

Photoproduction of the scalar mesons $f_0(500)$ and $f_0(980)$ off the nucleon

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We investigate photoproduction of scalar mesons off the nucleon, using the effective Lagrangians and the Regge approach. We first study $f_0(980)$ photoproduction, replacing the Feynman propagator with Regge ones for ρ -meson exchange in the t channel. The model parameters are fixed by reproducing the experimental data on the differential cross section. We then apply the same method to $f_0(500)$ or σ photoproduction with the same parameters for the Regge propagator. Since the threshold energy of the $f_0(500)$ production is rather low, N^* resonances, which can decay into a nucleon and two pions in the isoscalar and scalar states, can come into play in the s channel. To examine the effects of the N^* resonances, we set up two models to look into the respective contribution of the N^* resonances. We find that, in particular, $N(1440)$ and $N(1680)$ play important roles in $f_0(500)$ photoproduction. We discuss the physical implications of the present results.
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1. Introduction

Understanding the structure of low-lying scalar mesons has been one of the most challenging issues in hadronic physics. Their internal structure is still under debate. That the $f_0(500)$ scalar meson, which is also known as σ , is not an ordinary meson consisting of a quark and an antiquark is more or less in consensus. Recent studies suggest that these scalar mesons may belong to the flavor SU(3) non- $q\bar{q}$ nonet (see Refs. [1,2], a “note on scalar mesons below 2 GeV” in Ref. [3], and references therein. A recent review provides also various pieces of information on the structure of the scalar meson [4], including a historical background of the σ meson). The $f_0(500)$ is also interpreted as one of the glueballs or gluonia, mixed with the $\bar{q}q$ state [5–7], though this idea is criticized because the same analysis is rather difficult to apply to an explanation of the strange scalar meson $K_0^*(800)$ or κ , which is also considered as a member of the nonet. The $f_0(500)$ is often regarded as a tetraquark state in a broad sense [8]. The $f_0(500)$ as a tetraquark state has a multiple meaning: It can be described as a diquark–antidiquark correlated state [9,10], $\bar{q}q\bar{q}q$ state [11], or correlated 2π state [12,13] arising from $\pi\pi$ scattering. This non- $\bar{q}q$ feature has been employed in various theoretical approaches such as quantum chromodynamics (QCD) sum rules [14], effective Lagrangians [15], and lattice QCD [16–18].

The scalar mesons have also been extensively studied phenomenologically. There are two scalar–isoscalar mesons ($I^G(J^{PC}) = 0^+(0^{++})$) below 1 GeV, i.e., the lowest-lying $f_0(500)$ (or σ) and the first excited $f_0(980)$. Both the $f_0(500)$ and the $f_0(980)$ exist in $\pi\pi$ scattering and their pole positions have been investigated based on many different processes, such as, e.g., $\pi N \rightarrow \pi\pi N$ reactions [19–21], K_{14} decay [22,23], $D \rightarrow 3\pi$ [24,25], $J/\psi \rightarrow \omega\pi\pi$ [26], $\psi(2S) \rightarrow \pi^+\pi^-J/\psi$ [27], $\gamma\gamma \rightarrow \pi\pi$ [28], pp scattering [29], and so on (for details, we refer to Refs. [3,4]). While the mass and the width of the $f_0(980)$ are more or less known to be $m_{f_0} = 990 \pm 20$ MeV and $\Gamma = 40\text{--}100$ MeV, those of $f_0(500)$ are still far from consensus (see Ref. [3]). The upper bound of the $f_0(500)$ mass is given in the large- N_c limit in terms of the Gasser–Leutwyler low-energy constant [30], which suggests that the f_0 mass is quite possibly smaller than 700 MeV.

While there has been a great deal of theoretical work on the structure of the $f_0(500)$, its reaction mechanism has been less investigated. Recently, the CLAS Collaboration has reported the first analysis of S -wave photoproduction of $\pi^+\pi^-$ pairs in the region of the $f_0(980)$ at photon energies between 3.0 and 3.8 GeV and momentum transfer squared $-t$ between 0.4 GeV² and 1 GeV² [31,32]. While the differential cross section for the $\gamma p \rightarrow \pi^+\pi^-p$ process in the S -wave shows an evident signal for the $f_0(980)$ production, the $f_0(500)$ was not seen clearly. However, there is still a hint to the existence of the $f_0(500)$ in $\pi^+\pi^-p$ photoproduction measured under different kinematic conditions [33]. Thus, it is of great interest to study the $\gamma p \rightarrow \pi^+\pi^-p$ reaction in the scalar and isoscalar channels. Since these two pions are strongly correlated, one has to consider the rescattering effects of these two pions to describe $\gamma p \rightarrow \pi^+\pi^-p$ in the scalar and isoscalar channels, which are essential in order to explain the production mechanism of the scalar and isoscalar mesons f_0 quantitatively in $\pi^+\pi^-p$ photoproduction. Moreover, it is crucial to take into account the $K\bar{K}$ channel in addition [34], since its threshold is open in the vicinity of the $f_0(980)$ mass. In order to take into account the effects of the $K\bar{K}$ channel, one has to introduce the coupled-channel formalism, which requires fully coupled $\pi\pi$ and $K\bar{K}$ amplitudes.

However, before we carry out the investigation on the $\gamma p \rightarrow \pi^+\pi^-p$ reaction, we need to examine the related two-body process $\gamma p \rightarrow f_0 p$ as a first step toward more complicated correlated $\pi\pi$ photoproduction. Moreover, since the CLAS Collaboration has already presented the differential cross section for $f_0(980)$ photoproduction in the photon energy range $E_\gamma = (3.0\text{--}3.8)$ GeV, it is important to study f_0 photoproduction theoretically as well before we examine the $\gamma p \rightarrow \pi^+\pi^-p$ process with pion pairs in the S -wave. In addition to $f_0(980)$ photoproduction, we study in the present work the $f_0(500)$ production by photon beams, based on effective Lagrangians and a Regge approach. The Regge exchange in $f_0(980)$ photoproduction has already been applied in Ref. [35] with the same Regge trajectory but a different set of parameters. We will first compute the differential cross section for $f_0(980)$ photoproduction and compare the results with the CLAS experimental data such that we can fix the parameters for the t -channel Reggeon exchange. Since there is no excited nucleon that decays into $\pi^+\pi^-$ pairs in the S -wave beyond the $f_0(980)N$ threshold, we consider only the N exchange in the s channel. Then we will proceed to study $f_0(500)$ photoproduction with the same parameters for the Reggeon, which is fixed in $f_0(980)$ production. As far as $f_0(500)$ photoproduction is concerned, we need to consider several excited nucleons above the threshold energy, which can decay into $\pi^+\pi^-N$, where the pion pairs are in the isoscalar and scalar waves.

Upon computing the transition amplitude for the $f_0(500)$ photoproduction, there are ambiguities to which we have to pay close attention. Firstly, the width of the $f_0(500)$ is very large, so that the value of the $f_0(500)$ mass is quite uncertain. Thus, we have to look into the dependence of the results on the $f_0(500)$ mass. Secondly, little experimental information on the photocoupling constant $g_{\gamma\sigma\rho}$

exists. Though there are several theoretical suggestions about its value, an agreement has not yet been reached. Experimentally, two relevant decay channels are known: $\rho^0 \rightarrow \pi^+\pi^-\gamma$ [36,37] and $\rho^0 \rightarrow \pi^0\pi^0\gamma$ [38,39]. To see the contribution of the $f_0(500)$ in these decay processes, one has to rely on models. So, it is necessary to examine uncertainties arising from the coupling constant $g_{\gamma\sigma\rho}$. In principle, ω -meson exchange could be considered. However, the branching ratio of $\omega \rightarrow \pi^+\pi^-\gamma$ is not much known: its experimental upper bound is given as $< 3.6 \times 10^{-3}$ with a confidence level of 95 %. On the other hand, $\rho \rightarrow \pi^+\pi^-\gamma$ is experimentally known to be $(9.9 \pm 1.6) \times 10^{-3}$ [3]. Thus, the value of the $\gamma f_0 \rho$ coupling constant is expected to be much larger than that of $\omega \rightarrow f_0 \gamma$, based on the experimental data given above. We have confirmed numerically that the effect of ω exchange is indeed much smaller than that of ρ exchange. So, we will ignore in this work the contribution from ω -meson exchange. Thirdly, the final state in the $N^* \rightarrow (\pi\pi)_{S\text{-wave}}^{I=0} N$ decay should contain both the background $\pi\pi$ and the $f_0(500)$ resonance. This indicates that it is rather difficult to determine the coupling constants for the $f_0(500)NN^*$ unambiguously. Considering these points, which will bring about the uncertainties of the present work, we have to introduce certain assumptions before we proceed to investigate $f_0(500)$ photoproduction. Though we will take 500 MeV as a main value for the $f_0(500)$ mass in this work, we will carefully examine the dependence of the results for the total cross section on the mass of the lowest-lying scalar meson. Lastly, we will regard $\pi^+\pi^-$ pairs in the S -wave as the $f_0(500)$ meson, which are produced in the course of $N^* \rightarrow (\pi^+\pi^-)_{S\text{-wave}}^{I=0} N$ decays, so that we are able to determine the strong coupling constants for the $N^* \rightarrow f_0(500)N$ transitions. Since the $f_0(500)$ resonance is the most dominant one in $\pi\pi$ scattering in the scalar–isoscalar channel, this approximation is rather plausible.

