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Diffraction Dissociation and Eikonalization in High Energy pp and $p\bar{p}$ Collisions

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Abstract

We show that eikonal corrections imposed on diffraction dissociation processes calculated in the triple Regge limit, produce a radical change in the energy dependence of the predicted cross section. The induced correction is shown to be in general agreement with the recent Tevatron experimental data.

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Over the past few years, phenomenological investigations of Pomeron exchange processes have been almost exclusively confined to the study of elastic scattering and total cross sections [1-5]. Recently published Tevatron data [6, 7] on single diffraction dissociation (SDD), enables us to evaluate the compatibility of the parametrizations used to describe elastic and diffractive scattering, and whether, it is necessary to include screening corrections, to obtain a successful description of these processes.

A fundamental problem that that must be tackled when one attempts to make a comprehensive analysis of the published high energy data on SDD [6-11], is the fact that there is no unique, agreed upon, experimental definition of SDD. Experimental groups have used different, and not always mutually consistent methods of extracting the desired data. In addition, it is difficult to compare the values that the different experimental groups give for σ_{sd} , as in their evaluation of $\frac{d\sigma_{sd}}{dM^2 dt}$, they have used diverse integration limits for t and M^2 . Furthermore, their treatment of the correlations observed between M^2 and t are entirely different.

With the above limitations in mind, we present in this communication, a general study of SDD, which is compatible with the analysis of elastic scattering, and at the same time reproduces all the important features of the experimental data measured in SDD at high energies.

Even though the Pomeron was introduced into high energy physics more than 30 years ago, its exact definition and detailed substructure remain an enigma. In contrast to standard Regge trajectories, the Pomeron has no particles on the time-like sector of its trajectory. Nevertheless, it is required both phenomenologically, to describe the forward hadron-hadron scattering data, and theoretically to ensure that Regge theory is self consistent. Indeed, in a Reggeon field calculus the Pomeron is described as a ladder of Reggeons yielding [12] $\alpha(0) = 1$. We will refer to this as the "soft Pomeron".

A number of different models have been proposed to account for the rising hadron-hadron cross sections:

- 1) Donnachie and Landshoff [1] have advocated an ad hoc approach in which the soft Pomeron amplitude keeps its traditional form with $\alpha(0) = 1 + \Delta \simeq 1.08$. This simple model reproduces the qualitative features of the experimental data remarkably well.
- 2) Alternatively, one may perceive the Pomeron as a two gluon exchange [13], or more generally as a gluon ladder. Lipatov [14] has shown that such a ladder, when calculated within the framework of perturbative QCD, receives its major contribution from high p_{\perp} gluon exchanges. These give rise to a series of poles in the complex j -plane above unity. The summation of these poles yield the "hard Pomeron" with $\Delta = \frac{12}{\pi} \alpha_s \ln 2$. Bjorken has suggested [15] that the generic Pomeron may actually manifest itself in both soft and hard modes, each contributing in a different kinematical domain. Models based on a hybrid

Pomeron are very successful in reproducing the data [4, 5].

3) In the QCD inspired model [2, 3], the growth of the total cross section is associated with the greater probability of semi-hard gluons to interact with increasing energy. In this case, the need to describe the data over a wide energy range also requires a hybrid model [3] consisting of a soft q-q background and semi-hard q-g and g-g interactions.

All the above models of the Pomeron have a intrinsic powerlike s^Δ rise of the total hadronic cross section. We note [4, 16] that the Pomeron amplitude proposed in [1] violates s-channel unitarity, just above the Tevatron energy range, for small b . In general, we expect the unitarity bound to induce screening effects which saturate the growth of σ_{tot} , making $\sigma_{tot} \leq \ln^2 s$, which is compatible with the Froissart bound. Technically, this is most easily achieved through eikonalization [17], in which the amplitude discussed above serves as the lowest order input to the eikonal expansion. Even though in the eikonal model one only sums over elastic rescattering, ignoring diffraction in the intermediate states, it has the advantage of being simple to apply. In addition, it introduces the natural scale of the screening corrections, and allows one to explore different models of the Pomeron.

The main purpose of this letter is to examine the role played by eikonalization in SDD. This is investigated utilizing a simple Regge-like Pomeron [1]. Extending the same formalism to include an input Lipatov type Pomeron is straightforward. As the presently available diffractive data is not sufficiently refined to enable one to discriminate between these models of the Pomeron, we shall not discuss it in detail here.

