PHYSICAL REVIEW D VOLUME 6, NUMBER 3 1 AUGUST 1972

Phenomenology Based on Tachyon Exchange

A. M. Gleeson, M. G. Gundzik, and E. C. G. Sudarshan
Center for Particle Theory, University of Texas at Austin, Austin, Texas 78712

and

Antonio Pagnamenta
University of Illinois, Chicago Circle, Chicago, Illinois 60680
(Received 10 June 1971)

The observation that the existence of tachyonic resonances can lead to peaks in differential cross sections implies a phenomenology of high-energy processes. A large class of hadronic differential cross sections shows a common peak at $-t \simeq 1.0$ (GeV/c)^2. We present a detailed fit of the K^-p and π^-p charge-exchange (CEX) differential cross sections and determine that both processes are consistently fitted by using a tachyonic resonance with $m^2 = -1.0$ GeV^2. and $\Gamma = 580$ MeV. The spin of the resonance determines the energy shrinkage rate and is of the form $j = -\frac{1}{2} + i\sigma$. A value of $\sigma = 0.15$ leads to an adequate fit to the data. The other quantum numbers of this strongly interacting tachyon are $I = 1$, G parity even, and Y parity even. Other possible examples of tachyonic resonances are described.

I. INTRODUCTION

The existence of peaks in particle cross sections as a function of the pertinent squared-missing-mass (i.e., the invariant of the missing four-momentum) variable is the first step in the identification of resonance states.[1] This research has been naturally restricted to positive values of the squared missing mass. Here we propose a similar interpretation for the explanation of bumps at negative mass squared. Such a resonance is the unstable version of the tachyonic object first proposed by Bilaniuk, Deshpande, and Sudarshan.[2] Its possible exchange was suggested by Sudarshan.[3]

Since the early proposals for the existence of tachyons several preliminary experiments have been undertaken.[4-7] These experiments did not find the stable tachyons to which they were sensitive. However, the existence of unstable tachyons remains an open question.[8] It is to this question that we address ourselves. As a preliminary approach to this problem we propose that the success of a phenomenology based on tachyonic resonances may provide new insight on which to base a search for the effects of tachyonic matter. Our approach is applicable to the exchange of unstable tachyons and does not depend on the existence of freely propagating tachyonic states.

The most obvious candidates for tachyonic resonances are those which correspond to peaks in differential cross sections which persist from one process to the next. A study of the available hadronic differential cross sections indicates that a wide class of reactions shows a persistent peak at $-t \simeq 1.0$ (GeV/c)^2. The conventional Regge analysis of these reactions concentrates on parametrizing the dip at $-t \simeq 0.6$ (GeV/c)^2. In the following we show that π^-p and K^-p charge-exchange (CEX) reactions can be fit by a phenomenology based on the exchange of a tachyonic resonance with a negative squared mass of -1.0 GeV^2.

In Sec. II, we enumerate the differential cross sections which show the peak at $t = -1.0$ (GeV/c)^2 and discuss those in which it is apparently absent. We propose that the peak be attributed to a tachyonic resonance with the quantum numbers of the ρ meson. This hypothesis is investigated in detail in Sec. III, where we construct a model consisting of a smooth forward peak parametrized in a quasi-Regge fashion as background plus a tachyonic resonance contribution at $t = -1.0$ (GeV/c)^2. The model is then used to fit the π^-p and K^-p CEX reactions. The former requires the ρ meson and the tachyon to produce the secondary peak, while the latter, which consists of the ρ, the tachyon, and the A_2, shows a small shoulder in $d\sigma/dt$ near $t = -1.0$ (GeV/c)^2. Finally, in Sec. IV, we present some general conclusions based on this analysis.

II. REVIEW OF HADRONIC DIFFERENTIAL CROSS SECTIONS

The hadronic differential cross sections have long been characterized by the presence or absence of forward and backward peaks, which can be correlated with the existence of exchanges associated with positive-m^2 particles. Peaks in the interior regions of the hadronic differential cross section have long been explained by other mechanisms. Different authors relate this structure to various effects such as multiple scattering from fundamental constituents or interference between Regge cuts and poles or wrong-signature nonsense zeros. Reviewing the limited success of all these models, it is appropriate to ask whether tachyon exchange would not in fact be a simpler description.

