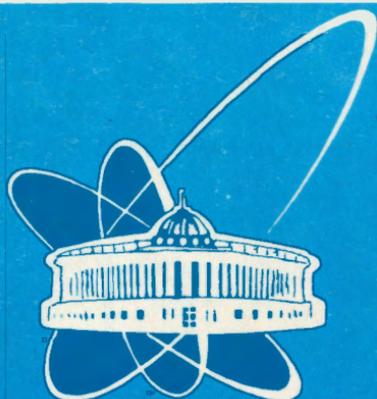


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DIELECTRON PRODUCTION IN PROTON-PROTON  
AND PROTON-DEUTERON COLLISIONS  
AT 1-5 GeV

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## I. INTRODUCTION

The available experimental data<sup>1-3</sup> on dielectron production in proton-proton or proton-nucleus or nucleus-nucleus collisions at 1-5 GeV·A bombarding energies have stimulated a series of theoretical investigations<sup>4-13</sup> of the elementary production process. The reasons for this interest are obvious. Dielectrons are thought to represent one of the promising signals which can directly probe the dense and hot nuclear matter produced in heavy-ion collisions at intermediate energies. The reliable description of various elementary reaction channels for dielectron production serve as an input to theoretical models and event generators for simulating heavy-ion collisions. These simulations are needed to unfold dielectron spectra and to get the wanted information about the compressed and heated nuclear matter. Also via the dielectron decay channel of vector mesons one can probe the behavior of such mesons in an excited nuclear environment. The second-generation precision spectrometers are devoted to these investigations.

The mentioned theoretical investigations of the elementary production mechanisms of dielectrons in  $pN$  reactions have step-wise improved the understanding of the relevant basic processes<sup>4-14</sup>. These investigations have their own right, also with respect to new hadron facilities (e.g., COSY in Jülich), which are devoted to deeper insight into hadron structure, hadronic reactions, and photon-hadron interactions. Concerning the dielectron production, the models, with appropriate parametrization, are in satisfactory agreement with available experimental data<sup>3</sup> which are still with low statistics. New data with high statistics are expected in the near future. Then, theoretical estimates and underlying assumptions can be tested better since they depend on certain model parameters which are difficult to fix without experimental data. For example, in dielectron processes the off-shell behavior of the strong interaction part is probed in a wider kinematic regime than in the case of real photon bremsstrahlung or elastic scattering of hadrons or light nuclei. The time-like (half off-shell) form factor of the nucleon is still unknown in the region where the transfer momentum is near the vector meson masses. Details of certain channels, e.g.,  $pn \rightarrow \eta X$ , are rather unsettled and reliable data do not yet exist. The latter fact is partially related to the difficulty in getting reliable information on  $pn$  reactions, in general, via light nuclei by subtracting masking many-body effects.

The aim of the present paper is to re-estimate dielectron production cross sections in elementary nucleon-nucleon subprocesses and to apply them in  $pp$  and  $pd$  reactions at 1, 2 and 5 GeV. We rely here on the vector dominance model (VDM) that has proven to be a useful guiding principle for hadron-photon interactions<sup>15</sup>. The VDM

form factor has been implemented, e.g., in Ref.<sup>14</sup> and it has been found to give reasonable results only if nuclear matter corrections are taken properly into account. However, in the  $pp$ , and  $pd$ , and  $p$ -light nuclei reactions such a nuclear correction is not operative, and one has to implement the VDM form factor in an alternative way<sup>16</sup>. Here, we present a detailed study of the underlying microscopy in elementary subprocesses<sup>12</sup> of dielectron production in nucleon–nucleon collisions.

Remember that at 1 to 5 GeV, the main dilepton sources are the following:  $\pi^+\pi^-$  annihilation,  $\pi^0, \Delta, \eta, \omega$  Dalitz decays,  $pN$  bremsstrahlung, and direct vector meson decays. For  $pp$  and  $pd$  reactions, when focusing on dileptons with invariant masses in the range of 0.15 – 1.2 GeV, only  $\Delta, \eta, \omega$  Dalitz decays,  $pN$  bremsstrahlung and direct vector meson decays are essential.

The  $\Delta$  Dalitz decay is one of the strongest dilepton channel. In Refs.<sup>5,6,13,14</sup>, it is used within a model, where the  $\Delta$  production cross section in  $pN$ -collisions is taken as a constant at fixed kinetic energy and is independent of the momentum transfer to the target nucleon, which is related to the  $\Delta$  mass directly. Experimental data<sup>17,18</sup> and theoretical models<sup>19</sup>, however, show such a dependence: the  $\Delta$  production cross section decreases with increasing values of the momentum transfer. The maximum of the dilepton invariant mass depends directly on the  $\Delta$  mass and therefore, one can expect some dynamical suppression of dileptons at large invariant masses. We find that, in spite of this suppression, it is almost compensated by the VDM form factor enhancement; the form of the spectra changes, and they obtain a resonance-like behavior at invariant masses near the rho-meson mass. In principle, such a resonance-like behavior at the initial energy of 5 GeV has been found in Refs.<sup>5,13</sup> but in those calculations a too large value for the  $\Delta$  production cross section was used; therefore, a re-estimate of the delta production and delta decay mechanisms is needed.

The analysis of the  $pN$  bremsstrahlung contribution in most previous papers is practically based on the so-called soft photon approximation. The soft photon approximation includes several approximations. A few of them are acceptable (e.g., keeping only the electric part of the hadron current and neglecting radiating from the virtual propagators and vertices), while others (e.g. integration over unobservable phase space kinematical region) result in an overestimation of the cross section. This overestimation is sometimes corrected by a phase space volume reduction factor<sup>5</sup>. Here, we improve this approach by correct phase space integration.

A further problem concerns the  $pp$  bremsstrahlung contribution. It is usually assumed that, because of destructive interference of direct and exchanged amplitudes

of the electric part of the bremsstrahlung matrix element, the amplitudes compensate each other. But that is not correct exactly, especially for high energy. The calculation of  $pp$  bremsstrahlung at 4.9 GeV in the soft photon approximation of Ref.<sup>13</sup> shows that it may overestimate the  $pn$  bremsstrahlung. So one expects only partial compensation of the electric part and both the  $pn$  and  $pp$  contribution should be taken into account. This conclusion is confirmed also in the recent paper<sup>16</sup> where both the proton nucleon bremsstrahlung and the delta Dalitz decay are considered simultaneously within effective one boson exchange model.

Another strong source of dileptons is the  $\eta$  Dalitz decay. Most interesting here is the dilepton production near 1 GeV. The threshold energy for  $\eta$  production is about 1.26 GeV, and in  $pp$  collision this channel is suppressed kinematically. For  $pd$  collisions, however, it is open, but one has to deal carefully with the mechanisms of subthreshold  $\eta$  production.

These are the main items we are going to analyze in some detail. Our paper is organized as follows. In Sec. 2, we analyze the Dalitz  $\Delta$  decay rate. In Sec. 3, we discuss  $pn$  and  $pp$  bremsstrahlung contributions in  $pd$  reactions within the soft photon approximation for the electromagnetic hadron current but with exact multiple integration of the resulting matrix element squared. In Sec. 4 and 5, we discuss, respectively, the contribution of the  $\eta, \omega$  Dalitz decays and  $\rho, \omega$  direct decays. In Sec. 6, we present the calculated cross sections of various reaction channels and compare the dilepton production in  $pp$  and  $pd$  interactions. A summary is given in Sec. 7.

## II. DALITZ DELTA DECAY

The cross section of the delta Dalitz decay is represented as follows:

$$\frac{d\sigma_{\Delta \rightarrow e^+e^-N}}{dM^2} = \int_{(m_N+m_\pi)^2}^{(\sqrt{s}-m_N)^2} dM_\Delta^2 \tilde{\sigma}_\Delta(s, M_\Delta) D(M_\Delta) \frac{1}{\Gamma_\Delta} \frac{d\Gamma_{\Delta \rightarrow e^+e^-N}}{dM^2}, \quad (2.1)$$

where

$$\tilde{\sigma}_\Delta(s, M_\Delta) = \frac{1}{16\pi s(s-4m_N^2)} \int_{t_{\min}(M_\Delta)}^{t_{\max}(M_\Delta)} dt |T_\Delta(s, t, M_\Delta)|^2 \quad (2.2)$$

and  $M_\Delta$  and  $\Gamma_\Delta$  are the mass and total width of an intermediate delta;  $T_\Delta$  stands for the delta production matrix element,  $t$  denotes, as usual, the momentum transfer

squared at the  $NN \rightarrow \Delta N'$  vertex, and  $(d\Gamma/dM^2)^{\Delta \rightarrow e^+e^-N}$  describes the differential width of the delta decay into a dilepton with invariant mass  $M$ . The “weight function”  $D(M_\Delta)$  is proportional to the  $\Delta$ -propagator squared which leads to the relativistic Breit-Wigner form

$$D(M_\Delta) = \frac{1}{\pi} \frac{M_\Delta \Gamma_\Delta}{(M_\Delta^2 - M_{\Delta 0}^2)^2 + M_\Delta^2 \Gamma_\Delta^2} \quad (2.3)$$

with the mean value of  $\langle M_\Delta \rangle = M_{\Delta 0} = 1.232 \text{ GeV}/c^2$ .