In addition to the Roper resonance, we want to consider other N^* resonances that can decay into $(\pi\pi)_{S\text{-wave}}^{I=0} N$. Referring to Ref. [3], we find that 10 excited nucleon resonances have a decay channel of $N^* \rightarrow (\pi\pi)_{S\text{-wave}}^{I=0} N$. However, there are not enough data for $N^* \rightarrow (\pi\pi)_{S\text{-wave}}^{I=0}$ except for $N(1440)$, $N(1680)$, and $N(1880)$. Since $N(1880)$ has an overall status of 2 stars, we will not consider it (see the review “ N and Δ resonances” in Ref. [3]). Thus, we expect that the main contribution will come from $N(1440)$ and $N(1680)$. As we will see later, $N(1680)$ provides predominantly a large contribution to the total cross section. Thus, we will set up two different models to delve into each contribution from the N^* resonances. In Model I, we will include those with spin 1/2 and an overall status of 3 or 4 stars. Thus, we consider $N(1535)1/2^-$, $N(1650)1/2^-$, and $N(1710)1/2^+$ in addition to $N(1440)1/2^+$, though their data are not much known. In Model II, we further take into account $N(1520)3/2^-$, $N(1675)5/2^-$, and $N(1680)5/2^+$, together with those included in Model I.

The present work is structured as follows: In Sect. 2, we explain the general formalism for the $f_0(980)$ and $f_0(500)$ photoproductions. In Sect. 3, we present the numerical results separately for Model I and Model II, and discuss their physical implications. The last section is devoted to a summary and conclusion of the present work. We also discuss perspectives for future work in that section.

2. General formalism

We start with the tree-level Feynman diagrams relevant to the $\gamma p \rightarrow f_0(980)p$ and $\gamma p \rightarrow f_0(500)p$ reactions based on the effective Lagrangian approach, as depicted in Fig. 1. Note that we will not consider any contribution from N^* resonances to $f_0(980)$ photoproduction, since the threshold energy of its production is rather high. There is no N^* above 1900 MeV, which can decay into $f_0(980)N$. Thus, we only take into account nucleon exchange both in the s and u channels. As far as the t -channel diagram is concerned, ρ -meson exchange comes into play. However, the $\gamma N \rightarrow f_0(980)N$ reaction

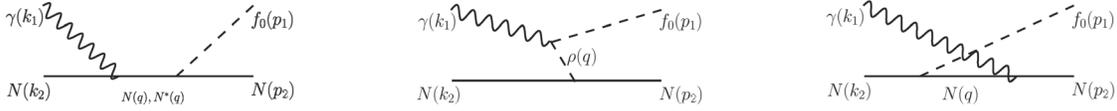


Fig. 1. Feynman diagrams for f_0 photoproduction at the tree level. Each diagram corresponds to the s channel, the t channel, and the u channel in order.

has been experimentally measured at higher photon energies ($E_\gamma = 3.0\text{--}3.8$ GeV) [31,32], which is quite far beyond the threshold. Since the method of effective Lagrangians is devised for describing the mechanism of hadronic reactions near the threshold, we need to revise it to explain $f_0(980)$ photoproduction at higher E_γ . We will employ a hybridized Regge approach in which the Feynman propagator in the t channel is replaced with the Regge propagator for ρ -meson exchange, while the coupling constants and spin structures are taken from the effective Lagrangians. This approach was successfully used for describing the productions of strange and charmed hadrons [40–42]. A virtue of using this hybridized Regge model is that we can also use the reggeized ρ meson in the t channel for $f_0(500)$ photoproduction with the same parameters fixed in $f_0(980)$ photoproduction.

As for the $\gamma N \rightarrow f_0(500)N$ reaction, we include the ρ Reggeon in the t channel with the same parameters used in $f_0(980)$ photoproduction. In the s and u channels, we consider nucleon exchange. Since the threshold energy of $f_0(500)$ photoproduction is about 1.4–1.5 GeV, we need to introduce in the s channel the N^* resonances that decay only into $(\pi\pi)_{S\text{-wave}}^{I=0}N$. To study the contributions of the N^* resonances, we develop two different models: In Model I, we include $N(1440)1/2^+$, $N(1535)1/2^-$, $N(1650)1/2^-$, and $N(1710)1/2^+$, while Model II further contains $N(1520)3/2^-$, $N(1675)5/2^-$, and $N(1680)5/2^+$ with higher spins in addition to those in Model I. The kinematics of the $\gamma p \rightarrow f_0 p$ reactions are given as shown in Fig. 1: k_1 and k_2 stand for the four-momenta, respectively, for the photon and nucleon in the initial state, whereas p_1 and p_2 designate those, respectively, for the f_0 and the nucleon in the final state.

2.1. Model I

As we have already briefly mentioned in the introduction, there are ten excited nucleon resonances that can decay into $f_0(500)N$. However, if we consider N^* with an overall status of 3 or 4 stars, then there are only seven N^* resonances. In order to scrutinize the effects of the N^* resonances carefully, we first introduce the pertinent N^* resonances only with spin 1/2 to describe $f_0(500)$ photoproduction. We call this *Model I*.

To compute the Feynman-invariant amplitudes for f_0 photoproduction, we use the following effective Lagrangians [43–46]:

$$\begin{aligned}
\mathcal{L}_{\gamma f_0 \rho} &= \frac{eg_{\gamma f_0 \rho}}{2m_\rho} \partial_\mu A_\nu (\partial^\mu \rho^\nu - \partial^\nu \rho^\mu) f_0, \\
\mathcal{L}_{\gamma NN} &= -\bar{N} \left(e_N \gamma_\mu A^\mu - \frac{e\kappa_N}{2m_N} \sigma_{\mu\nu} \partial^\nu A^\mu \right) N, \\
\mathcal{L}_{\rho NN} &= -g_{\rho NN} \bar{N} \left(\gamma_\mu \rho^\mu - \frac{\kappa_\rho}{2m_N} \sigma_{\mu\nu} \partial^\nu \rho^\mu \right) N, \\
\mathcal{L}_{f_0 NN} &= g_{f_0 NN} f_0 \bar{N} N, \\
\mathcal{L}_{\gamma NN^* \left(\frac{1}{2}^\pm\right)} &= \frac{ef_1}{2m_N} \bar{N} \Gamma^{(\mp)} \sigma_{\mu\nu} \partial^\nu A^\mu N^* + \text{H.c.}, \\
\mathcal{L}_{f_0 NN^* \left(\frac{1}{2}^\pm\right)} &= \pm g_{f_0 NN^*} f_0 \bar{N} \Gamma^{(\mp)} N^* + \text{H.c.},
\end{aligned} \tag{1}$$

where A , N , ρ , and f_0 denote the photon, the nucleon, the $\rho(770, 1^-)$, and the f_0 fields, respectively. N^* represents a field for the excited nucleon. The values of the coupling constants given in the Lagrangians will be discussed later. The matrix $\Gamma^{(\pm)}$ depends on the parity of the N^* resonance and is defined as

$$\Gamma^{(+)} = \gamma_5, \quad \Gamma^{(-)} = \mathbf{1}. \quad (2)$$

Based on the effective Lagrangians in Eq. (1), we can compute the Feynman-invariant amplitudes for each channel as follows:

