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Anisotropic Distribution of Secondaries in Extreme Energy Cosmic-Ray Stars

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The connection between the conservation law for angular momentum and the anisotropy observed in cosmic-ray stars produced by energetic primaries is investigated. A model is proposed which is probably reasonable for noncentral nucleus-nucleus collisions and may have some meaning for nucleon-nucleon collisions; it is based on assuming that the disk-shaped "interaction volume" consists of many zones acting independently. If equilibrium is reached within individual zones, the particles emerging from one zone ought to be distributed isotropically in its center-of-mass frame. In a noncentral collision these frames will in general, in accordance with the conservation of momentum, be different for different zones. Some anisotropy in the total angular distribution follows, but on this model it could never reach the degree frequently observed.

I. INTRODUCTION

IN Fermi's statistical theory of high-energy processes[1] it is postulated that the outgoing particles correspond to a state of equilibrium within the interaction volume V. Once the particles have left V their coupling is supposed to cease. The number and angular distribution of the particles depends on V and on the energy and angular momentum deposited therein. Fermi took V to be independent of the initial impact parameter, more especially to be a flattened ellipsoid of revolution with its axis, the length of which depends on the collision energy, parallel to the direction of approach of the particles. This simplification seems reasonable in the domain of small impact parameters; but as these approach the transverse radius of the ellipsoid, the predicted measure of anisotropy tends to infinity. This is clearly brought about by the simplifying geometrical assumptions, and one must ask oneself whether, on replacing them by less crude ones, the theory will still account for the considerable anisotropy found in recent experiments.[2,3] One might try to check this by having V depend on the impact parameter (increasing with it) while otherwise taking over Fermi's whole formalism. We are indeed convinced that on doing this and admitting even a rather weak increase of the transverse dimensions of V one would find the peaking predicted for large impact parameters greatly reduced. However, as Fermi has pointed out, it is difficult to specify shape and size of V in a realistic and unambiguous way. Therefore we have tried to find another method, not dependent on precise assumptions concerning V, and which may yet enable us to estimate what *maximum* anisotropy can reasonably be explained from the assumption of a flattened interaction volume and the conservation laws.

To avoid misunderstanding we should here say that we do *not* question Fermi's basic principle: we definitely follow him in asserting that the conservation of angular momentum will, for a flat interaction volume result in some anisotropy of the total angular distribution. However, we think we can show that this anisotropy may be very much weaker than on Fermi's original estimate. Consequently we doubt that an explanation of the observed, strong anisotropy is possible on this basis.

Below we carry out an investigation of a model similar to Fermi's but based on the hypothesis of "local interactions." (This is explained below.) We find that this model gives only negligible "peaking" (even for the largest impact parameters), although all conservation laws are satisfied.

II. EVIDENCE OF ANISOTROPY

We are concerned with the angular distribution in cosmic-ray stars produced by primaries with an energy of a few thousand Bev or greater. The shower or jet part of such a star consists of tracks which do not allow identification of the particles. The only data with direct bearing on the interaction are in most cases the angles between the tracks in the laboratory frame. These are not sufficient to determine the center-of-mass frame (c.m.f.) of the process. For the majority of observed jets there is no Lorentz-frame in which the angular distribution is spatially isotropic. Usually it proves possible to find a Lorentz-frame in which the distribution is more or less symmetric with respect to a plane normal to the path of the primary. It is often thought that this amounts to finding the c.m.f. In most cases, however, it is known that the colliding entities are by no means identical (nucleon-nucleus collision). It is then unlikely, *a priori*, that the angular distribution in the real c.m.f. should exhibit forward-backward symmetry. One could only, *a posteriori*, admit this symmetry if it were found that the distributions are isotropic in the c.m.f. But this is, experimentally, not the case. Therefore this procedure for guessing the c.m.f. and all conclusions based on it should be regarded with some caution.