The simplest form of the  $\Delta N\pi$  vertex is described by the interaction Lagrangian

$$\mathcal{L}_{\Delta N\pi} \sim \bar{\psi}_N(p_n) \psi_\Delta^\mu(P_\Delta) k_\mu \varphi(k), \quad (2.4)$$

where  $\psi_N(p_n)$ ,  $\psi_\Delta^\mu(P_\Delta)$ ,  $\varphi(k)$  are the nucleon, delta and pion wave functions, respectively, and  $k_\mu$  denotes the pion four momentum. Direct evaluation of the decay matrix element leads to a mass dependence of the  $\Delta$  width in Eq. (2.3),

$$\Gamma_\Delta(M_\Delta) = C \frac{(M_\Delta + m_N)^2 - m_\pi^2}{M_\Delta^2} |\mathbf{k}|^3. \quad (2.5)$$

Here  $\mathbf{k}$  denotes the *c.m.* momentum in the  $\pi N$ -channel, and the constant  $C$  is determined by the condition  $\Gamma_\Delta(M_{\Delta 0}) = 110 \text{ MeV}$ . The dependence  $\Gamma_\Delta(M_{\Delta 0})$  in Eqs. (2.3), (2.5) differs from the corresponding ones used in Ref.<sup>6</sup>; however, both of them coincide at a few per cent level.

The  $\Delta$ -decay probability  $(d\Gamma/dM^2)^{\Delta \rightarrow e^+e^-N}$  is calculated on the basis of the  $\Delta N\gamma$ -interaction Lagrangian

$$\mathcal{L}_{\Delta N\gamma} = e F_\gamma(M^2) \bar{\psi}_\Delta^\beta \Gamma_{\beta\mu} \psi_N A^\mu, \quad (2.6)$$

where  $A^\mu$  denotes the electromagnetic four-potential, and  $F_\gamma(m^2)$  is the vector dominance time-like electromagnetic form factor. For baryons this form factor in the kinematical region is still unknown. Following the vector dominance principle we use the minimal way to incorporate it: we assume that this form factor has a unique form for all hadrons, that is we use the  $\pi\pi\gamma$  VDM form factor. The physical meaning of this is quite clear: the virtual photon interacts with the pion cloud surrounding the nucleons and deltas. We employ the experimentally established parametrization<sup>4</sup>

$$F_\gamma^2(M^2) = \frac{m_\rho^4}{(M^2 - m_\rho^2)^2 + (m_\rho \Gamma'_\rho)^2}, \quad (2.7)$$

with  $m'_\rho = 761 \text{ MeV}$ ,  $\Gamma'_\rho = 118 \text{ MeV}$ ; the vertex function  $\Gamma_{\beta\mu}$  (2.6) is taken from Ref.<sup>20</sup>. The result of a direct calculation may be written as

$$\frac{1}{\Gamma_\Delta} \frac{d\Gamma}{dM^2} \Delta \rightarrow e^+e^-N = \frac{\alpha}{3\pi M^2} B_\gamma^\Delta(M_\Delta) R_\gamma^\Delta(M_\Delta, M), \quad (2.8)$$

$$R_\gamma^\Delta(M_\Delta, M) = \frac{\Gamma^{\Delta \rightarrow \gamma^* N}(M_\Delta, M)}{\Gamma^{\Delta \rightarrow \gamma N}(M_\Delta, 0)}.$$

Here  $B_\gamma^\Delta(M_\Delta)$  stands for the branching ratio of the electromagnetic width to the total delta decay width [ $B_\gamma^\Delta(M_\Delta = 1.232 \text{ GeV}) \simeq 0.6 \cdot 10^{-2}$ ] and  $R_\gamma^\Delta$  is the ratio of the electromagnetic delta decay widths for virtual to real photons,

$$R_\gamma^\Delta(M_\Delta, M) = \frac{M_\Delta M^2 + 5M_\Delta q_0^2 - 3M^2 m_N - 3M^2 q_0 - 3m_N q_0^2 - 3q_0^3}{q_0^2(5M_\Delta - 3m_N - 3q_0)} \times (\lambda(M_\Delta^2, m_N^2, M^2)/\lambda(M_\Delta^2, m_N^2, 0))^{1/2}, \quad (2.9)$$

and  $q_0 = (M_\Delta^2 + M^2 - m_N^2)/(2M_\Delta)$ ,  $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + yz + xz)$ .

For the calculation of the  $t$ -integrated cross section  $\tilde{\sigma}_\Delta(s, M_\Delta)$  in Eq. (2.2), we adopt the one-pion exchange model. By straightforward calculation we find the  $\Delta$  production matrix element in the form

$$\hat{T}_\Delta = a \frac{(\bar{\psi}_\mu^\Delta \vec{\tau}_1 \psi_{N_1} k^\mu)(\bar{\psi}_N \gamma_5 \gamma^\nu \vec{\tau}_2 \psi_{N_2} k_\nu)}{k^2 - m_\pi^2} \frac{\Lambda_{NN\pi}^2 - m_\pi^2}{\Lambda_{NN\pi}^2 - k^2} \frac{\Lambda_{\Delta N\pi}^2 - m_\pi^2}{\Lambda_{\Delta N\pi}^2 - k^2}, \quad (2.10)$$

where  $m_\pi = 140 \text{ MeV}/c$ , and  $k_\mu$  is the four momentum of the exchanged pion. The constant  $a$  is determined by the normalization condition

$$\int_{(m_N + m_\pi)^2}^{(\sqrt{s} - m_N)^2} dM_\Delta^2 \tilde{\sigma}_\Delta(s, M_\Delta) D(M_\Delta) = \sigma_\Delta(s). \quad (2.11)$$

In the above formula  $\sigma_\Delta(s)$  stands for the  $\Delta$  production cross section which we take as a product of the well-known analytical parametrization of the  $\Delta$  production cross section of Ver West and Arndt<sup>21</sup>,  $\sigma_\Delta^{V-A}$ , and the “high energy” correction factor

$$\sigma_\Delta(s) = \sigma_\Delta^{V-A} \frac{\theta(E - E_0)}{1 + A(E - E_0)}, \quad (2.12)$$

where  $E_0 = 1.3 \text{ GeV}$ ,  $A = 0.5 \text{ GeV}^{-1}$ , and  $E$  is the projectile kinetic energy in the laboratory system. The correction factor in Eq. (2.12) is introduced to ensure the reproduction of experimental data above  $E_0$ .

The  $t$ -integrated  $\Delta$  production cross section  $\tilde{\sigma}_\Delta$  in Eq. (2.1) at fixed  $\sigma_\Delta(s)$  depends on the cut-off parameters  $\Lambda_{NN\pi}$ ,  $\Lambda_{\Delta N\pi}$  in Eq. (2.10), which should be fitted to the differential cross section of the delta production. Fig. 1 shows the result of the

fitting procedure for the  $pp \rightarrow n\Delta^{++}$  reaction at initial kinetic energies 1.084 and 2.948 GeV. The solid lines correspond to the exclusive delta production cross section with the  $T$  matrix of Eq. (2.10) with  $\Lambda_{NN\pi} = \Lambda_{N\Delta\pi} = 0.7$  GeV. We calculate the spin averaged matrix element squared using relativistic Rarita–Schwinger propagator for spin 3/2 particles. The long dashed lines correspond to the calculation with the  $T$  matrix taken from Ref.<sup>19</sup> with taking into account the short-range correlations and cut-off parameters  $\Lambda_{NN\pi} = \Lambda_{N\Delta\pi} = 0.545$  GeV (set D of Ref.<sup>19</sup>). A similar result comes from the  $T$  matrix of Eq. (2.10) with the same cut-off parameters, see the dashed lines in Fig. 1. At higher energies, one can see that the predictions for these two last models practically coincide, while in the high- $t$  region they differ from the calculation with the cut-off parameters  $\Lambda \simeq 0.7$  GeV. This difference is also seen in the delta Dalitz decay rate in  $pp$  interactions at 1 and 2 GeV shown in Fig. 2. The short dashed lines represent the calculation with a constant  $t$ -weighted  $\Delta$  production cross section in Eq. (2.1) while the other lines correspond to the different  $\Delta$  production  $T$  matrix (with the same notation as in Fig. 1). One can see that at 1 GeV the  $M_\Delta$  dependence of the  $t$ -weighted cross section in Eq. (2.2) suppresses the delta Dalitz decay rate by a factor of 2 at  $M \sim 0.35$  GeV. The differences coming from different  $T$  matrix parametrizations are below 50%. At 2 GeV, the suppression is even larger, a factor of  $\sim 7$  at  $M \sim m_\rho$ . The difference between different models for the  $T$ -matrix may amount up to a factor of 2. Therefore, for a clear understanding of the  $\Delta$  Dalitz decay rate, and more detailed data on the differential delta production cross section at large momentum transfer are necessary. In the subsequent calculations we will use the delta production  $T$ -matrix as in Eq. (2.10) with  $\Lambda_{NN\pi} = \Lambda_{N\Delta\pi} = 0.7$  GeV, which seems to be preferable to reproduce the known experimental data on the delta production cross section.