The simplest way to write down the eikonal formulae is to consider the scattering process in impact parameter space. Our amplitude is normalised so that

$$\frac{d\sigma}{dt} = \pi |f(s, t)|^2 \quad (1)$$

$$\sigma_{tot} = 4\pi \text{Im} f(s, 0) \quad (2)$$

The scattering amplitude in b -space is defined as

$$a(s, b) = \frac{1}{2\pi} \int d\mathbf{q} e^{-i\mathbf{q}\cdot\mathbf{b}} f(s, t) \quad (3)$$

where $t = -q^2$.

In this representation

$$\sigma_{tot} = 2 \int db \text{Im} a(s, b) \quad (4)$$

$$\sigma_{el} = \int db |a(s, b)|^2 \quad (5)$$

s-channel unitarity when written in the diagonalised form implies

$$2\text{Im}a(s, b) = |a(s, b)|^2 + G_{in}(s, b) \quad (6)$$

where

$$\sigma_{in} = \int db G_{in}(s, b) \quad (7)$$

We list below several assumptions that we make regarding the eikonal model:

1) At high energy $a(s, b)$ is assumed to be pure imaginary and can be reduced to the simple form

$$a(s, b) = i(1 - e^{-\Omega(s, b)}) \quad (8)$$

where $\Omega(s, b)$ is a real function. Analyticity and crossing symmetry are easily restored to our oversimplified parametrization by substituting $s^\alpha \rightarrow s^\alpha e^{-i\pi\alpha/2}$.

2) From eq. (6) we can express $G_{in}(s, b)$ as

$$G_{in}(s, b) = 1 - e^{-2\Omega(s, b)} \quad (9)$$

where $e^{-2\Omega(s, b)}$ denotes the probability that no inelastic interaction takes place at impact parameter b .

3) We write the t-channel Pomeron exchange as

$$\Omega(s, b) = \nu(s) e^{-\frac{b^2}{R^2(s)}} \quad (10)$$

In the simple Regge pole model with a trajectory $\alpha_P(t) = 1 + \Delta + \alpha't$. We have

$$\nu(s) = \frac{\sigma_0}{2\pi R^2(s)} \left(\frac{s}{s_0}\right)^\Delta \quad (11)$$

where

$$R^2(s) = 4[R_0^2 + \alpha' \ln \frac{s}{s_0}] \quad (12)$$

and $\sigma_0 = \sigma(s_0)$. Agreement with the pp ($\bar{p}p$) data is obtained with $R_0^2 = 5.2 \text{ GeV}^{-2}$ and $\alpha' = 0.25 \text{ GeV}^{-2}$.

Eqs.(10-12) lead to simple expressions for the total and inelastic cross sections with $\sigma_{el} = \sigma_{tot} - \sigma_{inel}$ (see Fig. 2a).

$$\sigma_{tot} = 2\pi R^2(s) [\ln \nu(s) + C - Ei(-\nu(s))] \xrightarrow{\nu \gg 1} 2\pi R^2(s) [\ln \nu(s) + C] \quad (13)$$

$$\sigma_{in} = \pi R^2(s) [\ln 2\nu(s) + C - Ei(-2\nu(s))] \xrightarrow{\nu \gg 1} \pi R^2(s) [\ln 2\nu(s) + C] \quad (14)$$

where $Ei(x) = \int_{-\infty}^x \frac{e^t}{t} dt$, and $C = 0.5773$ is the Euler constant.

The standard approach to evaluate single diffractive dissociation is through the 3-body optical theorem [18] leading to the PPP and PPR diagrams of interest (see Fig. 1). The appropriate cross section is

$$M^2 \frac{d\sigma_{sd}}{dM^2 dt} = \left(\frac{s}{M^2}\right)^{2\Delta+2\alpha't} [G_{PPP}(t) \left(\frac{M^2}{s_0}\right)^\Delta + G_{PPR}(t) \left(\frac{M^2}{s_0}\right)^{-\frac{1}{2}}] \quad (15)$$

where all of the relevant couplings have been absorbed into $G_{PPP}(t)$ or $G_{PPR}(t)$. M^2 denotes the mass of the diffractive system, and for the Regge trajectory we have taken $\alpha_R(t) = \frac{1}{2} + t$.