One finds that many hadronic differential cross sections show a persistent conspicuous peak at $-t \simeq 1.0$ (GeV/c)^2. It appears distinctly in the following reactions[9,10]:

$$\pi^+ p \to \pi^+ p, \quad \text{Ref. 9(a)} \tag{2.1}$$

$$\pi^- p \to \pi^- p, \quad \text{Ref. 9(b)} \tag{2.2}$$

$$\pi^- p \to \pi^0 n, \quad \text{Ref. 9(c)} \tag{2.3}$$

$$\pi^+ p \to \pi^0 \Delta^{++}, \quad \text{Ref. 9(d)} \tag{2.4}$$

$$\gamma p \to \pi^0 p, \quad \text{Ref. 9(e)} \tag{2.5}$$

$$\bar{p} p \to \bar{p} p, \quad \text{Ref. 9(f)}. \tag{2.6}$$

There is no discernible peak at 1.0 (GeV/c)^2 in

$$pp \to pp, \quad \text{Ref. 10(a)} \tag{2.7}$$

$$K^- p \to \bar{K}^0 n, \quad \text{Ref. 10(b)} \tag{2.8}$$

$$\pi^- p \to \omega n, \quad \text{Ref. 10(c)} \tag{2.9}$$

$$\pi^- p \to \eta^0 n, \quad \text{Ref. 10(b)} \tag{2.10}$$

$$\pi^+ p \to \eta^0 \Delta^{++}, \quad \text{Ref. 10(d)} \tag{2.11}$$

and many others.

It is evident from the above reactions that the peak only appears distinctly in those channels where ρ exchange is possible. It is most prominent if ρ exchange is the only ($m^2 > 0$) exchange and is not so clear or completely absent if in addition to the ρ other exchanges are also possible. It is definitely absent in (2.10) and (2.11) where only A_2 exchange is allowed. We therefore conjecture that the tachyonic resonance, denoted by τ, which is responsible for the peak at $t = -1.0$ (GeV/c)^2 has the strong-interaction quantum numbers[11] of the ρ meson,

$$I_\tau = +1, \quad G_\tau = +1, \quad Y_\tau = +1.$$

The absence of a τ peak at this mass in pp scattering is qualitatively understandable by the presence of the many helicity amplitudes and the competition of other exchange quantum numbers. The reaction (2.9), $\pi^- p \to \omega n$, allows ρ, B, and τ exchange. $K^- p \to \bar{K}^0 n$ allows ρ, A_2, and τ exchange. In these cases the additional exchanges can easily mask the effects of the τ. In Sec. III we describe a detailed quantitative fit for the process $K^- p \to \bar{K}^0 n$ and find that this explanation is completely consistent with the experimental data.

We observe that in process (2.11), $\pi^+ p \to \eta^0 \Delta^{++}$, which is dominated by the exchange of the A_2, there are good indications of a secondary peak at $-t \simeq 2.5$ (GeV/c)^2. This is therefore a candidate for a tachyonic resonance with the quantum numbers of the A_2. In order for this assignment to be appropriate, additional data must show the persistence of this peak in all A_2-dominated processes and shrinkage of the peak consistent with $\alpha_{\text{eff}}(t) = -\frac{1}{2}$ in this region. This last requirement is a consequence of the spin assignments of the tachyon in our model and will be described in Sec. III. An indication of an ω-like tachyonic resonance could be inferred from the process $\pi^+ p \to \rho^+ p$. Since the differential cross-section data[12] extends to $-t \approx 0.9$ (GeV/c)^2 with no distinct secondary peak evident, the suggestion is that no ω-like tachyonic resonance is present with (mass)$^2 \geq -0.9$ GeV^2.

II. THE TACHYONIC RESONANCE MODEL

In our formulation of two-body scattering, the complete s-channel amplitude $F(s,t)$ is separated into two parts:

$$f(s,t) = f^b(s,t) + f^\tau(s,t).$$

We consider particles with initial four momenta p_1 and p_2, final momenta p_3 and p_4 with $s = (p_1 + p_2)^2$, $t = (p_1 - p_3)^2$. The corresponding helicities are λ_i, $i = 1, \ldots, 4$.

The amplitude $f^b(s,t)$ describes the scattering away from the tachyonic peak and is treated as background. The dominant feature of this background is the forward (or backward) peak which is identified with the exchange of particles with $m^2 > 0$. We therefore parametrize this amplitude with the t dependence generated by the nearest poles. The shrinkage of the forward peak is determined by a linear trajectory consistent with recent Regge phenomenology. This amplitude is responsible for the smooth behavior of $d\sigma/dt$ near the forward direction and is completely determined by the data in this region.