### III. BREMSSTRAHLUNG DILEPTON PRODUCTION

Dilepton radiation via  $pn$  bremsstrahlung has been extensively studied, cf. Refs.<sup>4,5,10,11,22</sup>. Explicit diagrammatic calculations of the  $pn$  bremsstrahlung are performed on the basis of the one-boson exchange model<sup>10,16,22,23</sup> where four mesons ( $\pi, \sigma, \omega, \rho$ ) are used for the description of the two-body  $pn$   $T$  matrix. It is found that the result depends on the two-body  $T$  matrix parameters which cannot be fixed uniquely only by fitting to the  $pn$  elastic scattering. This method is too complicated to be used as a convenient input in many-body kinetic calculations of dilepton production in nucleus–nucleus collisions. As has been mentioned in the introduction,

to avoid these difficulties, a method based on the soft photon approximation<sup>5</sup> has been used in Refs.<sup>6,7,14</sup>. One should have in mind that the soft photon approximation contains at least three approximations: (i) it keeps only the electric part of the electromagnetic current, (ii) it neglects the radiation from the internal charged meson exchange lines and the nucleon–nucleon–meson vertices, and (iii) contains an approximate integration over unobservable kinematic variables, where the momentum, energy and invariant mass of the virtual photon are assumed to be negligible as compared to the other variables (e.g., the initial and momentum transfer, etc.). The comparison with the exact diagrammatic calculation<sup>10,22,23</sup> shows that the first two approximations change the result not more than a few percent and really may be approved. But the third approximation appears crude. To improve the result, a phenomenological reduction factor has been introduced in Ref.<sup>5</sup>, which is aimed to reduce the remaining phase-space volume for the colliding hadrons in their final state. We must stress that this factor cannot be extracted explicitly from a multidimensional integral, and one should be careful in interpreting the final result within this model, especially at large invariant dilepton masses. For all these reasons, in the present paper we use a model that employs the first two approximations (i, ii) of the soft photon approximation, however, takes into account exact kinematic relations. The net result reads

$$\frac{d\sigma^{ab \rightarrow e^+e^- a'b'}}{dM} = \frac{\alpha^2}{16^2 \pi^4} \frac{1}{M^3 \sqrt{s(s-4m_N^2)}} F_\gamma^2(M) \int dy dq_\perp^2 dE'_b d\varphi \frac{1}{|\mathbf{q}|} J_\mu J_\nu \mathcal{P}^{\mu\nu}, \quad (3.1)$$

where  $\mathcal{P}^{\mu\nu} = -\frac{4}{3}(g^{\mu\nu}q^2 - q^\mu q^\nu)$  is a projector, and  $J^\mu$  is the hadron current. The upper and lower limits of the integral over  $dE'_b$  are defined from the condition

$$|\cos\theta_{\mathbf{q}, \mathbf{p}'_b}| = \left| \frac{s - 2\sqrt{s}q_0 + M^2 - 2(\sqrt{s} - q_0)E'_b}{2|\mathbf{q}||\mathbf{p}'_b|} \right| \leq 1. \quad (3.2)$$

Let us first discuss the structure of the hadron current. For the  $pn$  bremsstrahlung in the soft photon approximation it has the usual form

$$J_{pn}^\mu = -\frac{p_a^\mu}{(p_a q)^2} T(s', t, (p_a - q)^2) + \frac{p'_a{}^\mu}{(p'_a q)^2} T(s, t, (p'_a + q)^2) \quad (3.3)$$

which is gauge invariant in the on-shell limit

$$T(s', t, p^2) = T(s, t, m_N^2), \quad (3.4)$$

with

$$|T(s, t, m_N^2)|^2 = 16\pi s(s-4m_N^2) \frac{d\sigma^{ab \rightarrow a'b'}}{dt}(s, t) \quad (3.5)$$

where  $p_a$  and  $p'_a$  are the four momenta of the initial and outgoing protons;  $T$  is the strong interaction two-body  $T$  matrix;  $t = (p_b - p'_b)^2 = 2m_N^2 - 2E_a E'_b + 2p_a p'_b (\cos \theta_{\mathbf{q}, \mathbf{p}_a} \cos \theta_{\mathbf{q}, \mathbf{p}'_b} + \sin \theta_{\mathbf{q}, \mathbf{p}_a} \sin \theta_{\mathbf{q}, \mathbf{p}'_b} \cos \varphi)$ ,  $s = (p_a + p_b)^2$ ,  $s' = (p_a + p_b - q)^2$ .  $(d\sigma/dt)^{ab \rightarrow a'b'}$  ( $s, t$ ) denotes the elastic  $ab \rightarrow a'b'$  scattering cross section, the symbol  $a$  denotes a proton and  $b$  refers to a neutron.

The electric part of the hadron current in a  $pp$  collision within the soft photon approximation takes the form

$$J_{pp}^\mu = -\frac{p_a^\mu}{(p_a q)^2} T(s', t) - \frac{p_b^\mu}{(p_b q)^2} T(s', t') + \frac{p_a'^\mu}{(p_a' q)^2} T(s, t) + \frac{p_b'^\mu}{(p_b' q)^2} T(s, t'). \quad (3.6)$$

Here  $p_a$  and  $p_b$  are the four momenta of the projectile and target proton, and  $t' = (p_b - p'_b - q)^2$ . It is seen that the hadron current  $J_{pp}^\mu$  (3.6) does not vanish at finite values of  $q_0$  and  $M$ .

One of the still open and interesting questions here is the off-shell corrections to this process. Each of the  $T$  matrices in Eqs. (3.3), (3.6) are far off shell with  $\sum m_i^2 \neq m_N^2$ , where  $m_i$ ,  $i = a, b, a', b'$  is the mass of interacting particles. If we describe the nucleon-nucleon interaction on the basis of the effective one-boson exchange  $T$  matrix, we have to introduce vertex form factors, which, for the on-shell case, depend only on the momentum transfer squared  $t$ . For the "one half" on-shell  $T$  matrix, we have in the bremsstrahlung also effective vertex functions that must depend on an additional invariant variable. The momentum squared  $p^2$  of the off-shell nucleon may be chosen as this variable. For a qualitative analysis we can use also the dimensionless off-shell variable  $\xi = p^\mu p_\mu / m_N^2$  where  $m_N$  is the nucleon mass and  $p^\mu$  denotes the four momentum of the virtual off-shell proton after or before photon radiation. A kinematic analysis shows that, for large values of invariant masses as well as for high energies and momenta of the virtual photon,  $\xi$  is far from its on-shell value  $\xi = 1$ . A phenomenological analysis of the off-shell correction to the effective one boson exchange  $T$  matrix is performed in Ref.<sup>12</sup> where some additional off-shell suppression of the bremsstrahlung rate at higher energy is introduced. This suppression depends on the value of a dimensional cut-off parameter which should range on a typical hadron scale 1–2 GeV. Unfortunately, till now we have not at hand an appropriately well-founded generic theoretical model for the off-shell  $T$  matrices and form factors. In order to avoid in the present consideration such an additional parameter, in our further calculations we use the on-shell model and put in Eqs.(3.3), (3.6)  $t' = t$  and  $s' = s$  which are expected to give an upper estimate of the bremsstrahlung contribution. The procedure of including the off-shell dependence into the electromagnetic form factors and two body  $T$ -matrices discussed in<sup>16</sup>. But

the concrete calculation in<sup>16</sup> is performed with the on-shell form factors and the final results of Ref.<sup>16</sup>, and in our approach are very close to each other.

Fig. 3 shows separately the contribution of the  $pp$  and  $pn$  bremsstrahlung at 1 GeV. One can see that at 1 GeV the  $pp$  contribution is about 30-40% of the  $pn$  bremsstrahlung. For comparison, we also present the result of calculation of the  $pn$  bremsstrahlung within the traditional soft photon approximation<sup>5</sup>:

$$\frac{d\sigma^{pn \rightarrow e^+ e^- p' n'}}{dM} = \frac{\alpha^2}{6\pi^2 M} \int dy dq_1^2 dt \left( \frac{-t}{m_N^2 q_0^2} \right) \frac{d\sigma^{pn \rightarrow p' n'}}{dt}(s, t) \frac{R_2(s')}{R_2(s)}, \quad (3.7)$$

where  $R_2$  is the Lorentz invariant two-body phase-space integral of the final two nucleons of the energy  $\sqrt{s}$ . In calculation of Eq. (3.7) the expression for the  $pn$  elastic cross section is taken the same as in Eq. (3.3). One can see that the soft photon approximation with a phase-space correction results in a twice larger cross section than taking into account the explicit conservation law in calculations.

In all our calculations we use an energy dependent parametrization for the elastic  $pp$  and  $pn$  cross sections. These cross sections decrease with increasing energy, which leads to a decrease of the bremsstrahlung rate in the total dilepton production. In the  $pd$  interaction the  $pn$  and  $pp$  contributions are summed coherently.

