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A SPACE(-TIME) PICTURE OF HADRONIC DIFFRACTION

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ABSTRACT

Geometry is introduced into particle scattering by means of boundary conditions on the surface of an interaction region (in spirit similar to bag models of single hadrons). A unified picture of elastic scattering, diffraction dissociation and two-body exchange reactions emerges in terms of "bending rays" propagating along geodesics of this surface. The longitudinal extension of the interaction region plays an important role: the longer it is the steeper is the diffraction slope. Qualitative consequences of the picture are discussed. In particular, a natural explanation for the slope-mass correlation in diffraction dissociation emerges.

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The confinement problem, and in particular, the phenomenological successes of bag models, have recently created considerable interest in a space-time description of extended hadrons and their interactions.

The aim of this paper is to present the essential steps of a geometrical space(-time) formulation of (diffractive) interactions of hadrons. On the one hand, this picture has been largely influenced by the ideas, methods and results of bag models ¹⁾⁻³⁾ for single hadrons where geometry enters in the form of boundary conditions on the bag surface. On the other hand, it has been strongly inspired by the classical so-called ⁴⁾⁻¹⁰⁾ "geometrical theory of diffraction", which represents a fairly recent progress in advanced optics and acoustics ($\sim 1954-1970$): a solution for high-frequency diffraction can be given in closed form for an arbitrary shape ⁷⁾ of the scatterer. Diffraction appears as due to ⁴⁾⁻¹⁰⁾ "surface waves" or, in ray language, to ⁷⁾ "bending rays" propagating along a geodesic of the scatterer's surface into the geometrical shadow. In classical optics -- away from the forward directions -- the surface wave contributions to elastic scattering are completely dominated by the well-known geometrical optical contributions of reflected and refracted rays.

Not so in particle physics: a direct adaptation of this theory to particle physics is achieved in this paper. Geometry is introduced in a well-defined way in terms of boundary conditions on a finite size interaction region ("interaction bag"). An important result is: due to the presence of multiparticle production elastic scattering and diffractive dissociative processes appear to be dominated by surface waves. These depend on the geometrical shape of the interaction region -- and on nothing else. Let us emphasize that our approach has obviously very little in common with the conventional "optical model" in nuclear and particle physics.

Altogether a well-defined and physically intuitive subdivision of hadronic scattering into diffractive and non-diffractive phenomena is achieved, together with a qualitative, if not quantitative, description of elastic and inelastic diffraction as well as two-body exchange reactions.

THE ANALOGUE PROBLEM

Consider the collision of two hadrons in their centre-of-mass frame (cms)

The hadrons will have short-range interactions, i.e., interactions confined to a certain finite region of space(-time). The shape of this interaction region is determined by the detailed dynamics of strong interactions; it may, e.g., also depend on the cms energy $\sqrt{s}$. (Presumably it is not just the geometrical overlap of the two Lorentz contracted pancakes!) So far we have been fairly general.

Now we shall make the plausible assumption that the essential features of the hadronic scattering amplitude are preserved in the following reduced geometrical analogue situation: "scattering of a plane wave on a static object with sharp boundary", provided that

- i) the reduction is performed in the cms of the two colliding hadrons;
- ii) the wave number k of the plane wave is identified with the cms momentum k_{cm} of the two hadrons;
- iii) the scatterer has the shape of the genuine interaction region [in the cms ^{*)}], since the geometrical dimensions of the two problems should be the same.

Let us emphasize that, in general, neither the plane wave nor the scatterer should be identified with any of the two initial hadrons. [cf. an example in mechanics: scattering of two spheres with radii r_1, r_2 ; interaction region (trivially) = sphere with radius $r_1 + r_2$; reduced equivalent system: scattering of a mass point on a static sphere with radius $r_1 + r_2$.]

Let us now furthermore imagine that the interaction region -- and hence our analogue scatterer -- has the proportions of an intermediate excited hadronic state of mass $M = \sqrt{s}$ built up by the two colliding hadrons [a resonant fireball or "fire sausage" ²⁾]. Thus a bridge is cast to the results of bag models for excited high mass, high spin hadrons ^{2), 3)}. In fact in a subsequent paper ¹¹⁾ we shall choose the shape of our excited intermediate state (in its "rest frame"), as obtained in the framework of bag models ^{2), 3)}: a "cigar" ("fire sausage") with transverse dimension $R_{\perp} \approx 1$ fermi and longitudinal extension $R_{\parallel} \propto \sqrt{s} R_{\perp}^2$, which approaches a string for $s \rightarrow \infty$. We find it most encouraging that exactly such a shape turns out to give optimal results ¹¹⁾ for our scattering problem.