The amplitude $f^\tau(s,t)$ corresponds to the negative-mass-squared resonance or tachyonic resonance contribution. The parametrization of this amplitude is determined by the secondary peak structure of $d\sigma/dt$. We describe the parametrization of this part of the amplitude in the next section. Then we discuss the explicit forms of $f^\tau(s,t)$ and $f^b(s,t)$ in our model for $\pi^- p$ and $K^- p$ scattering.

A. Formulation of f^τ

If we wish to describe two particle scattering for $s > 0$ in terms of s-channel properties, e.g., resonances, etc., then the usual partial-wave decomposition is appropriate. That is, the decomposition of the scattering amplitude is in terms of the little group SO(3) of the Poincaré group. Since we wish to describe scattering in terms of t-channel exchanges for the case of $t<0$, the appropriate group-theoretic decomposition of the scattering amplitude in terms of representations of the Poincaré group is with respect to its little group SO(2, 1). That is, we determine the form of the t-channel amplitude and then construct the physical s-channel counterpart by a crossing transformation from the t to the s channel.

We assume that the part f^τ of the scattering amplitude, i.e., the full amplitude minus the forward "background" contribution f^b, is sufficiently well behaved in s and t such that it can be decomposed in terms of a single member of a unitary irreducible representation (UIR) of the Poincaré group. Specifically, we are interested in the UIR corresponding to $t<0$, i.e., those corresponding to the little group SO(2, 1). This decomposition in the t channel and transformation to the s channel leads to the form[13]

$$\begin{aligned}F^\tau_{\lambda_3\lambda_4,\lambda_2\lambda_1}(s,t) &= \langle p_3\lambda_3, p_4\lambda_4|\tilde{T}|p_1\lambda_1, p_2\lambda_2\rangle \\ &= \sum_{\mu_1\mu_3} A_s(\mu_1,\mu_3;\lambda_1\lambda_3)\Big(\int_{1/2-i\infty}^{1/2+i\infty} dj\,\frac{2j+1}{\tan\pi j}\,\langle\lambda_2\lambda_4||\tilde{T}^j(t)||\mu_1\mu_3\rangle d^j_{\lambda_4-\lambda_2,\mu_1-\mu_3}(\cosh\beta) \\ &\qquad + \sum_{-M\le j\le -1}(2j+1)\,\langle\lambda_2\lambda_4||\tilde{T}^{j\pm}(t)||\mu_1\mu_3\rangle\, d^{j\pm}_{\lambda_4-\lambda_2,\mu_1-\mu_3}(\cosh\beta)\Big)\,, \end{aligned} \tag{3.1}$$

where μ_1 and μ_3 are summed over the possible helicities of particles 1 and 3, respectively. Here $d^j_{\lambda,\lambda'}$ is the d function for the principal series, $d^{j\pm}_{\lambda,\lambda'}$ is that for discrete series, with d^{j+} appearing when both λ and λ' are positive and d^{j-} when both are negative. Further, we have

$$\begin{aligned}\cosh\beta &= \frac{2(s-m_1{}^2-m_2{}^2)+t}{[(4m_1{}^2-t)(4m_2{}^2-t)]^{1/2}}\,, \\ M &= \min(|\mu_1-\mu_3|, |\lambda_4-\lambda_2|)\,,\end{aligned} \tag{3.2}$$

and $A_s(\mu_1,\mu_3;\lambda_1\lambda_3)$ consists of spin rearrangement factors from the t- to s-channel transformation given in Ref. 13.

The essential ingredient of our model based on our parametrization in terms of tachyonic resonances is that the reduced matrix element in (3.1) has the form

$$\begin{aligned}F^{\tau j}_{\lambda_2\lambda_4;\lambda_1\lambda_3} &\equiv \langle\lambda_2\lambda_4||\tilde{T}^j(t)||\lambda_1\lambda_3\rangle \\ &= \frac{\operatorname{Res}(\langle\lambda_4\lambda_2||\tilde{T}^{j_0}(t)||\lambda_1\lambda_3\rangle)\delta(j-j_0)}{t+m^2-i|m|\Gamma}\,,\end{aligned} \tag{3.3}$$

where Γ is the full width at half maximum. The value of j_0 is a characteristic of the tachyonic resonance. It is selected from the range of values available in the expansion (3.1). The assumption that the spin of the tachyon has a definite value j is not necessary, but the range of values should be restricted. This restriction is, in a sense, the definition of the tachyon as a single particle with a definite spin whose exchange dominates the reaction in this range of t. The value of $\operatorname{Re} j_0$ for the principal series is $-\frac{1}{2}$, which determines the rate of shrinkage of the secondary peak in $d\sigma/dt$.