A convenient measure for the anisotropy has been

[1] E. Fermi, Progr. Theoret. Phys. **5**, 570 (1951); Phys. Rev. **81**, 863 (1951).
[2] M. F. Kaplon and D. M. Ritson, Phys. Rev. **88**, 386 (1952).
[3] Lal, Pal, Peters, and Swami, Proc. Indian Acad. Sci. **36**, 75 (1952).

introduced by Kaplon and Ritson.[2] Let ϕ_n denote the half-angle of the cone in the laboratory frame in which the tracks of nN of all the N particles in the jet are situated. Form the ratio $\phi_{3/4}/\phi_{1/4}$ for the observed distribution and, for comparison, for a distribution isotropic in the same c.m.f. The quantity,

$$X=\frac{(\phi_{3/4}/\phi_{1/4})_{\text{observed}}}{(\phi_{3/4}/\phi_{1/4})_{\text{isotropic}}}, \quad (1)$$

is equal to 1 for distributions for which a Lorentz-frame exists in which they are isotropic. For distributions peaked in forward and backward direction in that Lorentz-frame in which they are symmetrical with respect to the transverse plane, X is greater than 1. A large X manifests itself by the appearance of a "double cone" in the laboratory frame. The wider the outer, diffuse cone and the narrower the superimposed, inner, dense cone is, the larger is X. In the cases of interest the Lorentz-transformation relating the laboratory frame and the so-called c.m.f. is strong ($\gamma\gg 1$) so that, e.g., $\phi_{3/4}$ is still a small angle. For a distribution isotropic in the c.m.f. one has under these conditions, as can be proved with the help of Eqs. (5) and (8) below $\phi_n/\phi_{1-n}=n/(1-n)$. (Here n stands for any fraction not too near 0 or 1.) Putting $n=3/4$, one finds $(\phi_{3/4}/\phi_{1/4})_{\text{isotropic}}=3$. Therefore one can in practice evaluate X from

$$X=\phi_{3/4}/3\phi_{1/4}. \quad (2)$$

For the twenty events in copper described by Kaplon and Ritson[2] the quantity X has values varying from 1.17 to 66, and X is in this set of events not significantly correlated with either the energy or the multiplicity. The Lord-Fainberg-Schein star[4] and the very large star with an Mg primary described by Lal *et al.*[3] both show $X\approx 3$. Two events in copper with α primaries observed by Kaplon and Ritson[2] give $X=1.17$ and $X=2$. The Bradt-Kaplon-Peters star[5] has $X=2.33$. In Fermi's original paper[1] he obtained for median impact parameter a distribution in agreement with that of the Lord-Fainberg-Schein star, i.e., with an $X\approx 3$.

Since nothing is known about the interaction laws, there is no very convincing reason why peaking rather than isotropy should be regarded as surprising and as requiring a *special* explanation. The *a priori* motivation for requiring such an explanation is different for the different authors. Without going into too much detail, we should mention that besides Fermi's argument involving the conservation of angular momentum several other explanations of the anisotropy have been proposed. A number of investigators have put forward the most obvious idea that the double cone may be due to multiple scattering inside the target nucleus; those particles which are scattered more often being deflected more strongly from the direction of the primary. This explanation does away with the assumption that the jet is generated in a single act, and it denies the special role of either the real c.m.f. of the jet or of the frame (probably not identical with the c.m.f.) in which the distribution is most symmetric in backward and forward direction. Arguing from the opposite viewpoint that the peaking is a fundamental feature of interactions at high energies, already present and perhaps most prominent in the results of simple nucleon-nucleon collisions, Heisenberg[6] has pointed out that it might be explained in terms of the coherence of the waves going out from the flat-shaped zone of interaction.