Eq. (15) can be rewritten in the impact parameter representation

$$\begin{aligned} \frac{M^2 d\sigma_{sd}}{dM^2} &= G_{PPP} \sigma_0^2 \left(\frac{s}{M^2}\right)^{2\Delta} \left(\frac{M^2}{s_0}\right)^\Delta \frac{1}{[\pi \bar{R}_1^2(\frac{s}{M^2})]^2 \pi \bar{R}_1^2(\frac{M^2}{s_0})} \int db db' e^{-\frac{2(b-b')^2}{\bar{R}_1^2(\frac{s}{M^2})} - \frac{b'^2}{\bar{R}_1^2(\frac{M^2}{s_0})}} \\ &+ G_{PPR} \sigma_0^2 \left(\frac{s}{M^2}\right)^{2\Delta} \left(\frac{M^2}{s_0}\right)^{-\frac{1}{2}} \frac{1}{[\pi \bar{R}_1^2(\frac{s}{M^2})]^2 \pi \bar{R}_2^2(\frac{M^2}{s_0})} \int db db' e^{-\frac{2(b-b')^2}{\bar{R}_1^2(\frac{s}{M^2})} - \frac{b'^2}{\bar{R}_2^2(\frac{M^2}{s_0})}} \end{aligned} \quad (16)$$

where

$$\bar{R}_i^2\left(\frac{s}{M^2}\right) = 2R_{0i}^2 + r_{0i}^2 + 4\alpha' \ln\left(\frac{s}{M^2}\right) \quad (17)$$

$r_{0i} \leq 1\text{GeV}^{-2}$ denotes the radius of the triple vertex [19]. $\bar{R}_1^2(\frac{s}{M^2}) = 2B_{sd}$, where B_{sd} denotes the slope of the SDD cross section. Upon integrating eq. (16) we have

$$M^2 \frac{d\sigma_{sd}}{dM^2} = \frac{\sigma_0^2}{2\pi \bar{R}_1^2(\frac{s}{M^2})} \left(\frac{s}{M^2}\right)^{2\Delta} [G_{PPP} \left(\frac{M^2}{s_0}\right)^\Delta + G_{PPR} \left(\frac{M^2}{s_0}\right)^{-\frac{1}{2}}] \quad (18)$$

We will now comment on consequences of the above result and its relevance when compared to experimental data [6-11]:

1) We expect the forward SDD differential nuclear slope to be in the range $\frac{1}{2}B_{el} < B_{sd} < B_{el}$, where $B_{el} = 2R^2(s)$ denotes the appropriate elastic scattering slope. In general, B_{sd} is M^2 dependent. An explicit logarithmic dependence is implied by the definition of $\bar{R}_i(\frac{s}{M^2})$ in eq.(17). We also note that due to the different M^2 power dependences, the PPR contribution is concentrated at lower values of M^2 than the PPP. For energies in the ISR-Tevatron range, where $\ln^2 s \geq R_0^2$, we expect qualitatively, that $B_{sd} \geq \frac{1}{2}B_{el}$ with a very moderate $\ln(\frac{s}{M^2})$ dependence. This is in agreement with the data. We are unable to make a numerical fit due to strong correlations between M^2 and t , observed at small

values of M^2 . We strongly urge that measurements of B_{sd} be made for higher values of the mass spectrum, say $M^2 \geq 16 \text{ GeV}^2$.

2) The M^2 dependence of the SDD cross section is dominated by $[G_{PPP}(M^2)^{-(1+\Delta)} + G_{PPR}(M^2)^{-(1.5+2\Delta)}]$. If we express this dependence by $(M^2)^{-\alpha_{eff}}$, we expect $(\alpha_{eff} - 1) > \Delta$ and that $(\alpha_{eff} - 1)$ approaches Δ from above in the limit of very high s when the importance of the PPR term diminishes. This behaviour is corroborated by the two recent studies [6, 7] of the M^2 distribution at the Tevatron. In passing we note, that the experiments at the FNAL [8] and ISR [9] reported approximate scaling, i.e. a $(M^2)^{-1}$ behaviour. This is most probably due to the much narrower M^2 interval investigated. The approximation in which we only consider the PPP + PPR terms is obviously not sufficient to describe data at lower energies, where lower lying trajectories are important [20].