The form for $F^{\tau j}_{\lambda_2\lambda_4;\lambda_1\lambda_3}(t)$ is based on the simple single-level Breit-Wigner resonance amplitude. It is generally used for resonance enhancements for positive-mass-squared resonances because it is sharply peaked around the resonance mass and has a phase consistent with unitarity. Since in our case the t-channel amplitude has a unitarity relationship which requires that there be no phase, we modify the simple level form by adding its complex-conjugate amplitude. Thus in our case we use $F^{\tau j}_{\lambda_2\lambda_4;\lambda_1\lambda_3}$ to construct a reduced amplitude that is sharply peaked around the resonance mass but real, in keeping with the appropriate unitarity requirement. Since the s dependence of the tachyonic contribution is determined by a real analytic function, the total tachyonic amplitude satisfies

$$f^\tau(s,t) = \tfrac{1}{2}[F^\tau(s,t^*) + (F^\tau(s,t))^*]\,. \tag{3.4}$$

The explicit form of F^τ, and thus f^τ, is determined by the process under consideration. For simplicity we turn to the analysis of the scattering of a spin-$\frac{1}{2}$ and spin-0 particle into a spin-$\frac{1}{2}$ and spin-0 particle. In that case only the integral of (3.1) contributes and the spin rearrangement factor is unity. As examples of the above we consider the analysis of $\pi^- p \to \pi^0 n$ and $K^- p \to \bar{K}^0 n$.

B. Analysis of $\pi^- p \to \pi^0 n$

Here p_1, λ_1 and p_3, λ_3 refer to the initial and final four-momentum and helicity of the π's of mass m_π, while p_2, λ_2 and p_4, λ_4 refer to the initial and final nucleons of mass m_p. The parametrization of the background amplitude $f^b(s,t)$ is consistent with ρ dominance of the forward peak and follows our earlier preliminary approach[8] to this problem. The tachyon exchange contribution

$f^{\tau}(s,t)$, however, is modified from our earlier approach in that $f^{\tau}(s,t)$ is now a real analytic function.

The two independent amplitudes $f_{++}(s,t)$ and $f_{+-}(s,t)$ for this process correspond to the helicity-nonflip and -flip contributions, respectively. The nucleon helicities $\lambda=\pm\frac{1}{2}$ are denoted by $\pm$. We express each of the independent amplitudes in terms of the forward "background" contribution denoted $f^{\rho}_{+\pm}(s,t)$, due in this case to the ρ, and $f^{\tau}_{+\pm}(s,t)$. We have

$$f_{+\pm}(s,t)=f^{\rho}_{+\pm}(s,t)+f^{\tau}_{+\pm}(s,t)$$

as the possible amplitudes.

The phenomenological form for the amplitudes f^{ρ}_{++} and f^{ρ}_{+-} has the ρ pole, ρ-Regge pole asymptotic behavior, and the appropriate t-dependent kinematic factors. The kinematic factors are the conventional t-channel threshold terms and are included for completeness. They do not have much influence on the quality of the fit and we emphasize the most important features of the fit are qualitative; the peak position stability with various energies and the shrinkage rate of the shoulder. We assume the simple forms

$$f^{\rho}_{++}(s,t)=\frac{b_{++}(s/s_{0\rho})^{\alpha_\rho(t)}}{(t-4m_p^2)^{1/2}(t-m_\rho^2)}, \tag{3.5a}$$

$$f^{\rho}_{+-}(s,t)=\frac{ib_{+-}(s/s_{0\rho})^{\alpha_\rho(t)}}{t-m_\rho^2}\,[t(t-4m_\pi^2)]^{1/2}, \tag{3.5b}$$

where $\alpha_\rho(t)=0.58+t$, and $s_{0\rho}$, b_{++}, and b_{+-} are real parameters determined by fits to experiment.[14]