#### IV. MESONS DALITZ DECAY

The contribution of the  $\eta$  Dalitz decay takes the form

$$\frac{d\sigma^{\eta \rightarrow e^+ e^- \gamma}}{dM^2} = \sigma_{pd \rightarrow \eta X}(s) \frac{2\alpha}{3\pi M^2} 0.39 F_\gamma^2(M) \left( \frac{\lambda(m_\eta^2, M^2, m_0^2)}{\lambda(m_\eta^2, 0, m_0^2)} \right)^{3/2}, \quad m_0 = 0, \quad (4.1)$$

where the number 0.39 is the branching ratio for the  $\eta \rightarrow \gamma\gamma$  decay<sup>24</sup>. When calculating the  $\eta$  production in the  $pd$  scattering, we use a realistic deuteron wave function  $\phi_d$  obtained within the Paris potential model<sup>25</sup>

$$\sigma_{pd \rightarrow \eta X}(s) = \int \mathcal{R} [\sigma_{pp \rightarrow \eta pp}(s'(k)) + \sigma_{pn \rightarrow \eta pn}(s'(k))] |\phi_d(\vec{k})|^2 d\vec{k}, \quad (4.2)$$

where  $\vec{k}$  is the relative nucleon momentum in the deuteron, and  $\mathcal{R}$  denotes the flux factor. The internal nucleon motion in the deuteron is important near and below the  $\eta$  threshold. We also include short range correlations describing a simultaneous interaction of the proton with a correlated two-nucleon cluster in the deuteron wave function with a 5% probability as in Ref.<sup>26</sup>. Such an effect has been found important

for scattering processes near thresholds and at large angles<sup>26-28</sup>. In the present calculation, the  $\eta$  production cross section is taken in the form:

$$\sigma_{pd \rightarrow \eta X}(s) = (1 - \alpha)\sigma_{pd \rightarrow \eta X}^{(1)}(s) + \alpha\sigma_{pd \rightarrow \eta X}^{(2)}(s), \quad (4.3)$$

where  $\alpha$  is the correlation probability ( $\alpha = 0.05$ ),  $\sigma^{(1)}$  is determined by Eq. (4.2), while  $\alpha\sigma^{(2)}$  is the contribution of the correlated two-nucleon cluster, which we will discuss later.

The cross section for the  $\eta$  production in  $pp$  collisions via the intermediate  $N^*(1535)$  resonance has been studied within the one boson exchange model in Refs.<sup>29,30</sup>. The results of those calculations depend on the input parameters and are different in two papers just near the threshold. The assumption that the threshold behavior of the  $\eta$  production cross section in the collision of particles  $a$  and  $b$  is mainly determined by the phase space integral results

$$\sigma_{ab \rightarrow \eta ab}(s) \sim \sqrt{\frac{m_a m_b (m_a + m_b) m_\eta}{\lambda(s^2, m_a^2, m_b^2)}} \left(1 - \sqrt{\frac{s_0}{s}}\right)^2, \quad s_0 = (m_a + m_b + m_\eta)^2. \quad (4.4)$$

In our calculation, we use an analytical parametrization of  $\sigma_{pp \rightarrow \eta pp}(s)$  motivated by Eq.(4.4) and given by

$$\sigma_{pp \rightarrow \eta pp}(s) = A \frac{(1-x)^2}{\sqrt{\lambda(s, m_p^2, m_p^2)}} \left(1 + \left[\frac{17}{x}(1-x)\right]^\gamma\right)^{-1};$$

$$x = \sqrt{\frac{s_0}{s}}, \quad A = 4 \cdot 10^2 \text{ mb} \cdot \text{GeV}^2, \quad \gamma = 1.8. \quad (4.5)$$

This parametrization gives an average of Refs.<sup>29,30</sup> estimation for the cross section and numerically coincides with the prediction of Ref.<sup>31</sup>.

The  $\eta$ -production cross section in  $pn$ -collisions near the threshold is a subject of some debate. Usually, it is assumed that this cross section, should be scaled, when comparing with the  $pp$  cross section, i.e.  $\sigma_{pn \rightarrow \eta pn} = \kappa \sigma_{pp \rightarrow \eta pp}$ . The possible increasing of  $\sigma_{pn \rightarrow \eta pn}$  reflects the dynamics of the eta production which is beyond our "kinematical" consideration given by Eqs. (4.4,4.5). But if we introduce the enhancement factor  $\kappa$  into  $\sigma_{pn \rightarrow \eta pn}$ , then approximately the same factor ( $\sim (1 + \kappa)$ ) should be included into the cross section of interaction of the proton with a correlated two-nucleon cluster.

The one-boson exchange model prediction of Ref.<sup>29</sup> is  $\kappa \sim 3$ . Estimation of  $\kappa$  within a statistical string model<sup>32,33</sup> gives  $\kappa \sim 1$  in a wide energy range, starting from the threshold. An attempt of a direct extraction of  $\kappa$  from experiment near

the threshold indicates a large value<sup>34</sup>  $\kappa \sim 8-9$ . Unfortunately, we have no real information on the cross section  $\sigma_{pn \rightarrow \eta pn}$ . On the one hand, the energy dependent OBE model parameters of Refs.<sup>29,30</sup> are not fixed from independent experiments. On the other hand, extracting  $\sigma_{pn \rightarrow \eta pn}$  from the nuclear data one has to take carefully into account both the internal motion of the nucleons and short range nucleon-nucleon correlations in nuclei. The latter effect increases the total cross section near the threshold strongly (about one order of magnitude), and this increasing may be described phenomenologically as increasing of  $\kappa$ . Taking into account this indirect knowledge on  $\sigma_{pn \rightarrow \eta pn}$  we adopt for the later the same expression as for the  $pp$  case, Eq. (4.5) with  $A = 3 \cdot 10^3 \text{ mb} \cdot \text{GeV}^2$ ,  $b = 33$  and  $\gamma = 2.1$ . The contribution of the two-nucleon correlation becomes completely negligible at bombarding energy  $E \geq 1.3 \text{ GeV}$ . Eqs. (4.4) and (4.5) give prescription for  $\sigma_{pd \rightarrow \eta pd}^{(2)}(s)$  in Eq. (4.3). This expression has the same form  $\frac{\sqrt{3}}{2}(\sigma_{pp} + \sigma_{pn})$  with substitution:  $\lambda(s, m_p^2, m_p^2) \rightarrow \lambda(2s + m_p^2, 4m_p^2, m_p^2)$ .

At 1 GeV, the  $\eta$ -production cross section with taking into account only the internal nucleon motion is about  $2.6 \cdot 10^{-4} \text{ mb}$ , while with the two-nucleon correlation we obtain  $\sigma_{pd \rightarrow \eta}(s) \approx 5.1 \cdot 10^{-3} \text{ mb}$ . This strong effect of the subthreshold  $\eta$ -production is seen in the dilepton distributions at initial energy 1 GeV at large invariant masses near the kinematical limit. In this case, the contribution of the  $\eta$ -Dalitz decay is comparable with the contribution of the  $\Delta$ -Dalitz decay and  $pd$  bremsstrahlung, and is seen but is not dominant. So, we find that the total dilepton invariant mass distribution in  $pd$  collisions at 1 GeV is not very sensitive to the large uncertainty of the eta production cross section in  $pn$ -collision near the threshold. For higher energies ( $\sim 5 \text{ GeV}$ ), we have to take into account the total inclusive eta production cross section that is larger than the exclusive cross section discussed above. At 4.9 GeV, we use the upper limit for the eta production cross section<sup>13</sup>: 0.5 mb with  $\sigma_{pp} = \sigma_{pn}$ , our choice of  $\kappa = 1$  corresponds to the prediction which is in agreement with the statistical string model<sup>32</sup>.

Estimates of the  $\omega$  Dalitz decay may be performed on the basis of Eq. (4.1) with the substitutions  $\sigma_\eta \rightarrow \sigma_\omega$ ,  $m_\eta \rightarrow m_\omega$ ,  $m_0 \rightarrow m_\pi$ ,  $2\alpha \rightarrow \alpha$  and  $0.39 \rightarrow 0.08$ . Using the known experimental data on the  $\omega$  production cross section<sup>35</sup>, we find that the contribution of the  $\omega$  Dalitz decay to the dilepton production is several orders of magnitudes smaller than the contribution of other subprocesses even at 5 GeV.

We do not consider the Dalitz decay of pions because it contributes to the low invariant mass region  $M \leq m_\pi$  not investigated here.

## V. DIRECT DECAY OF THE VECTOR MESONS

In principle, it is difficult to distinguish between the “pure” bremsstrahlung discussed in Sect. 3 with the intermediate  $\rho$  meson formation and the contribution of the direct vector meson decay to the lepton pair with the same final states. But as has been discussed above, because of the energy dependence of the two-body  $T$  matrix (or nucleon–nucleon elastic scattering cross section) the contribution of pure bremsstrahlung decreases with increasing kinetic energy while the  $\rho, \omega$  production cross sections increase. To avoid a double counting problem, we use a simple recipe that at low energy  $E \leq 3$  GeV only the pure bremsstrahlung, defined in Sec. 3, contributes while for energies larger than the  $\rho, \omega$  production threshold the bremsstrahlung is considered only as a subprocess of the direct decay of vector mesons. To be correct, this mechanism and the Dalitz delta decay channel should be summed coherently. An exact solution of this complicated problem on the basis of microscopic decomposition of the subprocesses without double counting is beyond the scope of the present consideration. Instead, we put the strong interaction part in a vector meson production cross section and consider its decay into the dielectron channel. This is an analog to the handling of the delta Dalitz decay contribution in Sec. 2 which is commonly used<sup>8,9,36</sup>.