$$\begin{aligned} -i\mathcal{M}_{t(\rho)} &= \frac{ieg_{\gamma f_0 \rho} g_{\rho NN}}{2m_\rho} \bar{u}(p_2) \frac{1}{t - m_\rho^2} \left[\epsilon^\alpha (k_1 \cdot p_1) - k_1^\alpha (\epsilon \cdot p_1) \right] \left[\gamma_\alpha - \frac{i\kappa_\rho}{2m_N} \sigma_{\alpha\beta} q_t^\beta \right] u(k_2), \\ -i\mathcal{M}_{s(N)} &= ig_{f_0 NN} \bar{u}(p_2) \frac{q_s + m_N}{s - m_N^2} \not{\epsilon} \left[e_N - \frac{e\kappa_N}{2m_N} \not{k}_1 \right] u(k_2), \\ -i\mathcal{M}_{u(N)} &= ig_{f_0 NN} \bar{u}(p_2) \not{\epsilon} \left[e_N - \frac{e\kappa_N}{2m_N} \not{k}_1 \right] \frac{q_u + m_N}{u - m_N^2} u(k_2), \end{aligned} \quad (3)$$

where the momentum transfers are given as $q_t = p_1 - k_1$, $q_s = k_1 + k_2$, and $q_u = p_2 - k_1$. e_N is the electric charge of the nucleon and ϵ_μ indicates the polarization vector of the photon. $u(p_2)$ and $u(k_2)$ denote the Dirac spinors of the outgoing nucleon and the incoming nucleon, respectively. m_N and m_ρ stand for the masses of the nucleon and ρ meson, respectively. t , s , and u represent the Mandelstam variables and are defined as

$$t = (p_1 - k_1)^2, \quad s = (k_1 + k_2)^2, \quad u = (p_2 - k_1)^2. \quad (4)$$

The effective Lagrangian approach has been successfully used to describe hadronic reactions in the low-energy region. However, when the photon energy increases, the results from the effective Lagrangians start to deviate from the data and do not even satisfy the unitarity [41]. Since the CLAS data on $f_0(980)$ photoproduction were obtained at $E_\gamma = 3.0$ GeV and $E_\gamma = 3.8$ GeV, which is far from the threshold, the effective Lagrangian method is not suitable to explain them. Thus, we employ a hybridized Regge model, as already mentioned. In this approach, the Reggeon in the t channel is governed by the Regge trajectory of the ρ meson, which is well known already [47–52]. We replace the Feynman propagator with the Regge one \mathcal{P}_ρ^R in the t channel [47,51,52]:

$$\frac{1}{t - m_\rho^2} \Rightarrow \mathcal{P}_\rho^R(s, t), \quad (5)$$

where the Regge propagator is defined as

$$\mathcal{P}_\rho^R(s, t) = \left(\frac{s}{s_\rho} \right)^{\alpha_\rho(t)-1} \frac{1}{\sin[\pi\alpha_\rho(t)] \Gamma[\alpha_\rho(t)]}. \quad (6)$$

Here, $\alpha_\rho(t)$ denotes the Regge trajectory for the ρ meson. s_ρ indicates the energy scale parameter for the corresponding Reggeon and is set to be equal to 1.0 GeV^2 . The Regge trajectory for the t channel is taken from Ref. [52]:

$$\alpha_\rho(t) = 0.55 + 0.8t. \quad (7)$$

We adopt a degenerate propagator [52,53] with a constant phase (1). A rotating phase ($e^{-i\pi\alpha(t)}$) is also acceptable. Although the nondegenerate phase $((1 - \exp[-i\pi\alpha(t)])/2)$ is general for the

Table I. Spin-1/2 resonances [3] for $\gamma N \rightarrow f_0(500)N$ photoproduction in Model 1.

Particle	J^P	Mass [GeV]	Width [GeV]	Status
$N(1440)$	$\frac{1}{2}^+$	1.440	0.300	****
$N(1535)$	$\frac{1}{2}^-$	1.535	0.150	****
$N(1650)$	$\frac{1}{2}^-$	1.655	0.150	****
$N(1710)$	$\frac{1}{2}^+$	1.710	0.100	***

ρ trajectory [54], it is excluded from our consideration because the corresponding result for $d\sigma/dt$ reveals a dip structure near $-t \simeq 0.7 \text{ GeV}^2$ in the range $E_\gamma = (3.0\text{--}3.8) \text{ GeV}$, which definitely deviates from the experimental data [31].

In addition, we introduce the scaling factor for the ρ -meson Reggeon exchange:

$$C(t) = \frac{a_\rho}{(1 - t/\Lambda_\rho^2)^2}, \quad (8)$$

which is often included to explain the experimental values of the cross sections. However, we find that the results are not sensitive to the parameters a_ρ and Λ_ρ of the scaling factor.

While $f_0(980)$ photoproduction does not require any contribution from the N^* resonances because of its high threshold energy, we need to consider them for the explanation of the $\gamma p \rightarrow f_0(500)p$ reaction. As mentioned earlier, we will take into account the spin-1/2 resonances in Model I: $N(1440)1/2^+$, $N(1535)1/2^-$, $N(1650)1/2^-$, and $N(1710)1/2^+$, as listed in Table I. Since all of them have rather large widths, we will include the finite width in each propagator for N^* . Then, the Feynman amplitudes are derived as

$$\begin{aligned} -i\mathcal{M}_s(N^*(\frac{1}{2}^+)) &= \frac{ef_1 g_{f_0 NN^*}}{2m_N} \bar{u}(p_2) \frac{i(q_s + m_{N^*})}{s - (m_{N^*} - \frac{i}{2}\Gamma_{N^*})^2} \not{k}_1 \not{u}(k_2), \\ -i\mathcal{M}_s(N^*(\frac{1}{2}^-)) &= \frac{ef_1 g_{f_0 NN^*}}{2m_N} \bar{u}(p_2) \frac{i(-q_s + m_{N^*})}{s - (m_{N^*} - \frac{i}{2}\Gamma_{N^*})^2} \not{k}_1 \not{u}(k_2), \end{aligned} \quad (9)$$

where the first term is the amplitude for N^* with positive parity whereas the second one represents that for negative-parity N^* . The coupling constants f_1 and $g_{f_0 NN^*}$ denote, respectively, the generic photocouplings and strong coupling constants for the corresponding N^* resonances. m_{N^*} and Γ_{N^*} represent the corresponding masses and the decay widths of N^* , respectively. Note that, for the propagators in Eq. (9), we consider the pole positions of the N^* resonances in the complex plane.

2.2. Model II

In Model II, we additionally consider the N^* resonances with spins 3/2 and 5/2 in the s channel. We include $N(1520)3/2^-$, $N(1675)5/2^-$, and $N(1680)5/2^+$ in the s channel, as listed in Table II [3]. The relevant effective Lagrangians for the N^* resonances are given as follows:

$$\begin{aligned} \mathcal{L}_{\gamma NN^*(\frac{3}{2}^-)} &= \frac{ief_1}{2m_N} \bar{N}_\mu^* F^{\mu\nu} \gamma_\nu N - \frac{ef_2}{(2m_N)^2} \bar{N}_\mu^* F^{\mu\nu} \partial_\nu N + \text{H.c.}, \\ \mathcal{L}_{\gamma NN^*(\frac{5}{2}^\pm)} &= \pm \frac{ef_1}{(2m_N)^2} \bar{N}_{\mu\alpha}^* \partial^\alpha F^{\mu\nu} \Gamma_\nu^{(\mp)} N \pm \frac{ief_2}{(2m_N)^3} \bar{N}_{\mu\alpha}^* \partial^\alpha F^{\mu\nu} \Gamma^{(\mp)} \partial_\nu N + \text{H.c.}, \\ \mathcal{L}_{f_0 NN^*(\frac{3}{2}^-)} &= \frac{g_{f_0 NN^*}}{m_{f_0}} (\partial_\mu f_0) \bar{N} N^{*\mu} + \text{H.c.}, \\ \mathcal{L}_{f_0 NN^*(\frac{5}{2}^\pm)} &= i \frac{g_{f_0 NN^*}}{m_{f_0}^2} \bar{N} (\partial_\mu \partial_\nu f_0) \Gamma^{(\mp)} N^{*\mu\nu} + \text{H.c.}, \end{aligned} \quad (10)$$

Table II. The excited nucleon resonances [3] for $f_0(500)$ photoproduction in Model II.