The form of F^{τ} for this process is determined from (3.1) and (3.3). In addition we have introduced the same t-channel threshold factors as for f^{ρ}. This provides us with the following:

$$F^{\tau}_{++}(s,t)=\frac{2j_0+1}{(t-4m_p^2)^{1/2}}\,\frac{\lambda_{++}}{t+m_0^2-i\Gamma|m_0|}\times\frac{d^{j_0}_{00}(\cosh\beta)}{\tan\pi j_0}, \tag{3.6a}$$

$$F^{\tau}_{+-}(s,t)=(2j_0+1)[t(t-4m_\pi^2)]^{1/2}\times\frac{\lambda_{+-}}{t+m_0^2-i\Gamma|m_0|}\,\frac{d^{j_0}_{01}(\cosh\beta)}{\tan\pi j_0}. \tag{3.6b}$$

The final real analytic functions $f^{\tau}_{++}(s,t)$ and $f^{\tau}_{+-}(s,t)$ are determined from Eqs. (3.4) and (3.6). Equations (3.6) depend upon the mass m_0, the width Γ, and the spin j_0 of the tachyonic resonance, as well as the coupling constants λ_{++} and λ_{+-}. These five parameters are all determined from experimental fits. j_0 can take on a range of values permitted by $j_0=-\frac{1}{2}+i\sigma$, with σ determined from experimental data. $\mathrm{Re}\,j_0=-\frac{1}{2}$ occurs naturally from (3.1) without any further assumptions and is consistent with the shrinkage in s of $d\sigma/dt$ for the π^-p CEX reaction in the vicinity of the secondary peak, i.e., $t\simeq-1.0$ $(\mathrm{GeV}/c)^2$, corresponding to an $s^{-1/2}$ dependence in the amplitude. This feature of the data is most prominently displayed in Fig. 1, the well-known plot of $\alpha_{\mathrm{eff}}(t)$ for the π^-p CEX reaction. The observation that the effective trajectory flattens out around $t\approx-1.0$ $(\mathrm{GeV}/c)^2$ to the value of $\alpha_{\mathrm{eff}}\approx-\frac{1}{2}$ is one of the necessary features of tachyonic resonance dominance. The other is the persistence of the position of the peak at all energies and in all suitable reactions at the same t.

The data 9(c) for $\pi^-p\to\pi^0n$ were fitted to the above model. The fit with the best χ^2 for $d\sigma/dt$ is given in Fig. 2 with the corresponding fit for polarization[15] in Fig. 3. The parameters for the tachyonic resonance are

$$m_0^2=-1.0\ \mathrm{GeV}^2,$$

$$\Gamma=580\ \mathrm{MeV},$$

$$\sigma=0.15.$$

The coupling strengths λ_{++} and λ_{+-} are listed in Table I.

C. Analysis of $K^{\pm}p\to\overline{K}^0n$

In the case of the K^-p CEX reaction, exchanges with the quantum numbers of the A_2 as well as ρ

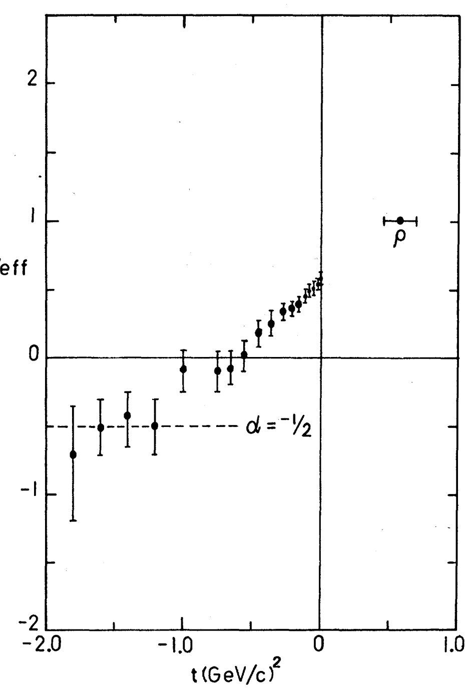

FIG. 1. The effective trajectory $(s/s_0)^{\alpha_{\mathrm{eff}}(t)}$ as extracted from fits to $d\sigma/dt\,(\pi^-p\to\pi^0n)$. The low-$t$ points extrapolate smoothly to the ρ meson. At larger negative t the effective trajectory flattens out to $\alpha_{\mathrm{eff}}=-\frac{1}{2}$.