3) Eq. (18) predicts a strong powerlike $s^{2\Delta}$ dependence of the differential as well as the integrated SDD cross section. This is a much stronger energy dependence than the predicted s^Δ behaviour of σ_{tot} , and clearly not compatible with either theory or data. Indeed, the CDF data [7] taken at $\sqrt{s} = 546$ and 1800 GeV show only a moderate 20% increase of the appropriate cross sections. This should be compared with an 80% increase expected from a $s^{2\Delta}$ behaviour with $\Delta = 0.125$, as reported by CDF [7].

Obviously, eq.(14) violates unitarity. Unitarity is restored, in the eikonal model, by multiplying the integrand of eq. (16) by $e^{-2\eta(s,b)}$ (see Fig. 2b). The resulting cross section is

$$\begin{aligned} \frac{M^2 d\sigma_{sd}}{dM^2} = & G_{PPP}\sigma_0^2 \left(\frac{s}{M^2}\right)^{2\Delta} \left(\frac{M^2}{s_0}\right)^\Delta \frac{1}{[\pi \bar{R}_1^2(\frac{s}{M^2})]^2 \pi \bar{R}_1^2(\frac{M^2}{s_0})} \\ & \int db db' e^{-\nu(s)e^{-\frac{b^2}{R^2(s)}}} \cdot e^{-\frac{2(b-b')^2}{R_1^2(\frac{s}{M^2})} - \frac{b'^2}{R_2^2(\frac{M^2}{s_0})}} \\ & + G_{PPR}\sigma_0^2 \left(\frac{s}{M^2}\right)^{2\Delta} \left(\frac{M^2}{s_0}\right)^{-\frac{1}{2}} \frac{1}{[\pi \bar{R}_1^2(\frac{s}{M^2})]^2 \pi \bar{R}_2^2(\frac{M^2}{s_0})} \\ & \int db db' e^{-\nu(s)e^{-\frac{b^2}{R^2(s)}}} \cdot e^{-\frac{2(b-b')^2}{R_1^2(\frac{s}{M^2})} - \frac{b'^2}{R_2^2(\frac{M^2}{s_0})}} \end{aligned} \quad (19)$$

where $\nu(s)$ is given by eq. (11) and $R^2(s)$ by eq. (12). After integration we have

$$\begin{aligned} \frac{M^2 d\sigma_{sd}}{dM^2} = & \frac{\sigma_0^2}{2\pi \bar{R}_1^2(\frac{s}{M^2})} \left(\frac{s}{M^2}\right)^{2\Delta} \cdot [G_{PPP}(\frac{M^2}{s_0})^\Delta a_1 \frac{1}{(2\nu(s))^{a_1}} \gamma(a_1, 2\nu(s)) \\ & + G_{PPR}(\frac{M^2}{s_0})^{-\frac{1}{2}} a_2 \frac{1}{(2\nu(s))^{a_2}} \gamma(a_2, 2\nu(s))] \end{aligned} \quad (20)$$

where

$$a_i = \frac{2R^2(s)}{\bar{R}_1^2(\frac{s}{M^2}) + 2\bar{R}_i^2(\frac{M^2}{s_0})} \quad (21)$$

and $\gamma(a, 2\nu)$ denotes the incomplete Euler gamma function $\gamma(a, 2\nu) = \int_0^{2\nu} z^{a-1} e^{-z} dz$.

We list below the important consequences of the expression we obtained in eq. (20).

- 1) The b -space SDD amplitude, which is the integrand of eq. (19), differs from the intrinsic integrand of eq. (16) by the corrective multiplicative factor $e^{-2\Omega(a,b)}$. Whereas, the unabsorbed b -space SDD amplitude is central and can be approximated by a Gaussian centered at $b = 0$, the corrected amplitude has a dip at $b = 0$, and its Gaussian approximation is centered at some $b = b_0 \neq 0$. This behaviour suggests that the generalized unitarity condition [20] is satisfied. This is consistent with the general pattern expected of SDD b -space amplitudes after screening has been included [21].
- 2) Our qualitative observation that $B_{sd} \geq \frac{1}{2} B_{el}$ is unchanged. We expect the ratio $\frac{B_{sd}}{B_{el}}$ to grow with energy, up to a limiting value of 1.
- 3) The dominant M^2 dependence of $\frac{d\sigma_{sd}}{dM^2}$ is identical to that determined from eq.(18). We stress, that the two properties of the triple Regge model, those concerning the t and M^2 dependence, which are in agreement with experiment, are essentially unchanged once the eikonal correction is made to the original SDD amplitude.
- 4) Eq. (20) exhibits a weak s -dependence. This is best seen if we examine our result in the high energy limit, where we have $a_i \rightarrow 2$, and $\gamma[a_i, 2\nu(s)] \rightarrow \Gamma(2)$. Thus the factor $s^{2\Delta}$ is compensated by $[\frac{1}{\nu(s)}]^{a_i}$ and eq. (20) reduces to