Particle	J^P	Mass [GeV]	Width [GeV]	Status
$N(1520)$	$\frac{3}{2}^-$	1.525	0.115	****
$N(1675)$	$\frac{3}{2}^-$	1.675	0.150	****
$N(1680)$	$\frac{5}{2}^+$	1.685	0.130	****

where f_1, f_2 , and $g_{f_0 NN^*}$ denote the photocouplings and the strong coupling constants, respectively. They can be determined by using the experimental data on the photon decay amplitudes and the decay widths $\Gamma_{N^* \rightarrow f_0(500)N}$. Then, the Feynman-invariant amplitudes for the s channel are derived as

$$\begin{aligned}
-i\mathcal{M}_{s(N(1520))} &= ip_1^\mu \bar{u}(p_2) \Delta_{\mu\alpha}(q_s, m_{N^*}) \left(k_1^\alpha \epsilon^\beta - k_1^\beta \epsilon^\alpha \right) \\
&\quad \times \left(\frac{ef_1 g_{f_0 NN^*}}{2m_N m_{f_0}} \gamma_\beta + \frac{ef_2 g_{f_0 NN^*}}{4m_N^2 m_{f_0}} k_{2\beta} \right) u(k_2), \\
-i\mathcal{M}_{s(N(1675))} &= \bar{u}(p_2) p_1^\rho p_1^\sigma \gamma_5 \Delta_{\rho\sigma; \mu\alpha}(q_s, m_{N^*}) k_1^\alpha \left(k_1^\mu \epsilon^\nu - k_1^\nu \epsilon^\mu \right) \\
&\quad \times \left(\frac{ief_1 g_{f_0 NN^*}}{4m_N^2 m_{f_0}^2} \gamma_\nu + \frac{ief_2 g_{f_0 NN^*}}{8m_N^3 m_{f_0}^2} k_{2\nu} \right) \gamma_5 u(k_2), \\
-i\mathcal{M}_{s(N(1680))} &= -\bar{u}(p_2) p_1^\rho p_1^\sigma \Delta_{\rho\sigma; \mu\alpha}(q_s, m_{N^*}) k_1^\alpha \left(k_1^\mu \epsilon^\nu - k_1^\nu \epsilon^\mu \right) \\
&\quad \times \left(\frac{ief_1 g_{f_0 NN^*}}{4m_N^2 m_{f_0}^2} \gamma_\nu + \frac{ief_2 g_{f_0 NN^*}}{8m_N^3 m_{f_0}^2} k_{2\nu} \right) u(k_2), \tag{11}
\end{aligned}$$

where $\Delta_{\mu\alpha}$ and $\Delta_{\rho\sigma; \mu\alpha}$ indicate the Rarita–Schwinger propagators for the N^* resonances with spin 3/2 and 5/2, respectively, defined as [55–58]

$$\begin{aligned}
\Delta_{\mu\alpha}(q, m_{N^*}) &= \frac{i(q + m_{N^*})}{s - (m_{N^*} - \frac{i}{2}\Gamma_{N^*})^2} S_{\mu\alpha}(q, m_{N^*}), \\
\Delta_{\alpha\beta; \mu\nu}(q, m_{N^*}) &= \frac{i(q + m_{N^*})}{s - (m_{N^*} - \frac{i}{2}\Gamma_{N^*})^2} S_{\alpha\beta; \mu\nu}(q, m_{N^*}). \tag{12}
\end{aligned}$$

Here, $S_{\mu\alpha}$, $\bar{g}_{\mu\alpha}$, and $\bar{\gamma}_\mu$ are expressed as

$$\begin{aligned}
S_{\mu\alpha}(q, m_{N^*}) &= -\bar{g}_{\mu\alpha} + \frac{1}{3} \bar{\gamma}_\mu \bar{\gamma}_\alpha, \\
S_{\alpha\beta; \mu\nu}(q, m_{N^*}) &= \frac{1}{2} (\bar{g}_{\alpha\mu} \bar{g}_{\beta\nu} + \bar{g}_{\alpha\nu} \bar{g}_{\beta\mu}) - \frac{1}{5} \bar{g}_{\alpha\beta} \bar{g}_{\mu\nu} \\
&\quad - \frac{1}{10} (\bar{\gamma}_\alpha \bar{\gamma}_\mu \bar{g}_{\beta\nu} + \bar{\gamma}_\alpha \bar{\gamma}_\nu \bar{g}_{\beta\mu} + \bar{\gamma}_\beta \bar{\gamma}_\mu \bar{g}_{\alpha\nu} + \bar{\gamma}_\beta \bar{\gamma}_\nu \bar{g}_{\alpha\mu}), \\
\bar{g}_{\mu\alpha} &= g_{\mu\alpha} - \frac{q_\mu q_\alpha}{m_{N^*}^2}, \quad \bar{\gamma}_\mu = \gamma_\mu - \frac{q_\mu \not{q}}{m_{N^*}^2}. \tag{13}
\end{aligned}$$

2.3. Parameters and form factors

The coupling constant for the ρNN vertex is the most important parameter to describe $f_0(980)$ photoproduction, since the t channel governs the production mechanism of the $\gamma N \rightarrow f_0(980)N$ reaction. The coupling constant $g_{\rho NN}$ and κ_ρ are well known from NN potentials. For example,

$g_{\rho NN} = 3.25$ was used in the full Bonn potential [59], while the Nijmegen group employed $g_{\rho NN} = 2.76$ [60]. On the other hand, smaller coupling constants $g_{\rho NN} = 2.6$ and $\kappa_\rho = 3.7$ have been exploited in Regge models for photoproduction of the pion and of the charged ρ^\pm meson [61,62]. We use $g_{\rho NN} = 3.25$ and the ratio of the vector and tensor couplings $\kappa_\rho = 6.1$ [59].

The $f_0(980)NN$ coupling constant is taken from Ref. [63]: $g_{f_0(980)NN} = 5.8$. The $f_0(500)NN$ coupling constant $g_{f_0(500)NN}$ is a crucial one that explains the mid-range strong attraction in the NN interaction. Its value is given in a wide numerical range. For example, the full Bonn potential suggests that $g_{\sigma' NN} = 8.46$ [59] with $m_{\sigma'} = 550$ MeV. The notation σ' in Ref. [59] was introduced to emphasize the fact that σ' represents only an effective description of correlated 2π exchange in the S -wave. In the one-boson-exchange (OBE) Bonn NN potential, two different σ were introduced, i.e., σ_1 in the isovector ($T = 1$) channel and σ_2 in the isoscalar ($T = 0$) channel. The charge-dependent OBE Bonn NN potential even has several different values for different partial waves in NN scattering [64]. One has to keep in mind that the coupling constant for the σNN vertex describes effectively whole 2π and even $\pi\rho$ exchanges. On the other hand, the Nijmegen soft-core (NSC) potential includes two different scalar mesons, flavor SU(3) symmetry and ideal mixing being used [65]. Depending on whether the scalar mesons constitute quark–antiquark pairs or tetraquarks, Ref. [65] suggested two different values for the σNN and $f_0(980)NN$ coupling constants with different flavor content. Since the $f_0(500)$ in the present work corresponds to the S -wave correlated 2π , we will take the value of $g_{\sigma' NN}$ from the full Bonn potential, i.e., $g_{f_0(500)NN} = 8.46$. This choice is reasonable, since $f_0(500)$ in the present work indeed corresponds to the resonance in the S -wave correlated 2π channel. Note that σ' exchange in the full Bonn potential was later replaced by the explicit S -wave correlated 2π exchange [66,67]. Of course, there is one caveat: Since the $f_0(500)$ meson has a very broad width, a single coupling constant is a rather crude approximation. A more complete work considering $f_0(500)$ as an S -wave correlated 2π resonance will be considered in future work.

The mass of the $f_0(500)$ meson brings out another ambiguity in dealing with $f_0(500)$ photoproduction. Because of the broad width of $f_0(500)$ ($\Gamma = (400\text{--}700)$ MeV), it is rather difficult to determine its mass exactly. The Particle Data Group estimated the pole mass of $f_0(500)$ to be $(400\text{--}500 - i(200\text{--}350))$ MeV [3]. This implies that we need to examine carefully the $f_0(500)$ mass dependence of the observables. If it is taken to be larger than $m_{f_0(500)} = 500$ MeV, the threshold energy can be larger than the mass of $N(1440)$. Thus, the total cross section of $\gamma N \rightarrow f_0(500)N$ could reveal different behavior from that with the lower value of the $f_0(500)$ mass. We will discuss in detail the dependence of the total cross section on $m_{f_0(500)}$ in the next section.