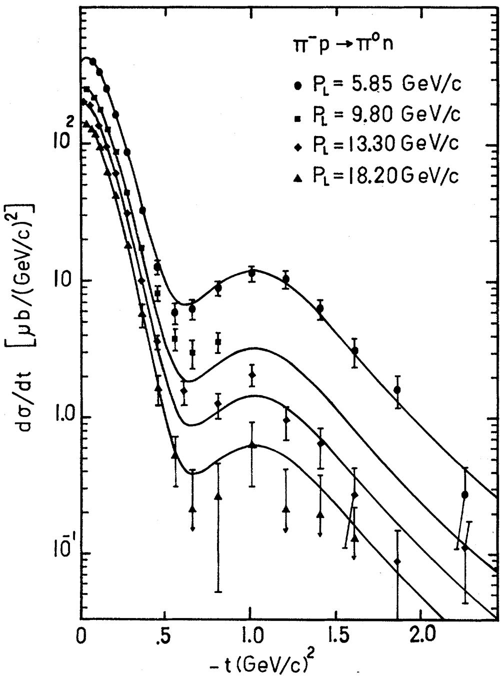

FIG. 2. The $\pi^- p \to \pi^0 n$ differential cross section data [Ref. 9(c)] fitted by $\rho + \tau$ exchange.

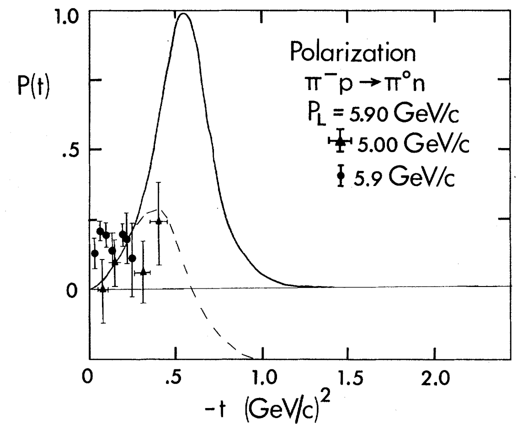

FIG. 3. The polarization (Ref. 15) in $\pi^- p \to \pi^0 n$ fitted with $\rho + \tau$ exchange. The solid line is obtained if f_{+-} has no factor $\alpha(t)$, the dashed line if a factor $\alpha(t)$ is added to f_{+-}. The experimental points for $P_L = 5.00$ GeV/c are those of Drobnis *et al.*, and for 5.90 GeV/c are those of Bonamy *et al.*

meson are possible. In accord with our discussion of Sec. II, we assume that the τ has the same quantum numbers as the ρ and that there is no A_2-like tachyonic resonance present in the process under consideration.

From this point of view the independent $K^- p$ CEX amplitudes $f_{++}(s,t)$ and $f_{+-}(s,t)$ can be written as the sum of three different contributions:

$$f_{++}(s,t) = f_{++}^{\rho} + f_{++}^{A}(s,t) + f_{++}^{\tau}(s,t)\,, \qquad (3.7)$$

with the corresponding form for $f_{+-}(s,t)$. Here the functions $f^{\rho}(s,t)$ are given by (3.5a) and (3.5b) with different parameters b_{++} and b_{+-} and the $f^{\tau}(s,t)$ are given by (3.4) and (3.6) with j_0, $m_0{}^2$, and Γ the same as for the $\pi^- p$ CEX reaction, but different parameters λ_{++} and λ_{+-}. The form for $f^{A}(s,t)$, in analogy to $f^{\rho}(s,t)$, is given as follows:

$$f_{++}^{A}(s,t) = \frac{c_{++}(s/s_{0A})^{\alpha_A(t)}}{(t-4m_p{}^2)^{1/2}(t-m_A{}^2)}\,, \qquad (3.8a)$$

$$f_{+-}^{A}(s,t) = \frac{c_{+-}(s/s_{0A})^{\alpha_A(t)}}{t-m_A{}^2}\,[t(t-4m_K{}^2)]^{1/2}. \qquad (3.8b)$$

Here we have new real parameters c_{++}, c_{+-}, and s_{0A} for the A_2, with m_A and m_K referring to the A_2 and K masses, respectively. In addition, we have taken for the A_2 trajectory function $\alpha_A(t) = 0.34 + t$.