The cross section of the vector meson decay into an electron pair may be written in the following form:

$$\frac{d\sigma^{V \rightarrow e^+e^-}}{dM^2} = \sum_{i=\rho,\omega} \sigma_{V_i}(s) \frac{1}{\pi} \frac{m_i \Gamma_i(m_i)}{(M^2 - m_i^2)^2 + (m_i \Gamma_i(M))^2} B_i \left( \frac{m_\rho}{M} \right)^3, \quad (5.1)$$

where  $\Gamma_i(M)$  is the mass-dependent total width of the  $i^{\text{th}}$  vector meson decay,  $B_i$  is the known electromagnetic decay branching ratio<sup>24</sup>, the cross section  $\sigma_{V_i}(s)$  is the total vector meson production cross section. For the narrow  $\omega$ -resonance, one can take  $\Gamma_\omega(M) \simeq \Gamma_\omega(m_\omega)$ . An interaction Lagrangian for the  $\rho\pi^+\pi^-$ -vertex

$$\mathcal{L}_{\rho\pi^+\pi^-} \sim g_{\rho\pi\pi} \pi^+ \pi^- (k_{\pi^+} - k_{\pi^-})^\mu \quad (5.2)$$

leads to the mass-dependent  $\rho$ -decay width  $\Gamma_\rho \simeq \Gamma_{\rho \rightarrow \pi\pi}$

$$\Gamma_\rho(M) = \Gamma_\rho(m_\rho) \frac{M}{m_\rho} \left[ \left(1 - \frac{4m_\pi^2}{M^2}\right) / \left(1 - \frac{4m_\pi^2}{m_\rho^2}\right) \right]^{3/2}. \quad (5.3)$$

Dileptons with invariant masses  $M < 2m_\pi$  are “produced” by the virtual  $\rho$ -mesons with zero width. The cross section for the  $\rho, \omega$  production near the threshold has

been studied within the Nambu–Jona-Lasinio model in Ref.<sup>37</sup>. In our calculation, we use the predicted value  $\sigma_{\rho,\omega}(E = 4.9\text{GeV}) \sim 0.09$  mb which is in agreement with available experimental data<sup>35</sup>.

## VI. RESULTS

In Figs. 4-7 we display our results. Invariant mass spectra for the  $pd$  reactions without and with an experimental filter are displayed in Figs. 4, 5. The acceptance Dilepton Spectrometer Collaboration (DLS) filter we have used Version 2.0. The filter suppresses the dilepton yield at all invariant masses, and the resulting suppression is different for different subprocesses because of different kinematic conditions and kinematical limits in each channel. For comparison, we also display in Fig. 5 the results of the DLS collaboration<sup>3</sup> for the  $p^9\text{Be}$  interaction scaled by a factor  $A_{p^9\text{Be}}^{-2/3}$ . If we assume that the absorption of an initial proton in a nucleus is proportional to  $A^{-1}$ , then the  $A$ -dependence should be  $A^{2/3}$ . The result of calculation in Ref.<sup>39</sup> shows that the absorption factor for Beryllium numerically coincides with  $A^{-1}$ . This means that one can expect that the dielectron production cross section for the  $p^9\text{Be}$  interaction, scaled by a factor of  $A^{-2/3}$ , may be considered approximately as dilepton production in a  $p$ -isoscalar nucleon interaction. Other medium effects (excluding internal motion) in the light Beryllium nucleus are expected to be negligible. Therefore, one must consider the scaled Beryllium data as some rough guide of what to be expected by proper  $pd$  data at 1-2 GeV. We do not attempt a fine tuning of our input to reproduce exactly the scaled data.

In calculating the  $\Delta$  Dalitz decay and bremsstrahlung contribution in  $pd$  reactions, we also take into account the internal motion of nucleons in the deuteron, as in Eq. (4.2). The integrations in Eqs. (2.1, 3.1, 4.2) are performed by a Monte Carlo method. The cross section of the elastic scattering in Eq. (3.1) is parametrized to reproduce the experimental data at each initial energy separately.

Our results for 1 and 2 GeV without the DLS filter are close numerically to the results of our previous paper<sup>12</sup>, where the above-mentioned off-shell suppression has been used. A small difference is explained by different parametrization of the  $\eta$ -production cross section (in Ref.<sup>12</sup> the prediction of Ref.<sup>29</sup> was used) and taking into account in Ref.<sup>12</sup> the phenomenological off-shell correction in the two-body  $T$  matrix in the bremsstrahlung channel. In Ref.<sup>12</sup>, antisymmetrization in the  $pd$  collision was overestimated. At 2 GeV, the bremsstrahlung is not a dominant source of the dileptons, and the off-shell effect is not seen in the total cross section. At 1 GeV,

the effect of the off-shell correction is much smaller. So, concerning the off-shell effects our present result may be considered as an upper limit for the bremsstrahlung contribution.

In Figs. 4 and 5 we present the result of calculation without and with the DLS acceptance filter. As a matter of fact, our study shows that the influence of the DLS filter is much stronger than the off-shell corrections discussed in Ref.<sup>12</sup> and it is strictly necessary to take the filter into account for a correct comparison with experimental data.

At 1 GeV, we find the bremsstrahlung contribution nearly as strong as the  $\Delta$  Dalitz decay. There is also a contribution of the  $\eta$  Dalitz decay. Near the threshold the  $\eta$  decay gives the same (or even) larger contribution than the  $\Delta$  Dalitz decay and bremsstrahlung. Subthreshold effects are responsible for larger invariant mass tails of the  $\Delta$  and bremsstrahlung contributions in the  $pd$  reactions. The vector dominance effects (i.e., the form factor) are not important.

On the contrary, at 2.1 GeV the VDM effect is important. However, the strong enhancement of the  $\Delta$  Dalitz decay at the  $\rho$  peak is reduced by the  $t$ -dependence of the delta production matrix element. The net result is a shoulder in the sum of all contributions. This is not so clearly seen in  $pp$  reactions due to the kinematic limit. However, in  $pd$  reactions, due to subthreshold effects it can be observed. Our net results are of the same order of magnitude as those obtained in Ref.<sup>6</sup>, but in Ref.<sup>6</sup> there is no shoulder behavior in the  $\rho$  region. In the intermediate invariant mass region  $0.2 \leq M \leq 0.4$  GeV the  $\eta$  Dalitz decay gives the main contribution. The contribution of the  $\omega$  Dalitz decay is very small and is not displayed here.

At 4.9 GeV, the available data of the dielectron production in the  $pd$  collision<sup>2</sup> are shown by open circles. One can see that even a peak appears at the  $\rho, \omega$  position which coincides with the predictions of Refs.<sup>8,9,39</sup>. Contrary to other investigations (e.g.,<sup>9</sup>) we get also a prominent delta decay contribution due to the vector dominance form factor which determines the broad peak width while the  $\rho, \omega$  contributions give a sharp peak on the delta bump. This broad peak has been predicted also in Ref.<sup>13</sup> within a more simplified model for the  $\Delta$  Dalitz decay channel though with a too large value for the delta production cross section. For more detailed comparison with the data at 4.9 GeV, one has to take into account also other sources of dielectrons: bremsstrahlung with the multipion final states, pion annihilation etc.<sup>13</sup>. A similar analysis of the dilepton production in the proton nucleons collision with taking into account proton nucleon bremsstrahlung and the effect of propagating the  $\Delta$  resonance ( $\Delta$  Dalitz decay) has been performed in<sup>16</sup>. The principal results

of those channels in Ref.<sup>16</sup> and in our study coincide, i.e., the main contribution to the dilepton spectrum comes from the  $\Delta$  decay. But there is some difference in the interpretation of the  $pd/pp$  ratio. In Ref.<sup>16</sup>, some enhancement of the ratio at lower energies may be explained by (i) different values of the  $\Delta$  production cross section (in Ref.<sup>16</sup>) this difference is of the factor 2–3) and (ii) relatively large destructive interference between bremsstrahlung and the  $\Delta$  decay channel in  $pp$  as compared to the  $pn$  collisions. In our model, the first effect exists but its contribution is smaller. The difference between the  $\Delta$  production cross section is controlled by the Ver West-Arndt parametrisation and it is of a factor of  $\sim 1.7$  at  $E=1.2$  GeV. The second effect is dropped here, however, we take into account the  $\eta$  decay contribution.