$$M^2 \frac{d\sigma_{sd}}{dM^2} = \pi \Gamma(2) \sigma_0^2 R^2(s) [G_{PPP}(\frac{M^2}{s_0})^{-\Delta} + G_{PPR}(\frac{M^2}{s_0})^{-(\frac{1}{2}+2\Delta)}] \quad (22)$$

Since σ_{sd} is not very sensitive to the high M^2 integration limit, we find that σ_{sd} depends on s only through $R^2(s)$. Our result indicated that the changes induced by eikonalization on σ_{tot} and σ_{sd} are quite different. For σ_{tot} the input s^Δ power behaviour is modified to $\ln^2 s$, the energy scale at which this change becomes appreciable is at $\sqrt{s} \approx 3$ TeV [4]. For σ_{sd} the input $s^{2\Delta}$ power behaviour is modified to $\ln s$, this occurs at an energy scale which is considerably lower i.e. $\sqrt{s} \approx 300$ GeV. In addition we expect that $\frac{\sigma_{sd}}{\sigma_{tot}} \xrightarrow{s \rightarrow \infty} 0$. To test this theoretical prediction, we need to know Δ and the ratio between G_{PPP} and G_{PPR} .

These two parameters are obviously correlated. Donnachie and Landshoff [1] suggest a global σ_{tot} fit with $\Delta = 0.08$. This choice is compatible with the CDF M^2 distribution, if the PPR contributes 40% of the integrated σ_{sd} at 546 GeV. The above value quoted for Δ , which was suggested in [1], used the E710 measurement [23] of σ_{tot} at $\sqrt{s} = 1800$ GeV. A recent CDF measurement [24] at the same energy has a considerably higher value for σ_{tot} which is consistent with a value of $\Delta = 0.11$. The corresponding PPR contribution to σ_{sd} at 546 GeV is now reduced to 15%. Irrespective of which value of Δ we use, we are not able to find an adequate overall fit to SDD data measured over the entire energy range [6-11]. We feel that this is due to the following experimental and theoretical difficulties:

1) As we noted previously, comparing σ_{sd} values obtained by different experiments is not very instructive, due to the diverse constraints and algorithms used by the different groups.

2) To minimize experimental uncertainties we consider the two CDF measurements at $\sqrt{s} = 546$ and 1800 GeV, where they find [7] $R = \frac{\sigma_{sd}(1800)}{\sigma_{sd}(546)} = 1.20 \pm 0.06$, with $1.4 \leq M^2 \leq 0.15s$. If we take $\Delta = 0.08$ we predict a ratio of $R_{PPP} = 1.35$ for the PPP term, and for the PPR term $R_{PPR} = 1.25$. Assuming the PPR contribution to account for 40% of the SDD cross section at 546 GeV, we have a theoretical prediction of $R = 1.31$. For $\Delta = 0.11$, we obtain $R_{PPP} = 1.35$ and $R_{PPR} = 1.20$. This gives us a prediction for $R = 1.33$, assuming the PPR to account for 15% of the SDD cross section at 546 GeV.

3) The CDF group start their M^2 integration at $M_{min}^2 = 1.4 \text{ GeV}^2$, which is much too low for any triple Regge analysis. To eliminate the region of low diffractive masses, we compare with the experimental ratio quoted by CDF [25] of $R = 1.24 \pm 0.10$ obtained with $M_{min}^2 = 16 \text{ GeV}^2$. For $\Delta = 0.08$ we obtain $R = 1.34$, while for $\Delta = 0.11$ we have $R = 1.37$.