The best fit to the $K^- p \to \bar{K}^0 n$ data[16] using the same values for $m_0{}^2$, Γ, and j_0 as in the $\pi^- p$ CEX reaction appears in Fig. 4. The effect of the τ is to produce a hardly noticeable shoulder in the differential cross section at $t = -1.0$ (GeV/c)^2, in good agreement with the data. The A_2 coupling is stronger than that of the ρ and τ. Since there are no polarization data for this problem and our model does not predict its sign, we have not included a $K^- p$ CEX polarization curve.

IV. DISCUSSION

The preceding analysis of the $\pi^- p$ CEX and $K^- p$ CEX reactions shows that the use of tachyonic resonance exchange provides an adequate fit to the angular distribution data available for both of these processes. The polarization data on the $\pi^- p$ CEX reaction, although not very precise, are consistent with this model. This model fits these data with a good χ^2 and no more parameters than the more conventional Regge approaches. In order to account for polarization in $\pi^- p$ CEX reactions, the Regge models require either two mesonlike nondegenerate trajectories or cuts. In the tachyon resonance model, the form that the polarization takes is dependent on the parametrization of the amplitudes. One finds small polarization with a

TABLE I. Coupling parameters for the $\pi^- p$ and $K^- p$ CEX reactions.

	b_{++}	b_{+-}	c_{++}	c_{+-}	λ_{++}	λ_{+-}	$S_{0\rho}$	S_{0A}
$\pi^- p$ CEX	−6.24	54.5	⋯	⋯	-1.1×10^3	−32.0	0.0082	⋯
$K^- p$ CEX	34.1	140.5	468.0	469.	691.0	0	0.517	0.228

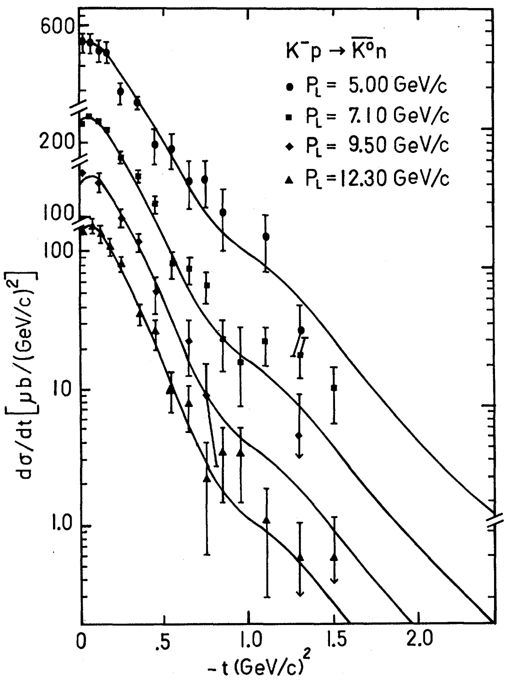

FIG. 4. The fit to the available data (Ref. 16) on $K^-p \to \bar{K}^0 n$ with $\rho + A_2 + \tau$ exchange.

no-nonsense wrong-signature zero and large polarization otherwise (see Fig. 3). In this sense, the presence of polarization, but not its form, is a feature of the tachyon exchange.

Besides the tachyonic resonance at a mass squared of -1.0 GeV^2 there are indications of a phenomenological tachyonic resonance with the A_2 quantum numbers and mass squared around -2.5 GeV^2. Data at different energies would test the shrinkage rate and peak position stability of this resonance.

Experiments have only searched for tachyons as stable particles with the absolute value of mass squared up to 1.0 GeV^2. The tachyons considered in this paper differ significantly from the ones looked for in these experiments. There is an intrinsic spread in the mass of the tachyons, a phenomenon which may be loosely referred to as intrinsic instability. As a consequence of this spread in mass, the tachyon cannot propagate freely for distances large compared with the inverse width, that is, for distances large compared with nuclear dimensions. This in turn means that for such a resonance to manifest itself in, say, a missing-mass experiment, another hadronic particle should be available within nuclear distances to which the tachyon can attach itself. The experimental realization of these conditions must therefore involve targets with composite structure. Analysis of reactions in which a fraction of a composite target (say one of the nucleons in a deuteron or carbon target) is taken as one of the colliding particles seems to be a natural method for searching for tachyons. Such tachyonic resonances should not therefore be expected to make visible tracks in bubble chambers or activate counters. These tachyonic resonances could therefore not have been observed in any of the previous tachyon searches.