Now let us consider the ratio of the cross section for  $pd$  to  $pp$  reactions

$$R = \frac{d\sigma^{pd}/dM}{d\sigma^{pp}/dM}, \quad (6.1)$$

which is displayed in Fig. 6 for three energies. The experimental data are taken from Ref.<sup>2</sup>. The difference between  $pp$  and  $pd$  interactions consists in (i) taking into account the internal nucleon motion in a deuteron, (ii) different expressions for the  $\Delta$  production cross section which follow from the Ver West-Arndt parametrization in Eq. (2.12), and (iii) absence of the  $\eta$  decay contribution in the  $pp$  case above the threshold. If the bremsstrahlung contribution in the  $pp$  and  $pd$  collision were switched off and the eta production cross section in the  $pn$  reaction is taken equal to the cross section in the  $pp$  collision, the ratio would be energy-independent and close to 2 except for the vicinity of the kinematic  $pp$  threshold. It is seen that the ratio rises towards the kinematic limits due to phase space limitations in the  $pp$  reactions. Except for this boundary behavior, the ratio decreases towards 2 with increasing initial energy, which reflects decreasing of  $\kappa = \sigma_{pn \rightarrow \eta} / \sigma_{pp \rightarrow \eta}$  towards 1 and relative decreasing of the short range correlation effect responsible for the subthreshold  $\eta$  production elastic  $NN$ -scattering cross section. At 1.26 GeV, the result is sensitive to the subthreshold  $\eta$  production mechanism. One can see that in spite of qualitative agreement with data, the theoretical prediction is twice smaller at  $M \sim 0.2 - 0.3$  GeV than the data. Our analysis shows that for the  $pp$ -collision at 1 GeV only the  $\Delta$  Dalitz and  $pp$  bremsstrahlung contribute. Theoretical uncertainties in the  $pp$  interaction are minimal because the total and differential  $\Delta$ -production cross section at 1 GeV are well known, as well as the elastic  $pp$  cross section operating in bremsstrahlung. So, the calculated dielectron production cross section for the  $pp$ -collision being multiplied by an experimentally measured quantity  $R$  results in an estimation for dielectron yield in the  $pd$ -collision at 1 GeV (cf. Eq. (6.1)). The

corresponding points are shown in Fig. 5 (open circles).

The Fig. 7 shows the ratio of the integrals

$$R_{\text{int}} = \frac{\int_{0.1 \text{ GeV}}^{M_{\text{max}}} (d\sigma^{pd}/dM)dM}{\int_{0.1 \text{ GeV}}^{M_{\text{max}}} (d\sigma^{pp}/dM)dM} \quad (6.2)$$

as a function on the initial energy  $E$  with the DLS filter. One can see some enhancement of the ratio at  $E < 1.4$  GeV because of the large sub-(and near-) threshold  $\eta$  Dalitz decay contribution in the  $pd$  collisions and  $pn$  bremsstrahlung contribution. Again, one can see the difference of the factor 2 between prediction and the data and the origin of this difference is the same as in Fig. 6. Then, the ratio goes to 2 as the contribution of the main channels in the  $pp$  and  $pn$  collision becomes the same.

## VII. SUMMARY

In summary, we present a detailed analysis of dielectron production in the  $pp$  and  $pd$  reaction at 1–5 GeV and find qualitative agreement with available experimental data. Our model relies on vector dominance, improves the soft photon approximation, and uses the correct  $\Delta$  production cross section.

We can conclude that the dilepton production cross section is sensitive to the very details of the elementary subprocesses which have been analyzed. The accuracy of the  $\Delta, \eta, \omega$  Dalitz and direct  $\rho, \omega$  decays depends on the knowledge of the unstable hadron production mechanisms. So, new precision measurements in the Bevalac, SIS, COSY energy region are needed. Also an independent verification of the two-body  $T$  matrix off-shell behavior and time-like nucleon form factor is needed. Only a clear understanding of the dilepton production in  $NN$  interaction can give a reliable possibility to use dileptons as an accurate probe for a more complex nuclear collision.

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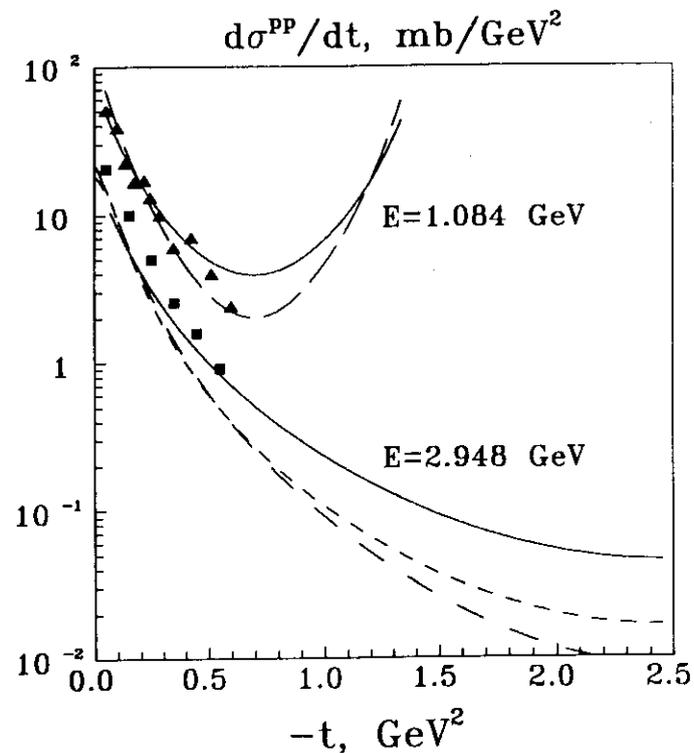


FIG. 1. Elastic  $\Delta$ -production cross section for the  $pp \rightarrow \Delta^{++}n$  reaction at bombarding energies 1.084 and 2.948 GeV. Experimental data at 1.084 (thringles) are taken from Ref.<sup>17</sup>, while data at 2.948 GeV (squares) are from Ref.<sup>18</sup>. The solid lines correspond to calculations with  $T$ -matrix (2.10) and  $\Lambda_{NN\pi} = \Lambda_{N\Delta\pi} = 0.7$  GeV. The curves representing calculations at 2.948 GeV correspond to the usage of the  $T$ -matrix from Ref.<sup>19</sup> with  $\Lambda_{NN\pi} = \Lambda_{N\Delta\pi} = 0.545$  GeV (long dashes) and to the exploiting of the  $T$ -matrix (2.10) with  $\Lambda_{NN\pi} = \Lambda_{N\Delta\pi} = 0.545$  GeV (dashes). At 1.084 GeV both latter curves practically coincide (long long dashes).

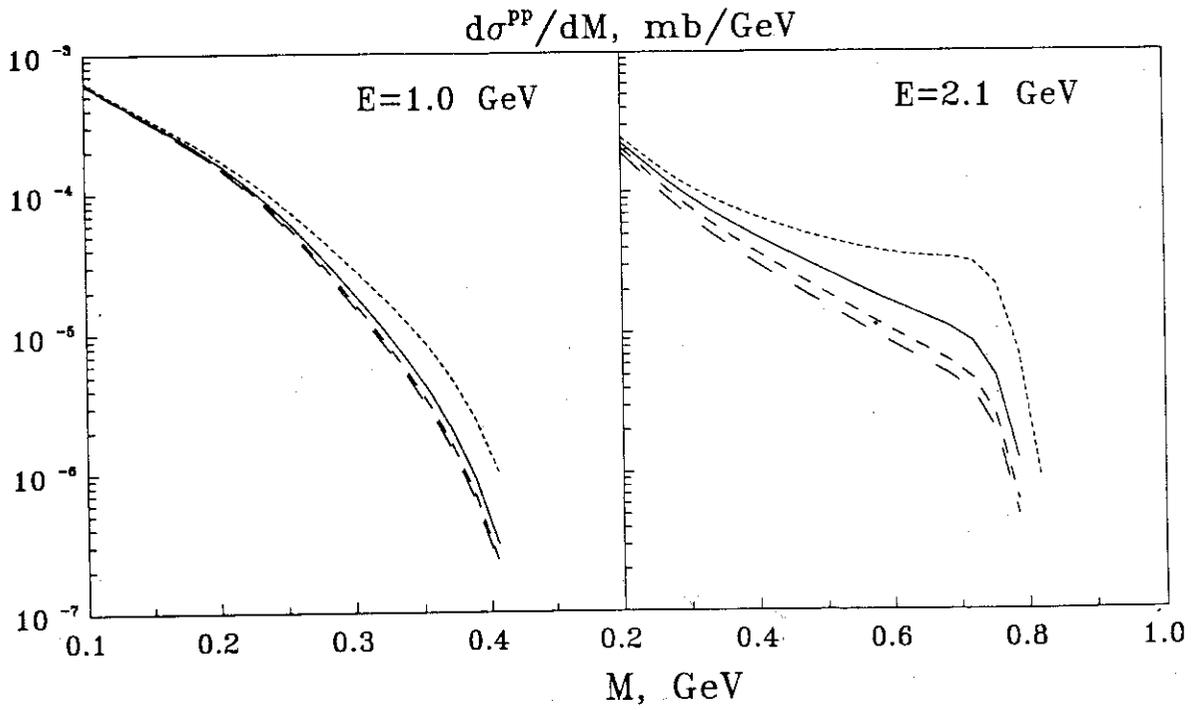


FIG. 2. The Dalitz delta decay contribution to the dielectron production at  $E = 1$  and  $2.1$  GeV. The results of calculations with the constant production cross section  $\bar{\sigma}_{\Delta}(s, M_{\Delta})$  in Eq. (2.1) are represented by the lines with short dashes while other lines correspond to different  $\Delta$  production  $T$ -matrices. Notation is the same as in Fig. 1.