In order to determine the strong coupling constants for the excited baryons, we assume that the $f_0(500)$ meson consists mainly of the correlated 2π state in the S -wave. Then, we can use the decay modes of $N^* \rightarrow (\pi^+\pi^-)_{S\text{-wave}}^{I=0}$ to fix the corresponding coupling constants by using the partial decay widths defined as

$$\begin{aligned}\Gamma\left(N^*\left(\frac{1}{2}^\pm\right) \rightarrow f_0 N\right) &= \frac{g_{f_0 NN^*}^2}{4\pi} \frac{|\mathbf{p}|}{m_{N^*}} (E_N \pm m_N), \\ \Gamma\left(N^*\left(\frac{3}{2}^\pm\right) \rightarrow f_0 N\right) &= \frac{1}{3} \frac{g_{f_0 NN^*}^2}{4\pi} \frac{|\mathbf{p}|^3}{m_{N^*} m_{f_0}^2} (E_N \mp m_N), \\ \Gamma\left(N^*\left(\frac{5}{2}^\pm\right) \rightarrow f_0 N\right) &= \frac{2}{15} \frac{g_{f_0 NN^*}^2}{4\pi} \frac{|\mathbf{p}|^5}{m_{N^*} m_{f_0}^4} (E_N \pm m_N),\end{aligned}\tag{14}$$

Table III. The strong coupling constants for the N^* resonances. The first four coupling constants correspond to those for the excited nucleons with spin 1/2, whereas the next three are those for the N^* resonances with spin 3/2 and 5/2.

$g_{f_0 NN(1440)}$	$g_{f_0 NN(1535)}$	$g_{f_0 NN(1650)}$	$g_{f_0 NN(1710)}$	$g_{f_0 NN(1520)}$	$g_{f_0 NN(1675)}$	$g_{f_0 NN(1680)}$
± 3.88	± 2.55	± 0.96	± 0.15	± 0.85	± 9.85	± 2.29

where the partial decay width $\Gamma(N^* \rightarrow f_0 N)$ is given as $\Gamma_{N^*}^{\text{BW}} \times \text{Br}(N^* \rightarrow f_0 N)$. $\Gamma_{N^*}^{\text{BW}}$ denotes the Breit–Wigner total decay width and $\text{Br}(N^* \rightarrow f_0 N)$ the branching ratio. $|\mathbf{p}|$ represents the magnitude of the final-state momentum defined as

$$|\mathbf{p}| = \frac{1}{2m_{N^*}} \sqrt{[m_{N^*}^2 - (m_N + m_{f_0})^2][m_{N^*}^2 - (m_N - m_{f_0})^2]}. \quad (15)$$

The results of the strong coupling constants are shown in Table III.

Concerning the photocoupling constants for ρ -meson exchange, we use $g_{\gamma f_0(500)\rho} = 0.25$ [38] from the measurement of the SND Collaboration, and $g_{\gamma f_0(980)\rho} \approx 0.21$ [68,69] from the molecular $K\bar{K}$ model for $\Gamma_{f_0 \rightarrow \rho\gamma}$ [70]. On the other hand, we can utilize the helicity amplitudes given by the Particle Data Group [3] to find the photocoupling constants f_1 and f_2 for excited nucleons. Since the helicity amplitude are expressed as [71]

$$\begin{aligned} A_{1/2} \left(\frac{1^\pm}{2} \right) &= \mp \frac{ef_1}{2m_N} \sqrt{\frac{k_\gamma m_{N^*}}{m_N}}, \\ A_{1/2} \left(\frac{3^\pm}{2} \right) &= \mp \frac{e\sqrt{6}}{12} \sqrt{\frac{k_\gamma}{m_N m_{N^*}}} \left[f_1 + \frac{f_2}{4m_N^2} m_{N^*} (m_{N^*} \mp m_N) \right], \\ A_{3/2} \left(\frac{3^\pm}{2} \right) &= \mp \frac{e\sqrt{2}}{4m_N} \sqrt{\frac{k_\gamma m_{N^*}}{m_N}} \left[f_1 \mp \frac{f_2}{4m_N} (m_{N^*} \mp m_N) \right], \\ A_{1/2} \left(\frac{5^\pm}{2} \right) &= \pm \frac{e}{4\sqrt{10}} \frac{k_\gamma}{m_N} \sqrt{\frac{k_\gamma}{m_N m_{N^*}}} \left[f_1 + \frac{f_2}{4m_N^2} m_{N^*} (m_{N^*} \pm m_N) \right], \\ A_{3/2} \left(\frac{5^\pm}{2} \right) &= \pm \frac{e}{4\sqrt{5}} \frac{k_\gamma}{m_N^2} \sqrt{\frac{k_\gamma m_{N^*}}{m_N}} \left[f_1 \pm \frac{f_2}{4m_N} (m_{N^*} \pm m_N) \right], \end{aligned} \quad (16)$$

where k_γ stands for the photon decay momentum in the center-of-mass (CM) frame and is expressed as

$$k_\gamma = \frac{m_{N^*}^2 - m_N^2}{2m_{N^*}}, \quad (17)$$

we obtain the photocouplings for the electromagnetic transitions $N^* \rightarrow \gamma N$ as listed in Table IV.

A hadron has a spatial size, which can be characterized by phenomenological form factors. Hence, one has to introduce them at each baryon–baryon–meson vertex. Note that each amplitude of N exchanges does not satisfy the gauge invariance in the s and u channels, but the sum does. To restore the gauge invariance, we modify the Born scattering amplitudes as

$$\mathcal{M}_{\text{Born}} = \mathcal{M}_t^{\text{Regge}} + (\mathcal{M}_s + \mathcal{M}_u) F_c^2(s, u), \quad (18)$$

Table IV. The photocoupling constants for the N^* resonances with spin 1/2 are listed in the first row. In the second and third rows the photocoupling constants for the N^* resonances with spin 3/2 and 5/2, respectively, are listed.

$f_{\gamma NN(1440)}$ 0.47	$f_{\gamma NN(1535)}$ 0.81	$f_{\gamma NN(1650)}$ 0.28	$f_{\gamma NN(1710)}$ -0.24
$f_{1\gamma NN(1520)}$ 4.63	$f_{2\gamma NN(1520)}$ -4.92		
$f_{1\gamma NN(1675)}$ -1.34	$f_{2\gamma NN(1675)}$ -2.28	$f_{1\gamma NN(1680)}$ 15.24	$f_{2\gamma NN(1680)}$ -13.47

where a common form factor is introduced as [72]

$$F_c(s, u) = F(s) + F(u) - F(s)F(u), \quad (19)$$

with

$$F(q^2) = \frac{\Lambda^4}{\Lambda^4 + (q^2 - m_{\text{ex}}^2)^2}. \quad (20)$$

Here, q indicates the off-shell four-momentum for an exchanged hadron, m_{ex} stands for its mass, and Λ denotes a cutoff parameter. As for the form factors in N^* exchange, we also use Eq. (20). However, it is not sufficient to preserve the unitarity for the N^* resonances with spin 3/2 and 5/2, since the corresponding amplitudes show much stronger q^2 dependence so that they are divergent as E_γ increases. Thus, in order to tame the divergence, we will employ a Gaussian type of form factor for $N(1520)3/2^-$, $N(1675)5/2^-$, and $N(1680)5/2^+$ exchange in the s channel, defined as

$$F(q^2) = \exp\left[-\frac{(q^2 - m_{N^*}^2)^2}{\Lambda^4}\right]. \quad (21)$$

3. Results and discussion

We are now in a position to present the numerical results of our work. Since there exist experimental data on the differential cross section of the $\gamma N \rightarrow f_0(980)N$ reaction [31], we start with $f_0(980)$ photoproduction.

3.1. $\gamma N \rightarrow f_0(980)N$

In order to describe the experimental data on $d\sigma/dt$ of the $\gamma N \rightarrow f_0(980)N$ process [31], we fix the scaling factor in Eq. (8) to be $a_\rho = 9.0$. As is well known, while the Regge approach describes the energy dependence very well, it cannot determine the absolute magnitude of the cross sections. Thus, it is inevitable to introduce the scale parameter a_ρ to fit the cross sections. The value of the cutoff parameter Λ_ρ is selected to be 1 GeV, which is a typical order for the cutoff value. We do not fit it to avoid additional ambiguity. The cutoff parameters in the N and N^* exchanges are chosen as $\Lambda_N = 0.8$ GeV and $\Lambda_{N^*} = 1.0$ GeV. The values of Λ_{ρ, N^*} are also typical ones. Note that the cutoff value for the backward-scattering amplitude is usually smaller than that for the forward-scattering one in the literature [44,45]. Furthermore, the experimental data for the backward region are currently not available. So, we simply choose these values without any fitting procedure.