*) We are aware that the shape of the interaction region is not Lorentz invariant. A manifestly covariant formulation has still to follow.

SURFACE WAVES

The analogue problem -- scattering of a plane wave (wave number k , propagation in z direction) on a body with the dimensions of the interaction region -- is described by the scalar, stationary wave equation .

$$\begin{aligned} \Delta \psi + k^2 \psi &= 0 && \text{outside of the scatterer (i.e. free motion)} \\ \Delta \psi + k^2 N^2 \psi &= 0 && \text{inside the scatterer} \end{aligned} \quad (1)$$

-- where $N = N(k)$ is the complex "refraction index" -- and the boundary conditions:

- i) $\psi(k, \vec{r})$ (or its derivative normal to the surface of the scatterer) is continuous on the surface (2)
- ii) $\psi(k, \vec{r}) \underset{|\vec{r}| \rightarrow \infty}{\approx} e^{+ikz} + f(k, \theta) \frac{e^{+ik|\vec{r}|}}{|\vec{r}|}$ (3)

where $f(k, \theta)$ is the scattering amplitude. In our relativistic case, Eq. (1) may be interpreted as a stationary Klein-Gordon equation with $k = k_{\text{cms}} = \sqrt{E^2 - m^2}$, $E = \sqrt{s}/2$.

The scattering problem is completely specified by a "geometrical" and a "structural" input:

- i) the geometrical input is the shape of the scatterer (= shape of the "interaction bag"). We shall discuss the case of an arbitrary, smooth, convex shape which, however, for simplicity is assumed to be rotationally symmetric around the beam direction (z axis)

Important parameters will be the two radii of curvature of the surface in the point met by a ray of tangential incidence, the transverse radius $R_{\perp}$ and the longitudinal radius of curvature $\rho_{\parallel}$ (at $z=0$);

- ii) the structural input is given in terms of the complex refraction index $N(k)$. In particle physics one is interested in a situation somewhere near the black body limit. This is the limit which maximizes absorption, while minimizing reflection [see, e.g., Ref. 12)]:

$$\text{Re } N \rightarrow 1, \text{Im } N \rightarrow 0 \text{ such that } k \cdot \text{Im } N \rightarrow \infty \quad (4)$$

The solution to our analogue problem is generally given in terms of a partial wave-type expansion. For large kR_1 and $k \cdot p_{||}$ it is badly convergent, and a Sommerfeld-Watson type transformation has to be performed instead [for a complete exposition of what follows, see Refs. 7), 9) and 10), for a summary and additional proofs, see Ref. 13)].

Consider first the solution away from the forward direction, i.e., $\theta \gg (k p_{||})^{-1/3}$:

$$f(k, \theta) \approx f_{\text{geom}} + f_{\text{s.w.}} \quad \text{for } kR_1 \text{ and } k p_{||} \gg 1 \quad (5)$$

f_{geom} constitutes the "background integral" in the Sommerfeld-Watson representation. In ray language f_{geom} contains the well-known geometrical optical contributions of reflected and refracted rays. In the black body limit it is approximately given in terms of reflected rays only.

The second term, $f_{\text{s.w.}}$, is a residue sum from moving poles in the direct channel complex angular momentum ^{*)} plane. The pole positions are asymptotically given by ^{*)}

$$\alpha_n(k) + \frac{1}{2} \underset{k \rightarrow \infty}{\approx} kR_1 + e^{i\pi/3} \left(\frac{k p_{||}}{2} \right)^{1/3} \times \left\{ x_n + O\left(\frac{1}{(k p_{||})^{1/3} \sqrt{N^2(k) - 1}} \right) \right\} \quad (6)$$

for $n = 1, 2, 3, \dots$

where x_n is the n^{th} zero of the Airy function $[A_1(-x_n) = 0: x_1 = 2.33811, x_2 = 4.08795, x_n \approx [(3\pi/2)(n - 1/4)]^{2/3} \text{ for } n \gg 1]$ and $N(k)$ has to obey the "technical" assumption ⁹⁾

$$(k p_{||})^{1/3} \sqrt{N^2(k) - 1} \gg 1. \quad (7)$$

This assumption guarantees firstly the existence of these poles. Furthermore, which is much more important, it implies that for sufficiently large k the pole positions $\alpha_n(k)$ become essentially independent of the refraction index $N(k)$. The same is true for the pole residues $r_n(k)$ near the black body limit. Thus the whole residue sum, $f_{\text{s.w.}}$, is almost entirely determined in terms of the shape of the scatterer, alias the interaction region,

*) For any non-spherical shape the poles are strictly in the complex plane of another variable and the pole positions are a more complicated function of the shape; however, α_n , as given in Eq. (6) and considered as complex poles in the J plane, may be shown ^{7), 13)} to lead to a good approximation for not too big angles θ . The correct formulae may be found in Refs. 7), and will be discussed by us in later papers ^{11), 13)}.

and essentially independent of its inner structure ⁹⁾. This is a very attractive feature.