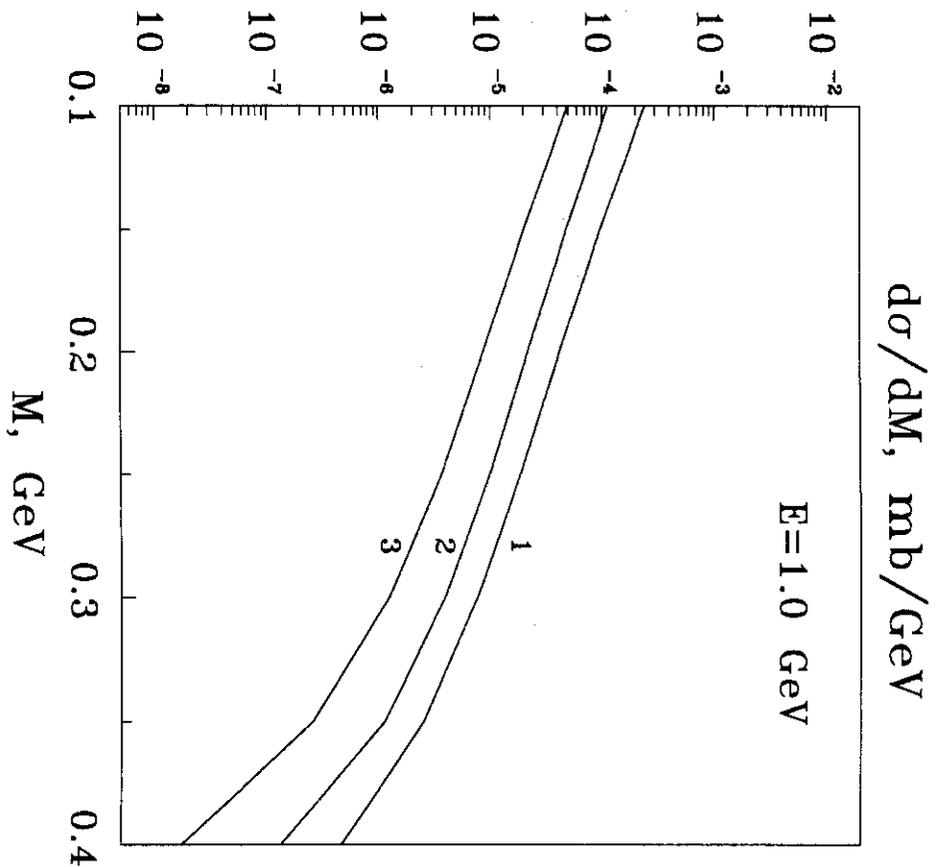


FIG. 3. Contributions of the  $pn$  (curve 2) and  $pp$  (curve 3) bremsstrahlung to the dielectron production at  $1$  GeV. The curve 1 represents a contribution of the  $pn$  bremsstrahlung calculated within the soft photon approximation [see Eq. (3.7)].

$d\sigma/dM, \text{ mb/GeV}$

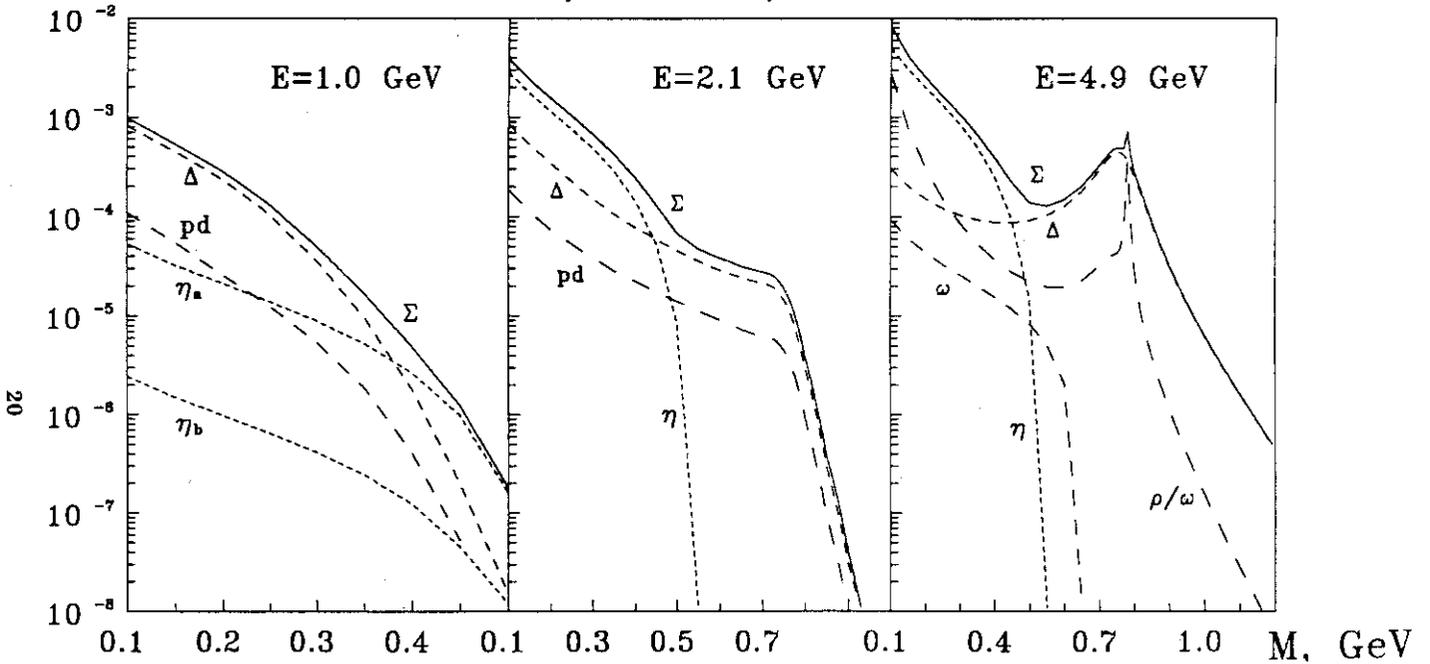


FIG. 4. Dielectron invariant mass spectra for the  $pN$ -collision ( $pN = pd/2$ ) at  $E = 1.0, 2.1, 4.9$  GeV calculated without the DLS filter. The “ $pd$ ” labels bremsstrahlung, “ $\eta$ ”, “ $\Delta$ ” and “ $\omega$ ” denote the corresponding Dalitz decay contributions, “ $\rho/\omega$ ” is the direct rho-omega decay, “ $\Sigma$ ” is the sum of all contributions. The line  $\eta_b$  at 1 GeV represents the subthreshold  $\eta$  decay contribution with taking into account the internal nucleon motion in a deuteron; the line  $\eta_a$  shows calculations with the two-nucleon short-range correlation.

$d\sigma/dM, \text{ mb/GeV}$

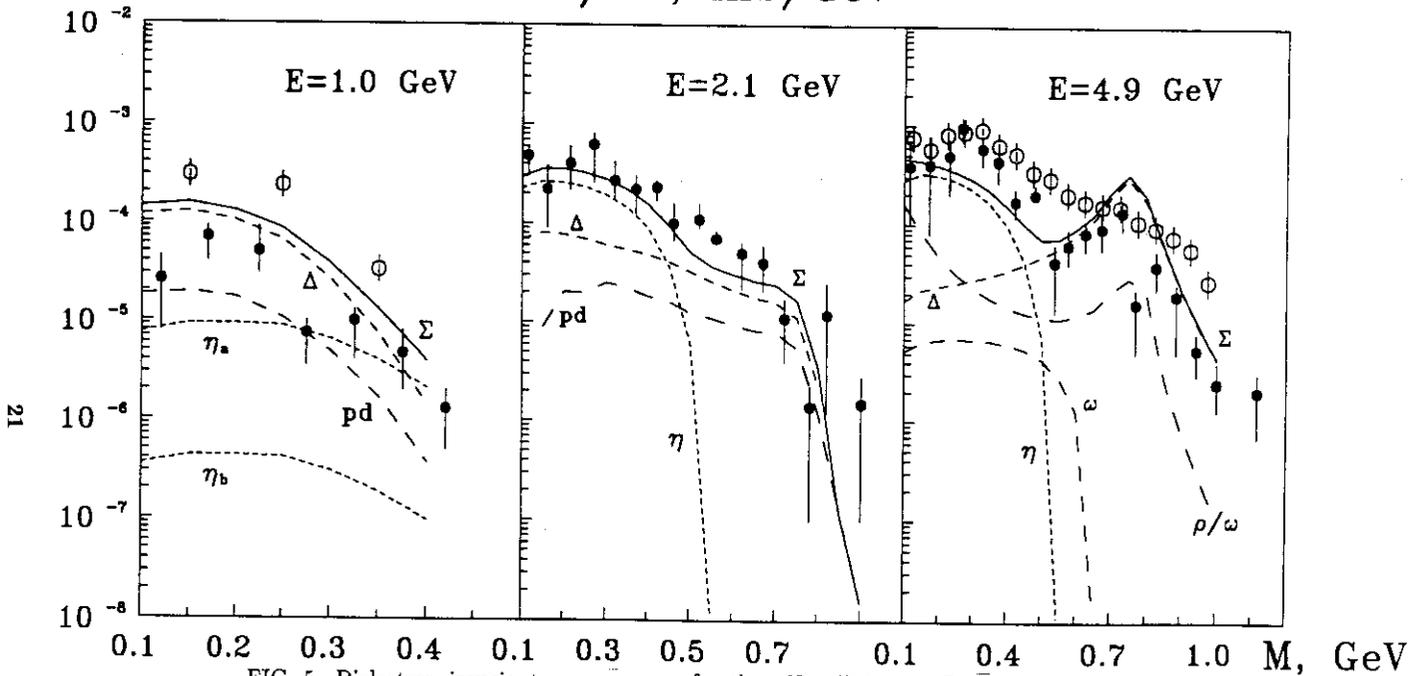


FIG. 5. Dielectron invariant mass spectra for the  $pN$ -collision at  $E = 1.0, 2.1, 4.9$  GeV calculated with the DLS filter. Notations for the curves are the same as in Fig. 4. Experimental data for  $p^9\text{Be}$  collisions, scaled by the factor  $A^{-2/3}$ , are taken from Ref.<sup>3</sup> (solid circles). Open circles at 1 GeV represent the product of calculated dielectron yield in  $pp$  collisions times the ratio of the  $pd$  to  $pp$  dielectron production (see the text). Open circles at 4.9 GeV are the experimental data taken from Ref.<sup>2</sup>.