4) Extrapolation of our model to ISR energies (using values of the parameters normalised to the CDF data) underestimates the measured values of σ_{sd} . This is not unexpected, as our simple model with only PPP + PPR contributions is clearly not sufficient at ISR energies, where a more detailed analysis [20] demonstrates the importance of lower lying trajectories at these energies. Examining SDD data over the whole energy range [6-11], it appears that screening corrections become important at energies lower than that predicted by our eikonal model. This is not surprising, as in our treatment of eikonalization we have only included elastic rescattering effects in the intermediate states, while completely ignoring diffractive effects or so called inelastic shadowing correction (see Fig. 2c)[26]. Such corrections cannot be considered to be small as the ratio σ_{sd}/σ_{el} is of the order of $\frac{1}{2}$ at the Tevatron energies. It means that dimensionless triple Pomeron vertex introduced in eq.(19) is about $\frac{1}{8}$ and diagrams of Fig. 2c should be taken into account at the next stage of our approach.

5) In contrast to point 4) we expect the extrapolation of our results to extremely high energies to be trustworthy. Integrating over $1.4 \text{ GeV}^2 \leq M^2 \leq 0.15s$, we predict that σ_{sd}

= 13.3 and 13.9 mb at $\sqrt{s} = 16$ and 40 TeV respectively, demonstrating the very weak s dependence predicted by our model.

In conclusion, we wish to emphasize that our model does reproduce the main features of SDD above 300 GeV, in particular the exceedingly moderate dependence of σ_{sd} on s . The model which does not include lower lying Regge trajectories is too simple to successfully describe the SDD data at lower energies.

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Figure Captions

Fig.1: SDD in the triple Regge approximation.

Fig.2:

- a) Screening corrections in the eikonal approximation to elastic scattering.
- b) Screening corrections in the eikonal approximation to SDD.
- c) Inelastic shadowing (screening) corrections to SDD.

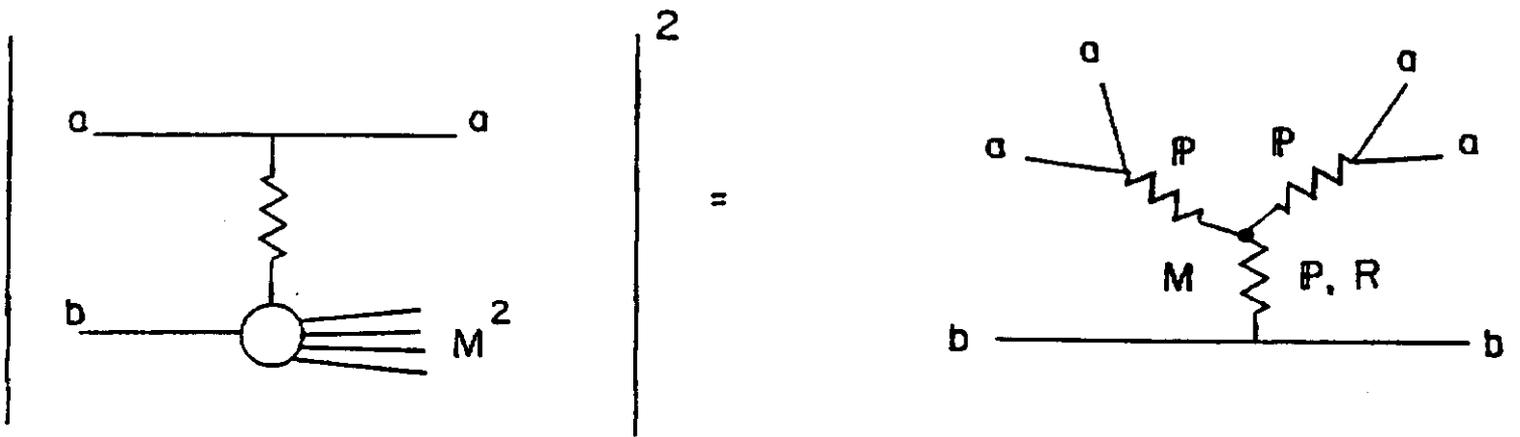
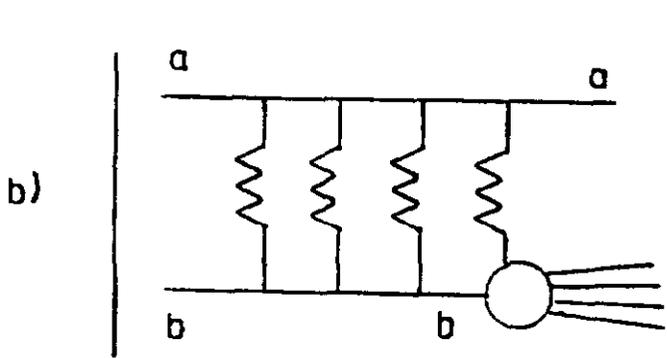
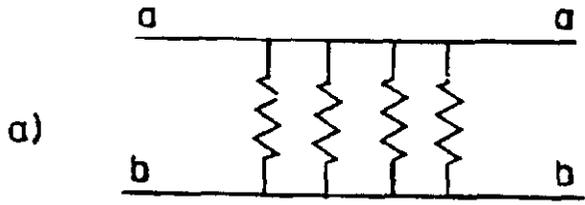