The Regge-type poles $\alpha_n(k)$ describe resonant states which -- in ray language -- are excited by incoming rays of tangential incidence, i.e., with an impact parameter $b \simeq (J + \frac{1}{2})/k \simeq (\text{Re } \alpha_n(k) + \frac{1}{2})/k \simeq R_\perp$. These "rays" creep ⁷⁾ along longitudinal geodesics of the scatterer's curved surface -- hence the name surface waves. Due to the curvature they suffer an intensity loss by radiating off tangentially into the geometrical shadow:

$$f_{\text{s.w.}} \simeq \sum_{n=1}^{\infty} (\alpha_n + \frac{1}{2}) \tau_n(k) \frac{P_{\alpha_n}(-\cos \theta)}{\sin \pi \alpha_n}$$

$$\simeq - \sum_{n=1}^{\infty} \left[\frac{2\alpha_n + \frac{1}{2}}{\pi \theta} \right]^{1/2} \tau_n(k) e^{i(\text{Re } \alpha_n + \frac{1}{2})\theta + i\frac{\pi}{4}} \times e^{-\text{Im } \alpha_n \cdot \theta}$$

(8)

Surface wave

According to Eqs. (6) and (7) the exponential radiation loss $e^{-\text{Im } \alpha_n(k)\theta}$ depends only ^{7), 13)} on the longitudinal radius of the curvature $\rho_\parallel$: the larger $\rho_\parallel$, i.e., the longer the interaction region, the steeper the angular distribution of the surface wave contribution $f_{\text{s.w.}}$.

For $\theta \gg (k \cdot \rho_\parallel)^{-1/3}$ the residue sum of the surface waves is an asymptotic expansion, which according to Eq. (8) is dominated by the first surface wave ($n=1$) having the smallest $\text{Im } \alpha(k)$.

For $\theta \ll (k \cdot \rho_\parallel)^{-1/3}$, however, all surface waves co-operate ¹³⁾, such that

$$f_{\text{s.w.}} \simeq - \sum_{n=1}^{\infty} (\alpha_n + \frac{1}{2}) \tau_n(k) \frac{P_{\alpha_n}(-\cos \theta)}{\sin \pi \alpha_n} \propto - i \int_{-\infty}^{\infty} d\alpha (\alpha + \frac{1}{2}) H_0^{(1)}(\alpha + \frac{1}{2}) \theta$$

$$\propto \frac{i H_1^{(1)}((\alpha_1(k) + \frac{1}{2})\theta)}{k \cdot R_\perp \cdot \theta} \quad [\tau_n \propto \frac{d\alpha_n}{dn}]$$

(9)

the real part of which diverges at $\theta=0$. This divergence is exactly cancelled by the "background term" f_{geom} , which behaves ¹³⁾ like

$$f_{\text{geom}}(k, \theta) \propto \frac{i H_1^{(2)}(k R_\perp \theta)}{k \cdot R_\perp \theta}$$

(10)

such that the sum produces ^{9), 12)} the well-known imaginary Fraunhofer diffraction for $\theta \ll (k \cdot \rho_\parallel)^{-1/3}$:

$$f(k, \theta) = f_{\text{geom}} + f_{\text{s.w.}} \approx \theta \alpha (k p_{\parallel})^{-1/3} \quad 2i \frac{f_1(k R_{\perp} \theta)}{k R_{\perp} \theta} \quad (11)$$

with $\text{Im } f_{\text{s.w.}} \approx \text{Im } f_{\text{geom}} \approx \frac{1}{2} \text{Im } f$ for $\theta \approx 0$.

Let us now discuss a crucial point: for larger angles $\theta \gtrsim (k p_{\parallel})^{-1/3}$ and any $N(k)$ in classical optics and acoustics, the surface wave contribution $f_{\text{s.w.}}$ is negligible as compared to the geometrical optical contribution f_{geom} . This explains why the surface waves have been discussed theoretically and experimentally ^{12), 14)} so late. The only possible exception is the backward direction ($\theta \approx 180^\circ$), where a subtle enhancement of surface waves may lead to a dominant effect ^{9), 14)}, which is known as the glory effect ^{*)} in scattering of light off rain drops.