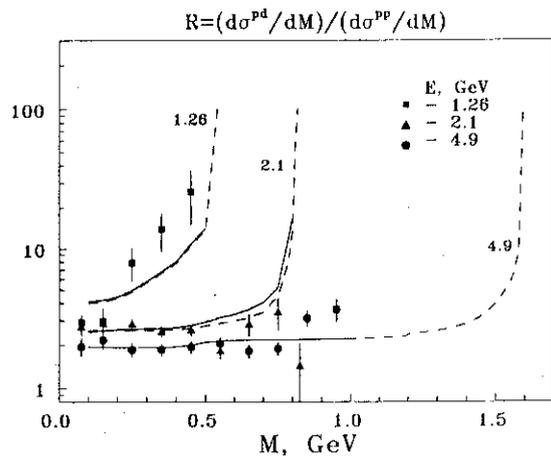


FIG. 6. Ratio of the cross section for  $pd$  to  $pp$  reactions at  $E = 1.26, 2.1, 4.9$  GeV. The preliminary experimental data are taken from Ref.<sup>2,40</sup>. Dashed and solid lines correspond to the calculation with and without of the DLS filter, respectively.

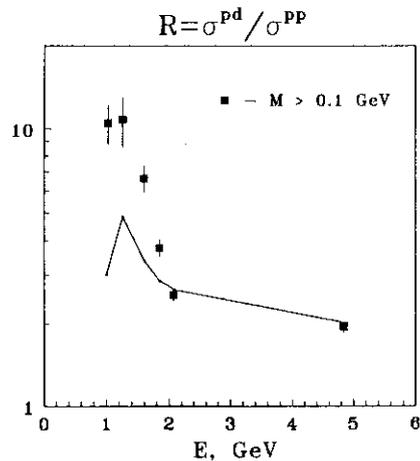


FIG. 7. Ratio of the integrated cross section for  $pd$  to  $pp$  reactions as function of the initial energy. Solid line corresponds to the calculation with the DLS filter. The preliminary experimental data are taken from Ref.<sup>2,40</sup>.

## REFERENCES

- <sup>1</sup> H.Z. Huang et al. Phys. Lett. **B297**, 233 (1992).
- <sup>2</sup> W.K. Wilson et al. Phys. Lett. **B316**, 245 (1993); Preprint LBL 33625, Berkeley, 1993.
- <sup>3</sup> G. Roche et al. Phys. Rev. Lett. **61**, 1069 (1988); Phys. Lett. **B226**, 228 (1989); C. Naudet et al. Phys. Rev. Lett. **62**, 2652 (1989).
- <sup>4</sup> C. Gale, and J. Kapusta. Phys. Rev. **C35**, 2107 (1987).
- <sup>5</sup> C. Gale, and J. Kapusta. Phys. Rev. **C40**, 2397 (1989).
- <sup>6</sup> G. Wolf, G. Batko, W. Cassing, U. Mosel, K. Niita, and M. Schäfer. Nucl. Phys. **A517**, 615 (1990).
- <sup>7</sup> L. Xiong, Z.G. Wu, C.M. Ko, and J.Q. Wu. Nucl. Phys. **A512**, 772 (1990).
- <sup>8</sup> K.K. Gudima, V.D. Toneev, and A.I. Titov. Phys. Lett. **B287**, 302 (1992).
- <sup>9</sup> L.A. Winckelmann, H. Stöcker, and W. Greiner. Phys. Rev. Lett. **B298**, 22 (1993).
- <sup>10</sup> M. Schäfer, T.S. Biro, W. Cassing, and U. Mosel. Phys. Lett. **B221**, 1 (1989).
- <sup>11</sup> K. Haglin, J. Kapusta, and C. Gale. Phys. Lett. **224**, 433 (1989); **221B**, 1 (1989).
- <sup>12</sup> B. Kämpfer, A.I. Titov, and E.L. Bratkovskaya. Phys. Lett. **B301**, 123 (1993).
- <sup>13</sup> K. Haglin and C. Gale. Preprint McGill/93-9 TPI-MINN-93/18-T.
- <sup>14</sup> G. Wolf, W. Cassing, and U. Mosel. Nucl. Phys. **A545**, 139c (1992).
- <sup>15</sup> J.J. Sakurai. *Currents and Mesons*, University of Chicago Press, Chicago, 1969.
- <sup>16</sup> M. Schäfer, H.C. Dönges, A. Engel and U. Mosel. *Dilepton Production in Nucleon-Nucleon Interaction*. Preprint Inst. Theor. Phys. University of Giessen, Germany, 1994.
- <sup>17</sup> A.M. Eisner et al. Phys. Rev. **B138**, 170 (1965).
- <sup>18</sup> S. Colletti et al. Nuovo Cim. **49**, 479 (1967).
- <sup>19</sup> G.F. Bertsch, G.E. Brown, V. Koch, and B.A. Li. Nucl. Phys. **A490**, 745 (1988).
- <sup>20</sup> H.F. Jones, and M.D. Scadron. Ann. of Phys. **81**, 1 (1973).
- <sup>21</sup> B.J. Ver West, and R.A. Arndt. Phys. Rev. **C25**, 1979 (1982).

- <sup>22</sup> K. Haglin. *Ann. of Phys.* **212**, 84 (1991).
- <sup>23</sup> E.L. Bratkovskaya, B.L. Reznik, and A.I. Titov. *Sov. J. Nucl. Phys.* **55**, 3058 (1992).
- <sup>24</sup> *Particle Data Group*. *Phys. Lett.* **B239** (1990).
- <sup>25</sup> M. Lacombe et al. *Phys. Rev.* **C21**, 861 (1980).
- <sup>26</sup> L.P. Kaptari, B.L. Resnik, and A.I. Titov. *Sov. J. Nucl. Phys.* **42**, 777 (1985);  
A.V. Efremov. *Sov. J. Part. Nucl.* **13**, 254 (1982);  
V.V. Burov et al. *Z. Phys.* **A306**, 149 (1982).
- <sup>27</sup> H. Pirner, and J.P. Vary. *Phys. Rev. Lett.* **46**, 1376 (1981).
- <sup>28</sup> T. Fujita. *Phys. Rev. Lett.* **39**, 174 (1977); *Phys. Lett.* **B72**, 16 (1977).
- <sup>29</sup> T. Vetter, A. Engel, T. Biro, and U. Mosel. *Phys. Lett.* **B263**, 153 (1991).
- <sup>30</sup> J.M. Laget, F. Wellers, and J.F. Lecomte. *Phys. Lett.* **B257**, 254 (1991).
- <sup>31</sup> A.L. De Paoli, K. Niita, W. Cassing, U. Mosel, and C.M. Ko. *Phys. Lett.* **B219**, 194 (1989).
- <sup>32</sup> N.S. Amelin et al. *Sov. J. Nucl. Phys.* **52**, 272 (1990).
- <sup>33</sup> V.D. Toneev, A.I. Titov, and K.K. Gudima. Preprint GSI-92-05 (1992).
- <sup>34</sup> E. Chiavassa et al. *Z. Phys.* **A344**, 345 (1993); *Phys. Lett.* **B** (to be published).
- <sup>35</sup> CERN-HERA 84-01 (data compilation), 17 April 1984, unpublished.
- <sup>36</sup> P. Koch. *Phys. Lett.* **B288**, 187 (1992); Preprint University of Regensburg TPR-92-18, 1992.
- <sup>37</sup> E.L. Bratkovskaya, and A.I. Titov. *Phys. Lett.* **B282**, 127 (1992).
- <sup>38</sup> A.I. Titov, B. Kämpfer, and E.L. Bratkovskaya. *Phys. Lett.* **B301**, 123 (1993).
- <sup>39</sup> L.P. Kaptari, V.K. Lukyanov, B.L. Resnik, and A. Titov. In: *Relativistic Nuclear Physics and Quantum Chromodynamics*. Ed. A.M. Baldin et al. World Scientific, 1991, p. 355.
- <sup>40</sup> G. Roche. Preprint PCCF RI 9316, Université Blaise Pascal, 1993.

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