THE GRIBOV–MORRISON RULE AS A KINEMATIC EFFECT

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For diffractive excitation of $I = 1/2$ nucleon resonances N* by pions, $N+\pi \to N^*+\pi$, the Gribov–Morrison rule is shown to follow from kinematics alone if t-channel helicity conservation is exact. Under this condition the production of unnatural-parity N* is strongly suppressed by a mass factor. A similar result for diffractive $N+N \to N^*+N$ follows from factorization. However, no corresponding effect can be found for diffractive excitation of mesons. The experimental situation is discussed in the light of this effect, which we conclude must play an important role in any realistic model for diffractive $N \to N^*$ transitions, at least beyond the low N* mass region.

1. *Introduction*

In this article we shall prove that for diffractive excitation of nucleons by pions, $N+\pi \to N^*+\pi$, the assumption of t-channel helicity conservation leads automatically to the Gribov–Morrison rule [1, 2]. Under the conditions assumed this rule is nothing more than a kinematic effect. The additional assumption of factorization leads to a corresponding result for diffractive excitation by nucleons, $N+N \to N^*+N$. Experimentally it is known the diffractive NN* vertex has a very strong component which conserves t-channel helicity, at least when the N* mass is not too small. Therefore our result appears to clarify the success of the Gribov–Morrison rule for diffractive excitation of nucleons in the N* region where this rule works best.

For diffractive excitation of pions and kaons the assumption of TCHC does not lead to a spin-parity selection rule, because of different kinematics. Even though TCHC is a characteristic of some (but not all) diffractive transitions $\pi \to \pi^*$ and $K \to K^*$ we find no reason to expect the Gribov–Morrison rule to be valid in this case. The present experimental situation suggests it is not.

The Gribov–Morrison (GM) rule [1, 2] for a diffractive transition $N \to N^*$ is the statement the nucleon, with natural parity, cannot be diffractively excited to any N* state with unnatural parity.It is well-known this rule cannot be derived from any generally-accepted dynamical principle, and yet it seems to be valid. Unnatural-parity resonances

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do not contribute strongly in diffractive $N\pi \to N^*\pi$ and $NN \to N^*N$. More precisely, natural-parity states appear to dominate the peaks which are observed at N* masses 1500, 1690 and 2190 MeV. Presumably some contributions from unnatural-parity states are in these peaks, but they cannot be very large. Another enhancement, usually associated with the $P_{11}(1470)$, is observed at about 1.4 GeV. This identification satisfies the GM rule. But the resonance interpretation of this enhancement is doubtful, and one should not use it to base any conclusion concerning spin-parity selection rules.

In Section 2 we find our result for $N\pi \to N^*\pi$ using the following elementary argument. TCHC is assumed for $N \to N^*$, leaving one independent amplitude. We discard all invariant amplitudes which break TCHC, retaining the one (there is only one) which does not. Calculating the nonzero t-channel helicity amplitude in terms of this invariant amplitude we find a strong kinematic suppression from a mass factor when the N* has unnatural parity. This effect weakens with increasing $|t|$. Nevertheless it can lead to a factor of, typically, 1/20 in the cross section for unnatural-parity N* production relative to the natural-parity cross section. An effect of this strength is quite enough to explain the success of the GM rule in diffractive excitation of nucleons.

The simplicity and generality of this result leave little to be desired. However, TCHC is not an exact symmetry of nature and one may ask, is this effect likely to be relevant for the construction of a quantitative model? Above N* mass ~ 1.6 GeV, where TCHC seems to become nearly exact [3–5], the answer is surely yes. For smaller N* mass, where TCHC is not exact, the answer is less clear. Here the Drell–Hiida–Deck effect [6, 7] seems to play a very important role, and resonance production is difficult to isolate, so that one cannot be sure whether TCHC and the GM rule are characteristic of resonance production or not.

Concerning TCHC breaking we can make another comment. More detailed investigations show that strong TCHC breaking *can* be arranged in such a way that the GM effect arising from TCHC is not weakened at all. TCHC breaking can also weaken the effect, of course. But it is difficult to arrange things in such a way the effect is entirely cancelled, within any realistic model.