In the previous sections we have described and analyzed the implications of a phenomenology based on tachyon exchange. The tachyons utilized to fit the data are resonance states associated with the principal series of SO(2, 1). The consistency of this phenomenology with the observed data is an indication that the assumption of the existence of tachyonic resonance states may provide a useful description of high-energy interactions. This suggests that experiments might be performed which seek to isolate other manifestations of these resonances.

---

[1]G. Pinski, E. C. G. Sudarshan, and K. T. Mahanthappa, in *Proceedings of the Tenth Annual International Rochester Conference on High-Energy Physics, 1960,* edited by E. C. G. Sudarshan, J. H. Tinlot, and A. C. Melissinos (Interscience, New York, 1960); B. C. Maglic, L. W. Alvarez, A. H. Rosenfeld, and M. L. Stevenson, Phys. Rev. Letters 7, 178 (1961); A. R. Erwin, R. March, W. D. Walker, and E. West, *ibid.* 6, 628 (1961).

[2]O. M. Bilaniuk, V. K. Deshpande, and E. C. G. Sudarshan, Am. J. Phys. 30, 718 (1969); O. M. Bilaniuk and E. C. G. Sudarshan, Phys. Today 22 (No. 5), 43 (1969); G. Feinberg, Sci. Am. 222 (No. 2), 68 (1970); G. Feinberg, Phys. Rev. 159, 1089 (1967).

[3]E. C. G. Sudarshan, Arkiv Physik 39, 585 (1969).

[4]T. Alväger and P. Erman, Nobel Institute report, Stockholm, 1966 (unpublished); T. Alväger and M. N. Kreisler, Phys. Rev. 171, 1357 (1968).

[5]C. Baltay, G. Feinberg, N. Yeh, and R. Linsker, Phys. Rev. D 1, 759 (1970).

[6]P. V. Ramana Murthy, Lett. Nuovo Cimento 1, 908 (1971).

[7]J. S. Danburg *et al.*, Phys. Rev. D 4, 53 (1971).

[8]A. M. Gleeson, M. G. Gundzik, E. C. G. Sudarshan, and Antonio Pagnamenta, Particles and Nuclei 1, 1 (1970).

[9]Here we list the sources of the experimental data. (a) D. D. Allen *et al.*, Phys. Letters 21, 468 (1966). (b) C. P. Coffin *et al.*, Phys. Rev. Letters 15, 838 (1966). (c) A. V. Stirling *et al.*, *ibid.* 14, 763 (1965); P. Sonderegger *et al.*, Phys. Letters 20, 75 (1966). (d) G. Gidal *et al.*, Phys. Rev. Letters 23, 994 (1969). (e) R. Anderson *et al.*, *ibid.* 21, 384 (1968). (f) M. Katz *et al.*, Phys.

Letters 19, 265 (1967).

[10](a) A. Allaby *et al.*, Phys. Letters 28b, 67 (1968). (b) R. O. Mathews, Nucl. Phys. B11, 339 (1969). (c) J. McLoed (private communication). (d) D. F. Grether and G. Gidal, Phys. Rev. Letters 26, 792 (1971).

[11]The operation of Y parity is defined as $Y \equiv e^{i\pi J_2}P$ and is the suitable extension of space reflection parity to the case of spacelike four-momenta. It is not the usual type of discrete symmetry operation since states of nonzero helicity are not eigenstates of Y. The Y parity of the ρ meson in the zero-helicity configuration is also even, since its space parity is odd.

[12]G. V. Dass and C. D. Froggatt, Nucl. Phys. B8, 661 (1968).

[13]J. F. Boyce, R. Delbourgo, A. Salam, and J. Strathdee, International Center for Theoretical Physics, Trieste Report No. IC/67/9, 1967 (unpublished).

[14]In addition to these kinematic factors, there is often included in b_{+-} an additional factor of $\alpha(t)$. This nonsense-wrong-signature zero does not affect the fit $d\sigma/dt$ in our model but does markedly change the nature of the polarization. Since the polarization data for the CEX processes is not very complete, no choice can be made about the appropriate form to utilize.

[15]P. Bonamy *et al.*, Phys. Letters 23, 501 (1966); D. D. Drobnis *et al.*, Phys. Rev. Letters 20, 274 (1968); P. Bonamy *et al.*, Nucl. Phys. B16, 335 (1970).

[16]P. Astbury *et al.*, Phys. Letters 23, 396 (1966).