III. MODEL OF THE INTERACTION

Fermi and Heisenberg have pictured "extreme-relativistic" ($E\gg Mc^2$) nucleons as disk-shaped objects, flattened in their direction of motion according to the ratio Mc^2/E. In dealing with the collision of two such disks and the insuing splash of matter they have used a mixture of arguments based on the "particle" and the "wave" picture of quantum processes and have amplified these by certain assumptions about the consequences of the "strong coupling" of the mesonic fields. With Fermi's approach, one difficulty which presumably has occurred to many and is mentioned by Lewis[7] is the clash between the assumption of equilibrium within the whole of V and the velocity of propagation of action within V, which one will be inclined to assume $\leq c$. The interaction volume has no "walls" which hinder the particles to escape; this is inherent in the physical picture and is also implied by Fermi's use of wave functions corresponding to free particle states. Hence the concentration of energy will only exist for a time of the order (R/c). (Mc^2/E), which is short compared to the time (R/c) approximately required so that action should spread across the disk. Thus, if one assumes that the coupling of the outgoing particles ceases (equilibrium frozen) when they leave the interaction volume, one can hardly avoid assuming at the same time that the processes in different zones of the interaction volume are physically independent if the transverse distance in between such zones is large compared to the depth of the interaction volume along the line of approach of the initial collision. (Whether such a split into independent zones might occur in Heisenberg's model depends on the field law assumed; more specifically on when the nonlinear character of the wave propagation ceases to be important. If this is soon after the waves start to spread out from the initial disk, the split into independent zones will occur.)

One may now consider amplifying the Fermi-Heisenberg disk-model by a postulate of "independent zones." This, in effect, we will do. However, the piling of more assumptions on to a perhaps rather uncertain basis must cease at some stage or be compensated for by

[4] Lord, Fainberg, and Schein, Nuovo cimento **7**, 774 (1950).
[5] Bradt, Kaplon, and Peters, Phys. Rev. **76**, 1735 (1949).
[6] W. Heisenberg, Z. Physik **133**, 65 (1952).
[7] H. W. Lewis, Revs. Modern Phys. **24**, 241 (1952).

establishing a more secure foundation. Since we are not in a position to provide additional support in the shape of theoretical arguments, we will try to get on firmer ground by restricting our claims as to the applicability of the model. We will consider (noncentral) nucleus-nucleus collisions and leave open the question of how far the model may apply to nucleon-nucleon interactions (see Appendix).

For a heavy nucleus, divisibility into subsystems is quite a reasonable assumption, and indeed independent of whether the nucleus is Lorentz-contracted or not. On the other hand, in a noncentral collision between two heavy nuclei the ratio of the initial angular momentum to the total energy and proper mass is even larger than for a noncentral nucleon-nucleon interaction at the same specific energy. Thus Fermi's fundamental idea should apply for complex collision partners at least as well as in the case of the nucleon-nucleon interaction. Finally, it is known from experiment[3] that energetic nucleus-nucleus collisions may give rise to a strongly anisotropic distribution of the outgoing particles. We may therefore expect that the more surveyable nucleus-nucleus process can be used to test the working of Fermi's basic idea. In carrying through the computations and making the comparison with experiment there is, of course, at the back of our minds the hope that this more surveyable problem may tell us something, by analogy, about the nucleon-nucleon case.

We make the following assumptions. The system of the two colliding nuclei consists of many independent zones. For every zone we assume isotropic evaporation, even if the total collision is noncentral. Two reasons can be given for this. First, the zone, in contrast to the whole system, is not essentially flat-shaped. Second, if the zones are small compared to the whole system, the angular momentum deposited in every zone, with respect to an axis through the center of mass of the zone, is negligibly small. (On decreasing in thought the size of the zone the angular momentum deposited in it goes more quickly to zero than the mass of matter in the zone.) Therefore we do not have to worry about conservation of angular momentum when dealing with an individual zone. The angular momentum of the whole system is automatically conserved if the energy-momentum four-vector is conserved for every zone individually.[8] If the total collision is noncentral, the c.m. frames of different zones are, in general, different.

A qualitative understanding of the way in which peaking can come out of these assumptions may be obtained by looking at the picture of a star described by Dainton and Kent.[9] In this event, a very heavy ($Z=13\pm1$) and slow ($\approx 0.55c$) primary closely passes a heavy nucleus in the emulsion. The resulting excitation leads to the disintegration of the primary into six α's as well as to the evaporation of particles from the nucleus in the emulsion. Both these processes can be assumed to proceed isotropically in the respective center-of-mass frames. Since the evaporation velocity of the α's is only a few percent of c one observes the six α's from the incident nucleus forming a narrow forward cone in the laboratory frame. The particles from the emulsion nucleus go off in haphazard directions. Because of the low velocity of the primary the geometry of the event is not obscured by relativistic effects.