In the following we shall argue -- and this is an important point of the paper -- that surface waves give the dominant contribution to diffraction scattering in particle physics and that this is due to the presence of multiparticle production!

THE PICTURE IN THE PRESENCE OF MULTIPARTICLE PRODUCTION

Let us remember: the analogue picture stands for hadronic scattering interpreted as formation and subsequent decay of a highly excited intermediate hadronic state. If we now visualize the enormous instability of this intermediate state we are led naturally to the following identifications in the analogue picture [which go beyond the familiar identifications of multiparticle production with (total) absorption]:

All rays which impinge on the scatterer (i.e., really hit it), "achieve" a breakup of the unstable intermediate state. Thus the rays with impact parameters $b < R_1$ contribute to non-diffractive multiparticle final states. Only those rays which at most touch the surface tangentially, i.e., with $b \gtrsim R_1$, contribute to elastic scattering and diffraction dissociation.

This prescription means that besides absorption, also the rays reflected into angles away from $\theta \approx 0$ [$f_{\text{geom}}(\theta)$] are supposed to contribute to (non-diffractive) multiparticle production: according to the laws of geometrical optics these reflected rays just come from impinging rays with $b < R_1$ (e.g., for a sphere $b = R_1 \cdot \cos \theta/2$). For hadronic diffraction we are then left with the full surface wave contributions [$f_{\text{s.w.}}(\theta)$] at all angles and $f_{\text{geom}}(\theta)$ around $\theta \approx 0$, since both result from rays of near grazing incidence.

*) It may well turn out that the backward peaks in particle physics and in scattering of particles off nuclei ¹⁵⁾ are due to the glory effect.

Let us indicate two steps of the proper mathematical treatment ¹³⁾ here.

- i) The "cut-off" of rays with $b \simeq (\text{Re } J + \frac{1}{2})/k < R_L$ has to be applied to the S matrix element $S(J) = S_{\text{geom}}(J) + S_{\text{s.w.}}(J)$ in the complex J plane, as appearing in the Sommerfeld-Watson representation ⁹⁾ of $f(k, \theta)$ and not to the usual impact parameter amplitude for real b. Notice: the impact parameter of a ray is defined as the real part of the complex point in $(J + \frac{1}{2})/k$, from which a dominant contribution to the path integral arises [saddle point at $(J + \frac{1}{2})/k = R_L \cdot \cos(\theta/2) < R_L$ for $f_{\text{geom}}(k, \theta)$ and surface pole positions with $(\text{Re } \alpha_n(k) + \frac{1}{2})/k \gtrsim R_L$ for $f_{\text{s.w.}}(k, \theta)$].
- ii) The "cut-off for $b < R_L$ " reduces the classical elastic $S(J)$ [with $|S(J)|^2 = 1 - |\text{total absorption}|^2$, J real] to the diffractive hadronic $S_{\text{Diff}}(J)$:

$$S_{\text{Diff}}(J) = e^{i\delta(J)} S(J) \quad (12)$$

such that for real J

$$|S_{\text{Diff}}(J)|^2 = 1 - \underbrace{(|\text{abs}|^2 + (1 - e^{-2\text{Im } \delta(J)})|S(J)|^2)}_{\text{nondiffractive overlap function}}$$

where $\delta(J)$ has to be chosen such as to "absorb" only the saddle point contributions of f_{geom} for $(\text{Re } J + \frac{1}{2})/k < R_L$.

Altogether: the "cut-off" leaves $f(k, \theta)$ unchanged at $\theta \simeq 0$ and leads to the dominance of the surface waves for all angles away from $\theta = 0$. The following discussion is restricted to the surface wave contributions only, leaving aside the region around $\theta \simeq 0$, which in fact turns out ¹¹⁾ to be very small.

The "translation" of the instability of the excited intermediate hadronic state into conditions in the analogue picture can be pushed even further to distinguish between elastic and inelastic diffraction. Equation (6) tells us that the surface waves are, in fact, excited on "orbits" around the surface corresponding to impact parameters

$$b_n \simeq \frac{\text{Re } \alpha_n + 1/2}{k} \simeq R_L + x_n \Delta R \quad (13)$$

which increase for increasing n. The rays following the innermost orbits are still in too close contact with the scatterer to contribute to elastic scattering [remember also the term $O(1/(k\rho_{\parallel}))^{1/3} \sqrt{N^2 - 1}$ in Eq. (6), which indicates a residual "interaction" of rays on the innermost orbits with the scatterer]. They are naturally made responsible for diffraction dissociation, while contributions of all higher surface waves make up elastic scattering.