Figure 2 shows the differential cross section $d\sigma/dt$ of the $\gamma N \rightarrow f_0(980)N$ reaction. The experimental data are taken from Ref. [31], where $d\sigma/dt$ were measured within the range of the photon

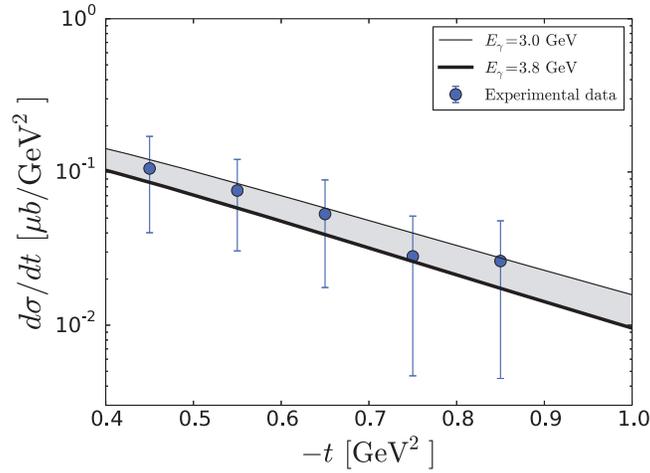


Fig. 2. Differential cross section of the $\gamma N \rightarrow f_0(980)N$ reaction as a function of t in the range of $E_\gamma = (3.0\text{--}3.8)$ GeV. The shaded band represents the present result. The experimental data are taken from Ref. [31].

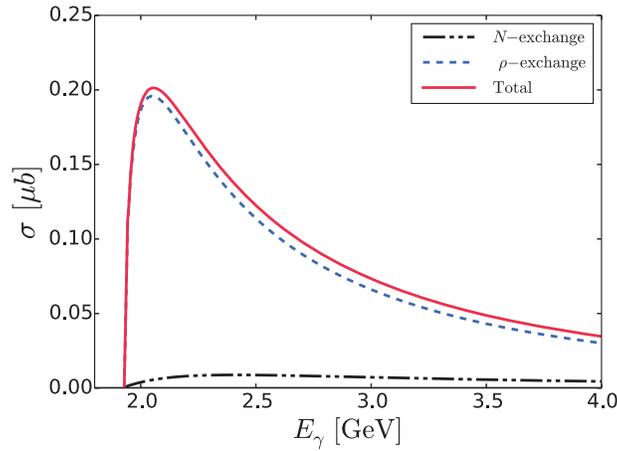


Fig. 3. Total cross section of the $\gamma N \rightarrow f_0(980)N$ reaction. The dashed curve depicts the contribution of ρ -Reggeon exchange in the t channel, whereas the dot-dashed and double-dot-dashed ones show those of N -Reggeon exchange and of N exchange in the u and s channels, respectively. The solid curve represents the total contribution.

energy $E_\gamma = (3.0\text{--}3.8)$ GeV. In order to compare the present results with the data, we present the results as a shaded band, the width of which represents the corresponding region of E_γ . Considering the large experimental uncertainty, the results describe the data very well. Note that the t dependence is governed by ρ -Reggeon exchange in the t channel.

In Fig. 3, we depict the contribution of each channel to the total cross section of the $\gamma N \rightarrow f_0(980)N$ reaction. The ρ -Reggeon exchange dominates over the s - and u -channel diagrams, whereas the N exchanges have small effects in the whole energy region.

3.2. $\gamma N \rightarrow f_0(500)N$

The parameters for $f_0(500)$ photoproduction in the t , s , and u channels are kept the same as those in the case of the $\gamma N \rightarrow f_0(980)N$ reaction. However, the N^* resonances play essential roles in describing the $\gamma N \rightarrow f_0(500)N$ reaction, in particular in the vicinity of the threshold.

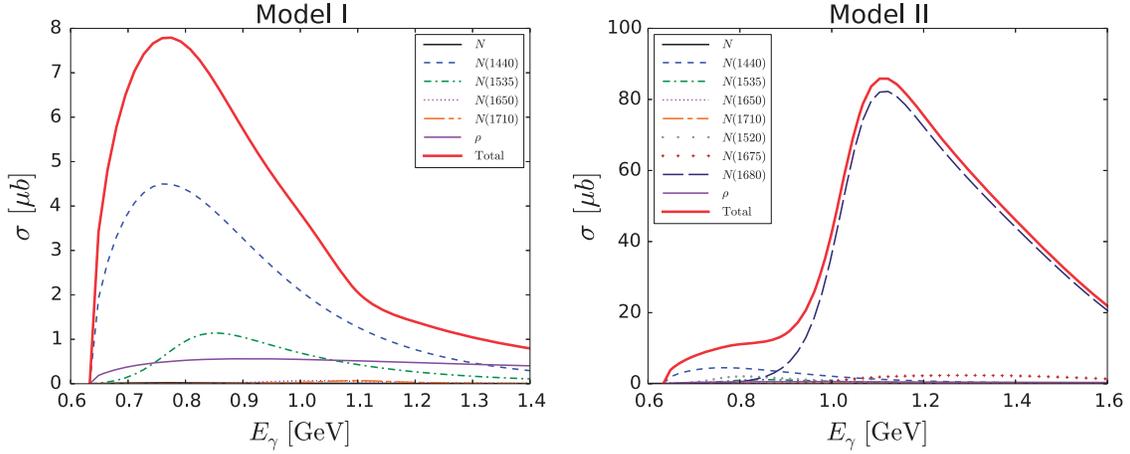


Fig. 4. Total cross sections for the $\gamma N \rightarrow f_0(500)N$ reaction as a function of E_γ . In the left panel, the results from Model I are shown, whereas in the right panel those from Model II are depicted. Each contribution is distinguished by different types of curves.

We need to delve into the physical reasons for Model I. The $N(1440)1/2^+$ and $N(1535)1/2^-$ resonances increase the total cross section of the $\gamma N \rightarrow f_0(500)N$ reaction near the threshold. As shown in the left panel of Fig. 4, the most dominant contribution arises from the $N(1440)1/2^+$ resonance. Since the $N(1650)1/2^-$ and $N(1710)1/2^+$ resonances have relatively smaller strong coupling constants as well as photocouplings, as listed in Tables III and IV, the effect of these resonances is tiny. Interestingly, the contribution of the ρ Reggeon in the t channel is rather suppressed. The effect of N exchanges in the s and u channels is much smaller than that of ρ -Reggeon exchange through the whole energy region. The magnitude of the total cross section for $f_0(500)$ photoproduction is about 40 times larger than that for $f_0(980)$ production in the case of Model I.

The right panel of Fig. 4 depicts the results of the total cross section for the $\gamma N \rightarrow f_0(500)N$ reaction from Model II, where the N^* resonances with higher spins are added to the s channel. The $N(1680)5/2^+$ resonance in the s channel yields a remarkably large contribution to the total cross section of $f_0(500)$ photoproduction, so that its magnitude reaches even about $80 \mu b$ around $E_\gamma \approx 1.1$ GeV. There is at least one clear reason for this large contribution of $N(1680)5/2^+$. Firstly, the photocouplings of $N(1680)5/2^+$ are very large, as shown in Table IV, which come from the large values of the experimental data on the photon decay amplitudes $A_{1/2}$ and $A_{3/2}$ [3]. The value of $f_{1\gamma NN(1680)}$ is even about 32 times larger than that of $f_{1\gamma NN(1440)}$. Moreover, the size of the strong coupling constant for the $f_0(500)NN(1680)$ vertex is comparable to that of $g_{f_0(500)NN(1440)}$. Note that even though the value of $g_{f_0(500)NN(1675)}$ is rather large, the $N(1675)5/2^-$ resonance has almost no effect on the total cross section because of its negative parity. The results are not so sensitive to variations of the cutoff masses. In this regard, the contribution of the $N(1680)5/2^+$ resonance discussed here seems to be robust. Thus, it would indeed be of great interest if one could justify experimentally whether the $N(1680)5/2^+$ resonance plays such a dominant role in describing the $\gamma N \rightarrow f_0(500)N$ reaction.