Diffractive $NN \to N^*N$ is discussed in Section 3. Exact TCHC at the NN* vertex still leaves four independent amplitudes, and if all of them are important there is no kinematic suppression of unnatural-parity N* production by protons. However, factorization seems to hold for inelastic diffractive scattering in the range $0 < |t| < 0.5\ \text{GeV}^2$ [29], and perhaps also for larger $|t|$. Assuming factorization is valid, the TCHC amplitudes for $NN \to N^*N$ contain the same mass factor which leads to the GM effect in $N\pi \to N^*\pi$.

In Section 3 we also discuss diffractive excitation of mesons, and the evidence for and against the GM rule in these reactions.

2. $N+\pi \to N^*+\pi$

Consider diffractive excitation of $I = 1/2$ nucleon resonances N* by pions, $N\pi \to N^*\pi$. The N* spin is J and its parity is arbitrary. One can also think of N* as a component of a diffractively-produced background with definite angular momentum. We assume exact

TCHC so just one amplitude is independent. In the t-channel $Db \to cA$ (where $c = \mathrm{N}$, $A = \overline{\mathrm{N}}$, $D = \pi$ and $b = \pi$) the helicity amplitudes are

$$f^{(t)}_{cAD0}(\pm) = \Psi_{c\nu_1 \dots \nu_{J-1/2}}(p_c)\Gamma(\pm)M^{\nu_1 \dots \nu_{J-1/2}}v_A(p_A), \tag{1}$$

where

$$\Gamma(+) \equiv 1,\ \Gamma(-) \equiv \gamma_5, \tag{2}$$

and the label (+), (−) means the N* has natural, unnatural parity. In general the M-fuction in Eq. (1) contains $2J+1$ invariant amplitudes. However, $2J$ of these break TCHC. Only one does not, namely

$$M_{\nu_1 \dots \nu_{J-1/2}} = (p_A \cdots p_A)_{\nu_1 \dots \nu_{J-1/2}} A(\pm). \tag{3}$$

To prove this statement one has to calculate the t-channel helicity coefficients of all $2J+1$ invariant amplitudes. Although we have done this calculation, the results are too long to reproduce here. The interested reader can, of course, verify Eq. (3) without too much difficulty using standard methods. The problem is simplified by the fact the t-channel spin structure is triangular. In other words, the $2J$ spin-flip amplitudes depend on the other $2J$ invariant amplitudes but *not* on $A(\pm)$.

Standard formulae [8, 9] for the spin-J wave-function in Eq. (1) enable one to calculate these helicity amplitudes quite easily. They are

$$f^{(t)}_{cAD0}(\pm) = \delta_{cA} C_J (m_c \pm m_a) \left(\frac{T_{ca}}{m_c\sqrt{2}}\right)^{J-1/2} [1 - t/(m_c \pm m_a)^2]^{1/2} A(\pm), \tag{4}$$

where

$$C_J = [(J+\tfrac{1}{2})!(J-\tfrac{1}{2})!/(2J)!]^{1/2},$$

$$T^2_{ca} = t^2 - 2t(m_c^2 + m_a^2) + (m_c^2 - m_a^2)^2.$$

In the amplitude (4) there is an overall factor $(m_c + m_a)$, $(m_c - m_a)$ for natural, unnatural parity N^*, where m_c is the N^* mass and m_a is the nucleon mass. For $m_c \sim 1.5$ GeV these factors make the amplitude for natural parity-production about five times larger than the one for unnatural-parity production, other things being equal. Here is the kinematic effect promised in the Introduction. Of course, this effect is t-dependent, as one sees from the ratio of cross sections

$$\frac{d\sigma(-)}{d\sigma(+)} = R(t)\frac{|A(-)|^2}{|A(+)|^2}, \tag{5}$$

where

$$R(t) = \frac{(m_c - m_a)^2}{(m_c + m_a)^2}\,\frac{1 - t/(m_c - m_a)^2}{1 - t/(m_c + m_a)^2}. \tag{6}$$

For large $|t|$ the function (6) is $R(t) \approx 1$. But the low $|t|$ region is much more important, and in this region $R(t)$ is rather small. Indeed, for $m_c \sim 1.5$ GeV, $R(t) \sim 1/25$ at $t = 0$, and even though this function increases with $|t|$ it is only $\sim 1/12$ at $|t| \sim 1$ GeV².