CONSEQUENCES

This picture has a number of immediate consequences which are in nice agreement with features in the data.

i) High-mass diffraction dissociation: we expect contributions from the first few surface waves with dominance of the "innermost" one ($n=1$). Its impact parameter profile will be peripheral, induced by the single direct channel Regge pole at $\alpha_1(k)$:

$$\left| \frac{1}{(j - \alpha_1(k))(j + \alpha_1(k) + 1)} \right| \approx \frac{1}{k^2} \left| \frac{1}{b^2 - \left(\frac{\alpha_1(k) + 1/2}{k} \right)^2} \right| \quad (14)$$

peaking at

$$b \approx \frac{\text{Re} \alpha_1(k) + 1/2}{k} \approx R_{\perp} \approx 1 \text{ fermi.}$$

From Eq. (6) we expect the flattest possible t distribution in this framework of surface waves

$$\frac{d\sigma}{dt} \propto \frac{1}{(t \neq 0) \sqrt{-t}} e^{-2 \frac{\text{Im} \alpha_1(k)}{k} \sqrt{-t}} \quad (\text{with } k \cdot \theta \approx \sqrt{-t}) \quad (15)$$

(remember $\text{Im} \alpha_1 < \text{Im} \alpha_n$ for all $n \geq 2$).

ii) Elastic scattering: we expect roughly speaking contributions from all surface waves with the exception of the first few ones (i.e., contributions from all waves with $n \geq n_0$, say). In particular, one obtains an enormous diffraction peak at small values of t due to the accumulation of all surface waves leading to [cf. Eq. (9)]

$$i \frac{H_1^{(1)}((\alpha_{n_0}(k) + 1/2) \theta)}{k R_{\perp} \theta} \quad (16)$$

For sufficiently large values of $|t|$ the lowest surface wave ($n = n_0$) starts to dominate and an $e^{-A p_{\perp}}$ like behaviour¹⁶⁾ ("Orear plot!") results due to

$$\frac{d\sigma}{dt} \propto \frac{1}{(t \neq 0) \sqrt{-t}} e^{-2 \frac{\text{Im} \alpha_{n_0}(k)}{k} \sqrt{-t}} \quad (17)$$

Although also elastic scattering is basically a peripheral phenomenon in our interpretation (i.e., due to surface waves), the elastic impact parameter profile is clearly "central", i.e., peaked at small b due to the accumulation of infinitely many surface waves at small t [leading to " $iJ_1(R_{\perp}\sqrt{-t}/R_{\perp}\sqrt{-t})$ " near $t=0$].

iii) Slope mass correlation: after having discussed the two extremes of elastic scattering and high-mass diffraction dissociation, it is clear what happens qualitatively for any diffraction dissociation process: varying the mass of the diffractively produced system from low to high values causes the diffraction peak at small values of t to disappear, while the slope of $d\sigma/dt$ becomes flatter until it reaches its minimum slope given in Eq. (15), since $\text{Im } \alpha_n < \text{Im } \alpha_{n+1}$ due to $x_n < x_{n+1}$. This phenomenon appears in the data and is known under name of slope-mass correlation.

iv) Two-body quantum number exchange reactions: they may be written (using isospin invariance or higher symmetries) as differences of elastic cross-sections which are almost equal at high energies. Making the near-at-hand assumption that the two elastic cross-sections differ only in their contributions to their lowest surface wave ($n=n_0$), we find the behaviour (17) uniformly for small and larger values of $-t$ ($t=0$ excluded) without a diffraction peak at small values of t . Correspondingly, the impact parameter profile of the imaginary part of the amplitude becomes peripheral, i.e., peaks near $b \approx R_1$.

Let us point out that the dominance of one single s channel Regge pole $\alpha(k)$ for quantum number exchange reactions with $\text{Re } \alpha(k)/k \simeq R_1 \simeq 1$ fermi had been proposed by us ¹⁷⁾ and other authors ¹⁸⁾ and had been successfully compared with a large number of reactions.

Let us finally mention: this paper will be backed up (i) by successful semiquantitative comparisons with data, basing on a cigar-like interaction bag, and (ii) by a more extensive mathematical treatment in Refs. 11) and 13).

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