As mentioned briefly in the introduction, the uncertainty in the mass of $f_0(500)$ is so large on account of its broad width, so it is quite difficult to settle the threshold energy. The PDG data have it that the pole mass of $f_0(500)$ is $[(400-550)-i(200-350)]$ MeV [3]. In fact, the estimated mass of $f_0(500)$ is given by a wide range of values, as listed in Ref. [3]. Thus, we have to examine the dependence of the total cross section on the mass of $f_0(500)$. In the left panel of Fig. 5, we show

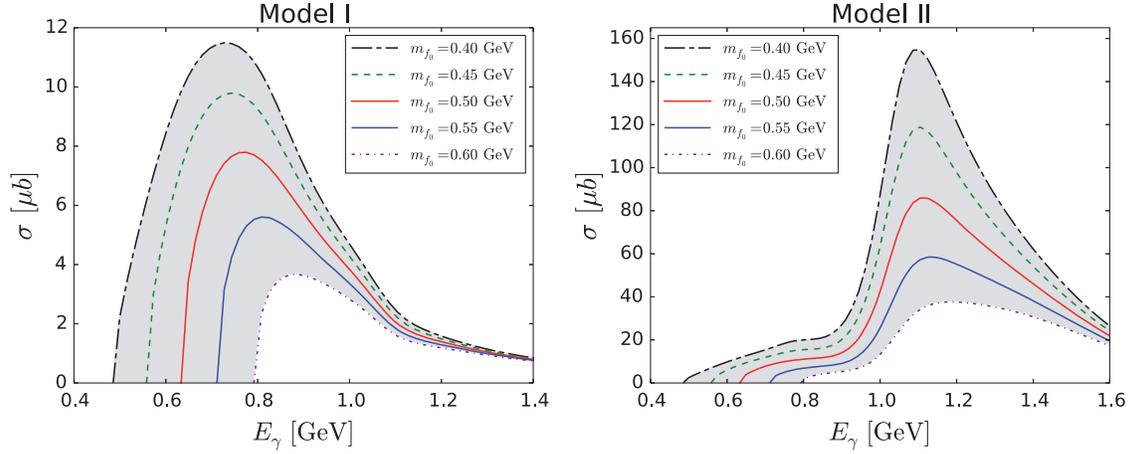


Fig. 5. $f_0(500)$ mass dependence of the total cross section for the $\gamma N \rightarrow f_0(500)N$ reaction both in Model I and Model II.

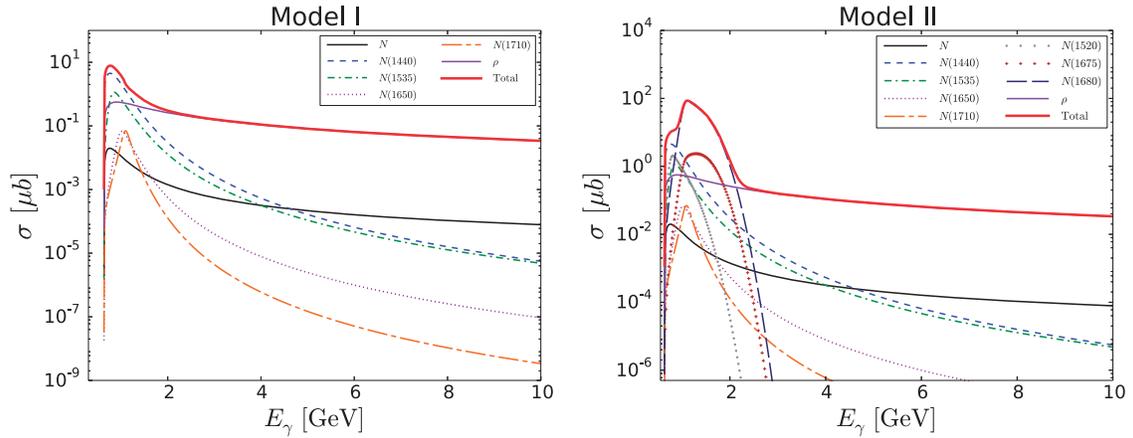


Fig. 6. Total cross sections for the $\gamma N \rightarrow f_0(500)N$ reaction as a function of E_γ in the log scale. In the left panel, the results from Model I are shown, whereas in the right panel those from Model II are depicted. Each contribution is distinguished by different types of curves.

the results of the total cross section of the $\gamma N \rightarrow f_0(500)N$ reaction from Model I with various values of m_{f_0} given between 0.4 GeV and 0.6 GeV. As expected, smaller values of m_{f_0} produce larger magnitudes of the total cross section. Note that, if one uses a value of m_{f_0} larger than 500 MeV, $N(1440)1/2^+$ will be excluded because of the larger threshold energy. Thus, the total cross section starts to be reduced when the value of m_{f_0} is larger than 500 MeV. In the present work, we will take $m_{f_0(500)} = 500$ MeV from now on.

In Figs. 4 and 5, we have mainly examined the total cross section of $f_0(500)$ photoproduction in the vicinity of the threshold energy. We now delve into the dependence of the total cross section on E_γ from the threshold energy through 10 GeV in the log scale. In the left panel of Fig. 6, we show the behavior of each contribution from Model I as E_γ increases. As expected, all the resonance effects are diminished quickly with increasing photon energy, while the ρ -Reggeon in the t channel takes over the contributions of all the N^* resonances around $E_\gamma \approx 2$ GeV and then dictates the dependence of the total cross section on E_γ . The t -channel Reggeon ensures the unitarity of the total

cross section, as shown in Fig. 6. In the limit of $s \rightarrow \infty$, the unpolarized sum of the Regge amplitude complies with the following asymptotic behavior:

$$\lim_{s \rightarrow \infty} \sum_{\text{pol}} |\mathcal{M}_t^{\text{Regge}}(s, t)|^2 \propto s^{2\alpha_\rho(t)}. \quad (22)$$

Thus, if one further increases E_γ , the contribution of the ρ meson starts to decrease, satisfying the asymptotic behavior of Eq. (22). The right panel of Fig. 6 shows the results from Model II. Similarly, the contributions of most N^* resonances fall off very fast as E_γ increases. The effect of $N(1680)5/2^+$ lessens continuously after $E_\gamma \approx 1.2$ GeV, and then becomes smaller than those of the ρ Reggeons around $E_\gamma \approx 2$ GeV. Thus, the N^* resonances come into play only near the threshold region as anticipated.

In the upper panel of Fig. 7, the results of the differential cross section $d\sigma/d \cos \theta$ from Model I are plotted as functions of $\cos \theta$, as the photon energy E_γ is varied from 0.8 GeV through 3.5 GeV. Usually, the s -channel contributions including all N^* resonances with spin 1/2 do not show any $\cos \theta$ dependence. Note that Model I does not contain any N^* resonances with higher spins. Nevertheless, the results of $d\sigma/d \cos \theta$ exhibit different peculiarities. The differential cross section $d\sigma/d \cos \theta$ in the forward region grows as E_γ increases. However, the value of $d\sigma/d \cos \theta$ almost vanishes at a very forward angle at higher values of E_γ . This arises from the structure of the amplitude of ρ exchange in Eq. (3). Model II yields rather different results from those based on Model I. Since the $N(1680)5/2^+$ resonance is the most dominant one from the threshold energy through 1.5 GeV, as shown in Fig. 4, and it has spin 5/2 with positive parity, we expect that it will have a certain effect on the $\cos \theta$ dependence of the differential cross section. Indeed, it steers $d\sigma/d \cos \theta$ up to $E_\gamma \approx 2.1$ GeV, as shown in the lower panel of Fig. 7. As E_γ increases to more than 1.7 GeV, the ρ Reggeon gains control of the $\cos \theta$ dependence of the differential cross section, so that we have more or less the same results as in the upper panel of Fig. 7.

Figure 8 shows the results of the differential cross section $d\sigma/dt$ as a function of $-t$ in the range of $E_\gamma = (3.0\text{--}3.8)$ GeV, so that we can directly compare them with those of $f_0(980)$ photoproduction. The t dependence is almost the same as that of the $f_0(980)$ case, because we have exactly the same ρ -Reggeon and N exchanges. The only differences come from the coupling constants. The contributions of the N^* resonances are all suppressed in this region of the photon energies. The magnitude of $d\sigma/dt$ is approximately 2 times larger than that of $f_0(980)$ photoproduction on account of different coupling constants. Note that the differential cross section $d\sigma/dt$ obeys the following asymptotic behavior:

$$\lim_{s \rightarrow \infty} \frac{d\sigma}{dt}(t \rightarrow 0) \propto s^{2\alpha_\rho(0)-2} \quad (23)$$

and the result shown in Fig. 8 satisfies Eq. (23).