IV. ANALYSIS OF THE MODEL

We carry out most of the analysis in the laboratory frame and, unless otherwise specified, all quantities used below refer to it. The nucleus initially at rest in this frame will be called the target and the bombarding particle the primary. All our considerations in this section refer to a single zone. Let μ_t and μ_p denote those parts of the mass of target and primary which belong to this zone and put $\mu_p/\mu_t=q$. M shall denote the nucleon mass, and we define $\Gamma=E'/Mc^2$, where E' is the total primary energy. Assuming $\Gamma\gg1$ one finds from a simple application of the relativistic transformation laws that the center-of-mass frame of the zone moves with respect to the laboratory frame with velocity βc such that

$$\beta=\Gamma\mu_p/(\Gamma\mu_p+\mu_t). \tag{3}$$

Under the condition $1/\Gamma\ll q\ll\Gamma$ we get for $\gamma=1/(1-\beta^2)^{\frac{1}{2}}$ the approximation

$$\gamma=(\Gamma\mu_p/2\mu_t)^{\frac{1}{2}}=(\Gamma q/2)^{\frac{1}{2}}. \tag{4}$$

As explained, we assume that the particles originating in a zone are emitted isotropically in its c.m.f. Then their normalized integral distribution law is, in this frame,

$$n(\vartheta)=(1-\cos\vartheta)/2. \tag{5}$$

We have to transform this distribution into the laboratory frame, where the angles will be denoted by ϕ. By vc we denote the velocity of the particles in the c.m.f. of the zone and we put $G=1/(1-v^2)^{\frac{1}{2}}$. We will for the moment proceed as if one could ascribe to all particles the same value of v and G.[10] The velocity four-vector to be transformed is

$$Gv\cos\vartheta,\quad Gv\sin\vartheta,\quad 0,\quad G. \tag{6}$$

The transformation parameters are β, γ and the trans-

[8] This is, in fact, only true if the individual zone is infinitesimally small. Since, physically, this is not allowable, and is not meant here, there is a residual angular momentum to be conserved in every zone. Still, this neglected part is small compared to the part we take into account.

[9] A. D. Dainton and D. W. Kent, Phil. Mag. 41, 963 (1950).

[10] The average value of G can be estimated from the energy balance of the event if the primary energy is known. In events with large primary energy and large multiplicity $1\ll\langle G\rangle_{\mathrm{Av}}\ll\gamma$. We will show later that certain properties of the angular distribution are not sensitive to the individual G's.

formation reads

$$\begin{aligned} G'v'\cos\phi &= \gamma(Gv\cos\vartheta+\beta G), \\ G'v'\sin\phi &= Gv\sin\vartheta, \\ 0 &= 0, \\ G' &= \gamma(G+\beta Gv\cos\vartheta). \end{aligned} \tag{7}$$

Dividing, we find

$$\gamma\frac{\sin\phi}{\cos\phi}=\frac{\sin\vartheta}{\cos\vartheta+\beta/v}. \tag{8}$$

We put $\sin\phi=z$, express $\cos\phi$ by it and express $\sin\vartheta$ by $\cos\vartheta$. We now have a relation between z and $\cos\vartheta$ which we can use to substitute a function of z for $\cos\vartheta$ in (5) and thus obtain the transformed distribution. Actually (8) is a quadratic equation for $\cos\vartheta$, with two roots, say,

$$\cos\vartheta=A(z)\pm\sqrt{B}(z). \tag{9}$$

The two roots arise in a natural fashion: To a given value of $z=\sin\phi$ must correspond two values of ϕ between 0 and π; and to these, two values of ϑ and of $\cos\vartheta$. In the high-energy stars one has $\beta/v>1$; then the function $B(z)$ is positive and correspondingly the roots are real only for values of z up to a certain $z=z_{\max}<1$. This means that the transformed angular distribution reaches the value 1 ("all particles") at the corresponding angle $\phi=\arcsin z_{\max}$. Two different domains of the $\cos\vartheta$ space are mapped on the domain $0<z<z_{\max}$ such that $\cos\vartheta=1$ and $\cos\vartheta=-1$ both correspond to $z=0$. Therefore the transformed angular distribution reads