The factor $R(t)$ approaches one in the limit $m_c \to \infty$, but the approach to this limit is not rapid. Thus for N* mass $m_c \sim 2$ GeV one still has $R(0) \sim 1/9$ and $R(|t| = 1) \sim 1/5$. Of course, for large enough N* mass the kinematic suppression of unnatural-parity N* production will disappear, whether or not TCHC holds in this mass region. This is an interesting prediction.

So far we have said nothing about the invariant amplitudes $A(\pm)$. These amplitudes are defined in such a way they are subject to no kinematic constraint, and contain no kinematic singularity, whether or not TCHC is imposed. Therefore one must take the mass factor $(m_c \pm m_a)$ in Eq. (4) seriously; by construction, invariant amplitudes are minimally dependent on external masses and mass differences. It is, of course, possible that $A(-)$ is larger than $A(+)$ for some dynamical reason. However, it would have to be several times larger to cancel the mass factor effect, and this seems quite improbable. The amplitudes $A(\pm)$ are proportional to the covariant couplings of the pomeron at the NN* and pion vertices. There is no reason why the unnatural-parity covariant couplings should be several times larger than the natural-parity covariant couplings.

Perhaps it is worth emphasizing that amplitudes of the form (4) will emerge in *any* covariant theory which is constrained to satisfy TCHC. A good example is the covariant Regge model (see e. g. [25]). Imposing exact TCHC leaves only one covariant coupling at the NN* vertex and the mass factor $(m_c \pm m_a)$ emerges automatically.

We have chosen to discuss TCHC in terms of t-channel amplitudes. However, one of the main approaches to diffractive excitation puts the emphasis on s-channel amplitudes (see e. g. Kane [28]). In this approach TCHC arises "accidentally" when one arranges that all s-channel helicity amplitudes be large away from $t = 0$. (In our covariant notation the corresponding statement is the invariant amplitude $A(\pm)$ in the M-function (3)contributes to *every* s-channel helicity amplitude, as one can easily verify.) Needless to say, the spin-parity effect, which arises from TCHC alone, exists whether or not TCHC is "accidental".

One should know where the mass factor $(m_c \pm m_a)$ originates. It comes from the matrix element

$$\bar{u}_c(p_c)\Gamma(\pm)v_A(p_A) = -\delta_{cA}(\pm)^{1/2+A}(m_c \pm m_a)\,(1-t/(m_c \pm m_a)^2)^{1/2} \tag{7}$$

which also appears in t-channel amplitudes for NN → N*N and for N* electroproduction Ne → N*e. Therefore one can expect similar results for these processes. However, for diffractive excitation of pions and kaons there is no corresponding result. In fact, there is no kinematic preference for or against a change in naturality in this case. The connection between TCHC and the GM rule holds only for nucleon excitation.

Now let us sketch the experimental situation in diffractive Nπ → N*π. Numerous experiments have shown there are peaks in the N* mass spectrum which can be identified with the natural-parity states $D_{13}(1520)$, $F_{15}(1690)$ and $G_{17}(2190)$ [3–5, 10, 11]. While these resonances may dominate the peaks, smaller contributions from unnatural-parity resonances are not excluded. An enhancement at 1.4 GeV is also seen, and for some years this enhancement was tentatively identified with the $P_{11}(1470)$ in spite of the large mass difference. However, this identification may be incorrect. Recent work [12, 19] based on the Drell-Hiida-Deck effect [6, 7] has been quite successful in explaining N*(1400) production in NN collisions. This is surely relevant here because NN → N*(1400)N and Nπ → N*(1400) π closely resemble one another, while both reactions exhibit characteristics which are very different from those of other N* production reactions (e. g. much steeper t-dependence). If Deck type models are realistic then resonance production is not very important in the lower N* mass region, and it is not appropriate to discuss the GM rule in this region. For larger N* mass, on the other hand, a Deck-type model is unlikely to give an important contribution [12]. Therefore we are led to consider diffractive N* production as consisting of two distinct N* mass regions in which different dynamics are important. The N* mass dividing these two regions lies somewhere between 1.5 GeV and 1.6 GeV.