4. Summary and conclusion

We aimed in this work to investigate $f_0(500)$ and $f_0(980)$ photoproduction, based on a hybridized Regge model. We first described the differential cross section $d\sigma/dt$ for the $\gamma N \rightarrow f_0(980)N$ reaction, compared with the recent experimental data on it. We fixed the relevant Regge parameters by reproducing the data. We introduced the N^* contribution in the s channel to study the production mechanism of the $\gamma N \rightarrow f_0(500)N$ reaction. Since its threshold energy is much smaller than $f_0(980)$ photoproduction, there exist several N^* resonances that can decay into $(\pi\pi)_{S\text{-wave}}^{I=0} N$. Assuming that

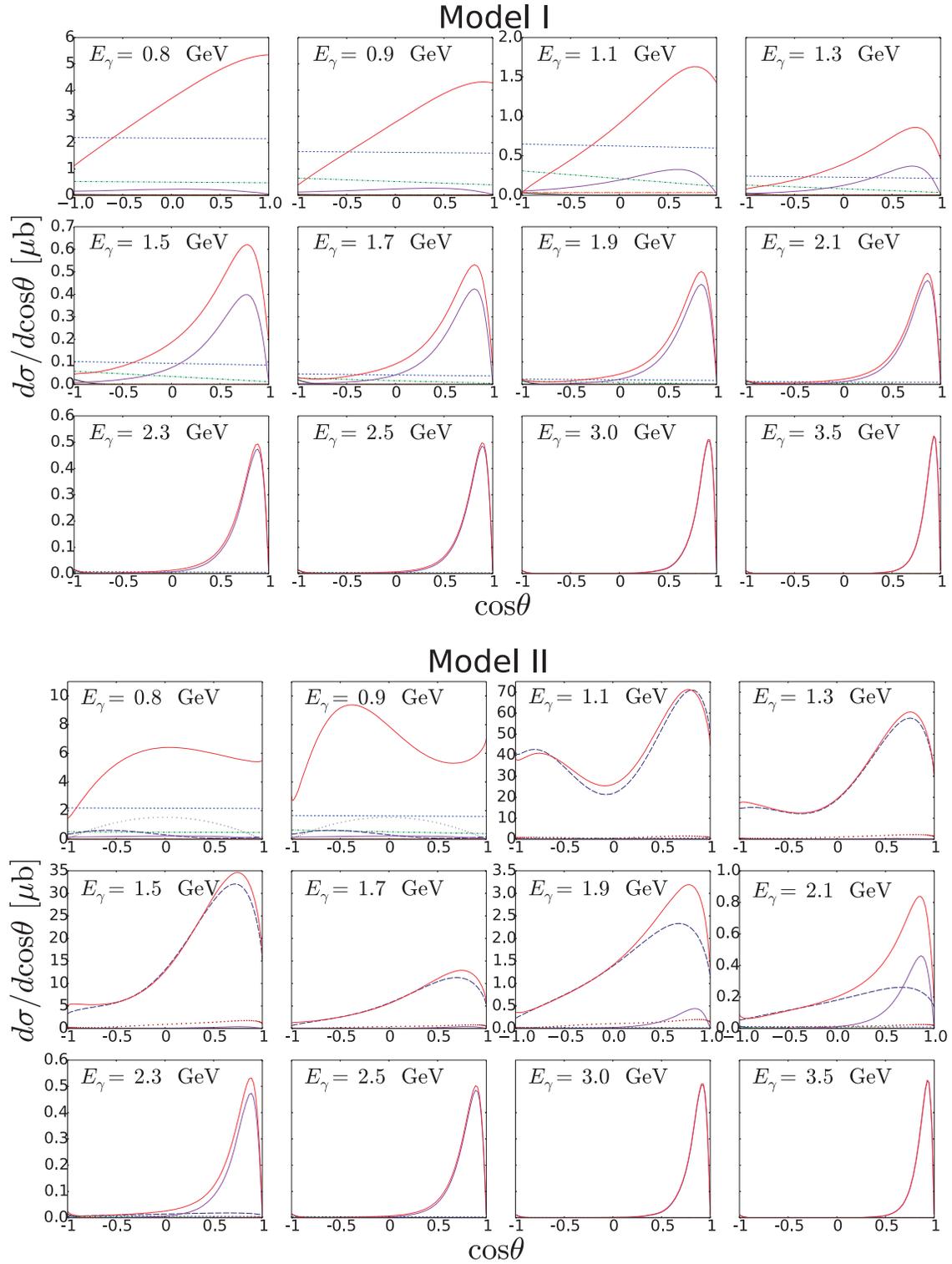


Fig. 7. Differential cross section for the $\gamma N \rightarrow f_0(500)N$ reaction as a function of $\cos\theta$ with the photon energy E_γ changing from 0.8 GeV to 3.5 GeV. The results from Model I are shown in the upper panel, whereas those from Model II are depicted in the lower panel. Notations are the same as in Fig. 4.

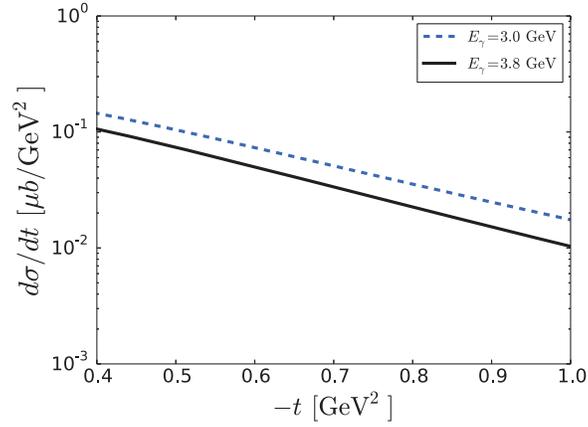


Fig. 8. Differential cross section for the $\gamma N \rightarrow f_0(500)N$ reaction as a function of t at two different photon energies, i.e., $E_\gamma = 3.0$ GeV and $E_\gamma = 3.8$ GeV. The solid curve represents the result at $E_\gamma = 3.8$ GeV, whereas the dashed one plots that at $E_\gamma = 3.0$ GeV.

$f_0(500)$ is much stronger than the background of the $(\pi\pi)_{S\text{-wave}}^{I=0}$ channel, we were able to find the strong coupling constants for the $f_0(500)NN^*$ vertices. The photocouplings of the N^* resonances were determined by using the experimental data on the corresponding photon decay amplitudes. The cutoff masses for the form factors were fixed at 1.0 GeV to avoid additional ambiguity. In dealing with these N^* resonances, we constructed Model I and Model II. Model I included those with spin 1/2 only, while Model II was built in such a way that more N^* resonances with higher spins were added to Model I. Near the threshold, we found that $N(1440)1/2^+$ and $N(1535)1/2^-$ were dominant in Model I, whereas the effects of other N^* resonances were almost negligible. In Model II, the contribution of $N(1680)5/2^+$ dictates the total cross section of the $\gamma N \rightarrow f_0(500)N$ reaction. Remarkably, the $N(1680)5/2^+$ resonance enhances the magnitude of the total cross section up to about 80 μb . The main reason for this comes from the large value of its photocouplings. The strong coupling constant of $N(1680)5/2^+$ is also relatively large. Since the mass of the $f_0(500)$ meson is not precisely fixed because of its large width, we examined the dependence of the total cross section on its mass in the range of (400–600) MeV. As expected, the smaller the $f_0(500)$ mass was, the larger the total cross section was in Model I and Model II. If E_γ increases, then the N^* contribution fades away very fast, so that ρ -Reggeon exchange takes control, as in the case of $f_0(980)$ photoproduction. We also computed the differential cross section $d\sigma/d\cos\theta$. While the contributions of the N^* resonances in the s channel are rather flat in the case of Model I, $N(1680)5/2^+$ governs $\cos\theta$ dependence again because of its high spin. Finally, we computed the differential cross section $d\sigma/dt$ for the $\gamma N \rightarrow f_0(500)N$ reaction. The results showed that the t dependence and the magnitude looked very similar to those for $f_0(980)$ photoproduction.

Though it is very difficult to study $f_0(500)$ photoproduction experimentally, it is still of great importance to study the production mechanism of the $\gamma N \rightarrow f_0(500)N$ reaction, since it casts light on the structure of the N^* resonances as investigated in the present work. It also provides a certain clue to studying more complicated processes with three-particle final states, such as $\gamma N \rightarrow (\pi\pi)_{S\text{-wave}}^{I=0} N$, in the future.

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