$$n(\phi)=\tfrac{1}{2}[1-(A+\sqrt{B})]+\tfrac{1}{2}[1+(A-\sqrt{B})], \tag{10}$$

or

$$n(\phi)=1-\sqrt{B}. \tag{10a}$$

The actual expression for B is found from (8) by simple algebra. We obtain then

$$n(z)=\frac{z^2(\gamma^2-1)+1-(1-z^2)^{\frac{1}{2}}[1-z^2(\gamma^2-1)/(G^2-1)]^{\frac{1}{2}}}{z^2(\gamma^2-1)+1}. \tag{11}$$

All the particles appear within the angular domain

$$z=\sin\phi\leq[(G^2-1)/(\gamma^2-1)]^{\frac{1}{2}}. \tag{12}$$

The relation (11) is cumbersome but exact, that is, it holds for all values of γ and G provided $\gamma\geq G$. For $\gamma\gg1$ it is easily simplified to read

$$n(z)=\frac{z^2\gamma^2+1-[1-z^2\gamma^2/(G^2-1)]^{\frac{1}{2}}}{z^2\gamma^2+1}. \tag{13}$$

If we restrict ourselves to the dense center-part of the distribution, that is, to the angular domain $z\ll(G^2-1)^{\frac{1}{2}}/\gamma$, we can further simplify and obtain

$$n(z)=z^2\gamma^2/(z^2\gamma^2+1). \tag{14}$$

The simple approximation (14) gives $n=1$ only for $z=\infty$, instead of for $z=(G^2-1)^{\frac{1}{2}}/\gamma$. However, the numerical error in the whole region where the exact distribution does not vanish is very small. If we take, e.g., $G=10$ (which is a small-to-moderate value) and γ very large, the limit of the exact distribution defined by $(G^2-1)^{\frac{1}{2}}/\gamma=\sqrt{(99)}/\gamma$ gives, on using (14), $n=0.99$ instead of the exact value $n=1$. It is of interest to note that the approximation (14) conserves the property $\phi_n/\phi_{1-n}=n/(1-n)$, or, more specially, $\phi_{3/4}/\phi_{1/4}=3$, $X=1$ of an isotropic distribution subjected to a strong Lorentz transformation; thus we may use (14) in a treatment of the X measure of anisotropy of Kaplon and Ritson. Finally, after having obtained (14) we may now drop the assumption of a single value of G. The central part of the distribution (13) is approximated by (14) and this evidently does not depend on the G's. Also, since (14) is a small-angle approximation anyhow, we may put $z=\sin\phi\approx\phi$.

V. LIMITS FOR THE ANISOTROPY BY SUPERPOSITION

If we wanted to obtain the total angular distribution of the event in the laboratory frame we would have to sum or integrate over all the zones. Every zone would give a contribution according to (14), with a value of γ depending by (4) on the q of the zone. The contribution from every zone would have to be multiplied by a *weight* equal to that fraction of all the particles in the event which comes from the zone. For a central collision of two nuclei of equal weight all q's are 1 and (14) gives the total distribution, X being equal to 1. The wider the variation in q for the different zones is and the larger the contributions to the total multiplicity (i.e., the weights) of those zones which have extreme values of q the larger an X (the stronger anisotropy) one may expect for the distribution resulting from the summing or integrating procedure. If q varies within a certain range only, from q_1 to $q_2=t^2q_1$ say, the maximum anisotropy formally possible will arise if there are just two zones, of weight $\frac{1}{2}$ each, with values of q as defined by the limits of this range, i.e., q_1 and q_2. The corresponding X is an upper limit for the anisotropy possible on our model in an event where q varies within the range from q_1 to q_2. We will now find such upper limits for X, based on shifting all the weight to the extreme values of q. (We have met, in the Dainton-Kent star, with an event—albeit at much lower energy—where two zones actually suffice for a reasonably exact description.) To the limiting values q_1 and $q_2=t^2q_1$ will from (4) correspond two values γ_1 and γ_2 such that $\gamma_2^2=t^2\gamma_1^2$. The distribution of maximum anisotropy compatible with this range of q or γ is, using (14) and writing γ for γ_1,

$$n(\phi)=\frac{1}{2}\left[\frac{\phi^2\gamma^2}{\phi^2\gamma^2+1}+\frac{\phi^2t^2\gamma^2}{\phi^2t^2\gamma^2+1}\right]. \tag{15}$$