Another reason for dividing diffractive excitation into two mass regions with different dynamics is the fact TCHC is a good approximation in the higher mass region but not in the lower [3–5]. Furthermore, a clear violation of the GM rule has been observed [21, 22] in the lower mass region, namely the presence of a $J^P = 1/2^-$ state. This is just what we expect from the fact that TCHC is strongly broken there. In the higher mass region TCHC seems to be nearly exact [5]. Our interpretation of the GM rule as an effect resulting from TCHC is therefore valid in this region. It is possible, of course, that TCHC and the GM rules are also characteristic of *resonance* production in the low mass region, but a nonresonant contribution prevents this from being seen. The situation may be especially complicated in the intermediate region where the N*(1500) peak is dominant.

Data on diffractive NK → N*K has also been collected [23, 24], and the N* excitation spectra observed is quite similar to the N* spectrum in Nπ → N*π. All of the remarks we have made for πN collisions apply also to KN collisions.

We touched upon the question of TCHC breaking in the Introduction. A careful discussion of this point requires an extensive kinematic analysis, which has been carried out. The main result is that TCHC can be broken in essentially two ways, such that (1) unnatural-parity suppression is retained full strength, (2) or weakened. It depends on which TCHC breaking invariant amplitudes become significantly nonzero. When good models of diffraction dissociation become available this point may be interesting to pursue.

3. Other diffractive processes

(1) For the reaction NN → N*N the assumption of TCHC at the N → N* vertex leaves four independent amplitudes. Two of these (those with an extra factor γ_5 in each vertex) are proportional to $(m_c \mp m_a)$, while the other two are proportional to the factor $(m_c \pm m_a)$ which gives us the GM effect. Clearly if all four amplitudes are important then there is no preference for any N* spin-parity. However, the first two amplitudes violate factorization. Since the latter property seems to be valid [29], these amplitudes can presumably be neglected, and we are left with the two which lead to the GM effect for NN → N*N.

Turning to the data on diffractive NN → N*N [11, 13–18, 26] one is impressed by its similarity with the diffractive Nπ → N*π data. The GM rule seems to work; peaks are seen which correspond to $D_{13}(1520)$, $F_{15}(1690)$ and $G_{17}(2190)$. Clear evidence for contributions to these peaks from N* with unnatural-parity has not been found. An $I = 1/2$ enhancement near 1.4 GeV is present similar to the one seen in $N\pi \to N^*\pi$. It seems probable [12, 19] this is due to the Drell-Hiida-Deck effect, and not to production of the $P_{11}(1470)$. As for $N\pi \to N^*\pi$, strong violation of TCHC is observed for N* mass below 1.6 GeV, while for larger N* mass TCHC is a good approximation [18]. TCHC also seems to hold at Fermilab energies [27]. At ISR energies there is evidence [20] TCHC holds even for very large diffractively produced masses.

As for $N\pi \to N^*\pi$ it seems appropriate to consider the larger and smaller mass regions separately. In the latter, resonance production is presumably not dominant, whereas in the former it is. Assuming factorization we conclude there is a direct connection between TCHC and the GM rule for larger N* masses.

(2) It has already been mentioned that TCHC in meson transitions like $\pi \to \pi^*$, $K \to K^*$ does not imply the GM rule. Therefore, if this rule is valid for diffractive meson excitation, then our argument concerning baryon excitation is weakened. Let us now briefly review the status of the GM rule in $\pi \to \pi^*$, $K \to K^*$.

The transitions $\pi \to A_1, A_3$ and $K \to Q$, L satisfy the GM rule. We have to look for diffractive transitions which violate the rule. A prime candidate is $\pi \to A_2$, which is dominated by isoscalar, natural-parity exchange, and is only weakly energy-dependent [30, 31]. (See however Leith [32] for a possible objection to the pomeron interpretation.) The transition $\pi \to A_2$ does not satisfy TCHC, because of factorization. Indeed, the zero helicity state of the A_2 is absent because of factorization and parity invariance [33, 34]. Therefore, the $\pi \to A_2$ cross section has a dip in the forward direction. This dip is a characteristic of all transitions $\pi \to \pi^*$, $K \to K^*$ which violate the GM rule. A forward dip decreases the cross section, of course. Nevertheless, A_2 production is a well-established resonance production reaction which seems to be diffractive. If it is, then the GM rule is not valid for meson excitation.

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