation is even found if one obtains the pinch contribution following the approach of Ref. 6, giving proper attention to the asymptotic limit. In that method, the logarithmic term obtained is $\ln(+\infty+i\epsilon)(+\infty+i\epsilon)/(-\infty-i\epsilon)(-\infty-i\epsilon)=-2\pi i$ for diagrams 2(b) and 2(c) and $\ln(-\infty-i\epsilon)(-\infty-i\epsilon)/(+\infty+i\epsilon)(+\infty+i\epsilon)=+2\pi i$ for diagrams 2(b′) and 2(c′), and again the pinch terms cancel.

$\alpha=1+\gamma$, to get

$$
\begin{aligned}
f_{\mu\nu}{}^{\pm} &\sim -\alpha^2\eta_\mu{}^{\pm}\eta_\nu{}^{\pm}(1+\gamma \ln z+\cdots)\,, \\
f_{10}{}^{\pm} &= -f_{01}{}^{\pm}\sim\alpha\eta_1{}^{\pm}\eta_0{}^{\pm}(1+\gamma \ln z+\cdots)\,, \\
f_{0-1}{}^{\pm} &= -f_{-10}{}^{\pm}\sim\alpha\eta_0{}^{\pm}\eta_{-1}{}^{\pm}(1+\gamma \ln z+\cdots)\,, \\
f_{00}{}^{\pm} &\sim -\eta_0{}^{\pm}\eta_0{}^{\pm}z(1+\gamma \ln z+\cdots)\,,
\end{aligned}
\tag{30}
$$

except for an over-all multiplicative factor.

The Born-approximation diagrams appear in Fig. 3; in the asymptotic limit, the negative parity amplitudes arise only from the pole term 3(a). The independent amplitudes are at large z

$$
\begin{aligned}
f_{11}{}^{-} &\rightarrow (-g^2/8\pi W)(s-\lambda^2)^{-1}(-8E^2)\,, \\
f_{10}{}^{-} &\rightarrow (-g^2/8\pi W)(s-\lambda^2)^{-1}(8mE)\,, \\
f_{00}{}^{-} &\rightarrow (-g^2/8\pi W)(s-\lambda^2)^{-1}(-8m^2z)\,.
\end{aligned}
\tag{31}
$$

(E= nucleon CM energy, λ= vector meson mass, m= nucleon mass). This of course looks like the lowest order term of Eq. (32), with

$$\alpha(\eta_1{}^{-}/\eta_0{}^{-})=E/m\,. \tag{32}$$

We are now prepared to look at sixth order in the spinor nucleon scattering problem. Rather than $(-8t^3)$ in the numerators as before, the dominant terms for diagram 2(a) at large z are the "explicit" terms

$$
\begin{aligned}
&(\bar{u}_3|\gamma_\mu(p_3+m)\gamma_\nu(p_3+m)\gamma_\sigma|v_4) \\
&\qquad\times(\bar{v}_2|\gamma_\sigma(p_1+m)\gamma_\nu(p_1+m)\gamma_\mu|u_1)\,.
\end{aligned}
\tag{33}
$$

For diagram 2(b), again the "explicit" terms dominate at large z:

$$
\begin{aligned}
&(\bar{u}_3|\gamma_\mu(p_3+m)\gamma_\nu(p_3+m)\gamma_\sigma|v_4) \\
&\qquad\times(\bar{v}_2|\gamma_\sigma(p_1+m)\gamma_\mu(p_1+m)\gamma_\nu|u_1)\,.
\end{aligned}
\tag{34}
$$

However, both these forms may be reduced to

$$16(p_1\cdot p_3)^2(\bar{u}_3|\gamma_\sigma|v_4)(\bar{v}_2|\gamma_\sigma|u_1)\,, \tag{35}$$

so that the amplitude for Fig. 3(a) explicitly factors out of each of the sixth-order asymptotic forms. We thus

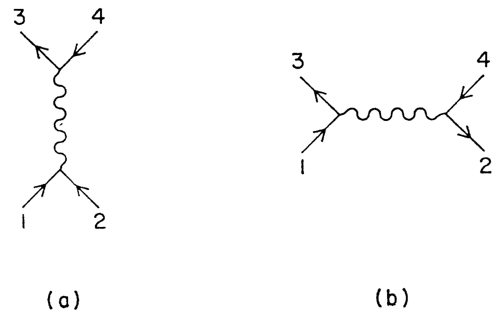

FIG. 3. Second-order terms.

find that for large z, the diagrams of Fig. 2 give

$$
\begin{aligned}
f_{11}{}^{-} &\rightarrow (-1/2t)(-8E^2)f^{(6)}\,, \\
f_{10}{}^{-} &\rightarrow (-1/2t)(8mE)f^{(6)}\,, \\
f_{00}{}^{-} &\rightarrow (-1/2t)(-8m^2z)f^{(6)}\,,
\end{aligned}
\tag{36}
$$

where $f^{(6)}$ is the asymptotic form of the sixth-order scattering amplitude for *spinless* nucleons:

$$f^{(6)}=f_{\mathrm{a}}+f_{\mathrm{b}}+f_{\mathrm{c}}+f_{\mathrm{a}'}+f_{\mathrm{b}'}+f_{\mathrm{c}'}\,. \tag{37}$$

Therefore, through sixth order, the amplitude *does* factor in the manner of Eq. (30) at large z, and again the Reggeization depends upon there being surviving terms of order $t\ \ln t$ in $f^{(6)}$. Since the $t\ \ln t$ terms cancelled for spinless nucleons, we may conclude that the vector meson is not Reggeized in spinor nucleon-antinucleon scattering through sixth order. This result agrees with Mandelstam's treatment of the same problem by quite different methods.[7]

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