Fig. 1.



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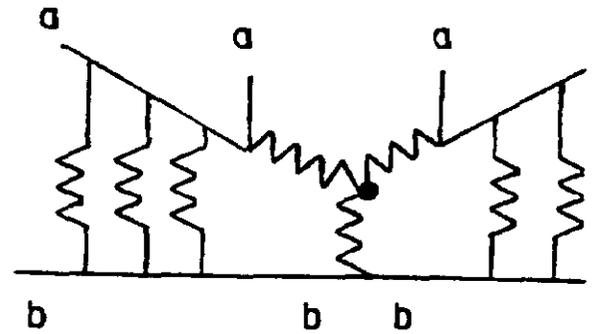


Fig. 2.

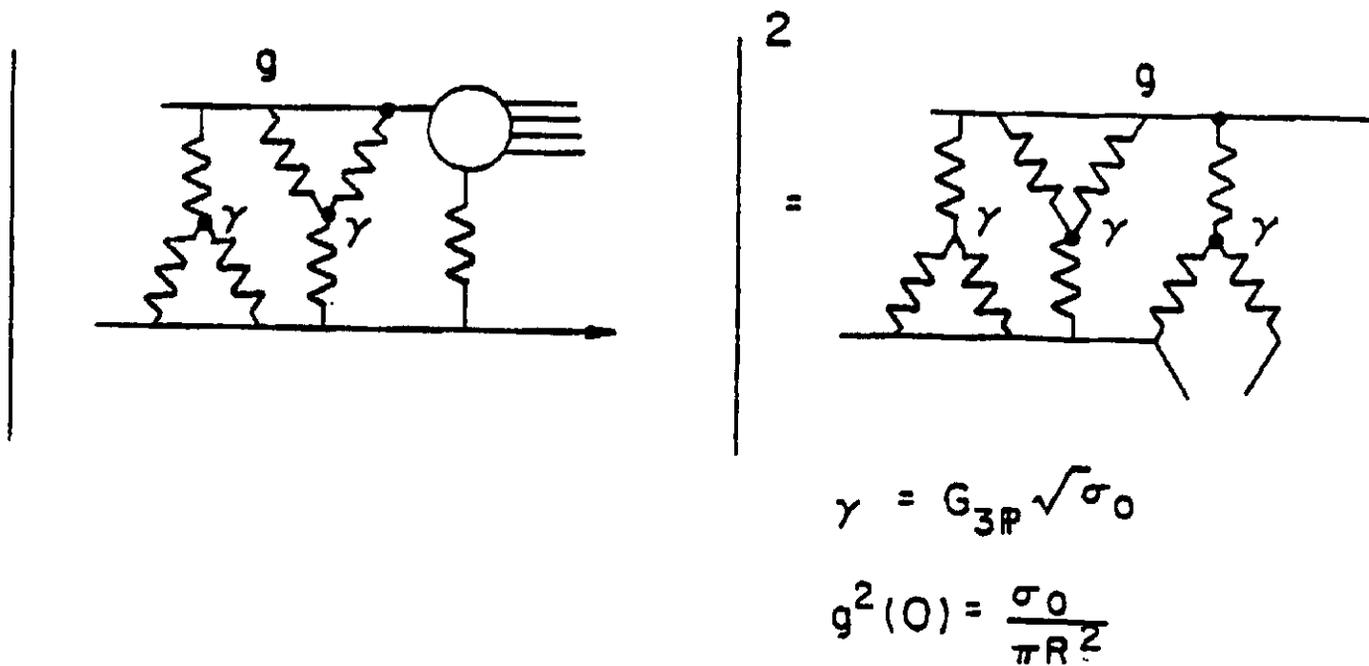


Fig. 2.