We may now find the value of X corresponding to a certain value of t. We may presuppose that t^2 is fairly large compared to 1 so that the second term in the bracket will with increasing ϕ tend towards 1 much faster than the first term. The values of ϕ for which $n=\frac{1}{4}$ and for which $n=\frac{3}{4}$ can then be found in close approximation by putting the second term in the bracket and the first term in the bracket equal to $\frac{1}{2}$, respectively. The results are $\phi_{1/4}=1/(\gamma t)$, $\phi_{3/4}=1/\gamma$. Therefore

$$X=\phi_{3/4}/3\phi_{1/4}=t/3. \quad (16)$$

Checking these last approximations shows that the error of (16) is smaller than 1 percent for all $t\geq 6$, $X\geq 2$.

We must now consider what effective range of q is physically reasonable. This depends on the distribution of matter density one chooses to attribute to the disks, and on the dependence of the number of particles produced in a zone on the collision energy and the proper mass initially deposited in it. One can hardly avoid assuming that the number of particles from a zone depends on these quantities in such a manner as to become very small when they do. If this is granted, and furthermore if the matter density in the disks is chosen so as to give the original collision-partners a more or less uniform density inside and a fairly sharp boundary, the effective range of q cannot be large. We estimate on the basis of some graphical work that in a collision with large impact parameter $(r\approx R)$[11] appreciable contributions to the total multiplicity may at most come from zones with q values ranging from $\frac{1}{6}$ to 6. Correspondingly, the upper limit for X is, from (16): $X<2$. (By contrast $X\approx 3$ is, on Fermi's original estimate, the expected value for median impact parameter.) Thus, if we make what seem the most natural assumptions about the weight of the zones as a function of q, we get very little peaking. Turning to the large values of X observed in nucleon-nucleus collisions[2] and considering a value $X=20$ (which is quite frequent), it is evident from (16) that to explain it with our model we would have to require q to vary from a value $q_{\min}$ $<1/60$ to a value $q_{\max}>60$. That the effective range of q should have such a huge span seems quite impossible, however strange and (from our present knowledge) unexpected a law for the weighting of the zones one might choose.

VI. CONCLUSIONS

We have pointed out a model of the extreme energy nucleus-nucleus collision which (a) operates with a flattened interaction volume and (b) satisfies the conservation laws. This model fails to give appreciable peaking, even for the largest impact parameters. Consequently we feel it is doubtful whether conservation of angular momentum is a good explanation for the peaking, even for events which may be nucleon-nucleon collisions. In terms of our model, the origin of the peaking must be in the individual zones: Our initial assumption of isotropy for the particles from one zone must be incorrect. This seems to imply that equilibrium is not even reached within a zone. If this conclusion holds for nucleon-nucleon-collisions, it contradicts the basic assumptions of the statistical theory.

Any theory of the peaking must not only explain the strong anisotropy, but also the large variation in anisotropy which is observed. This was indeed a strong point of the Fermi theory. If we can no longer ascribe the variation in anisotropy to different impact parameters, we must attribute the fluctuations in the angular distribution to the nature of the interaction itself. For this reason we believe that Heisenberg's explanation of anisotropy (wave coherence and a flat interaction zone) is not satisfactory. It does lead to strong anisotropy, but it cannot explain, so far as we can see, strongly varying anisotropy.

It seems to us that the best hope for an explanation of the angular distribution in extreme energy stars lies, after all, in assuming a cascade process and renouncing the idea that the jets are generated in a single act.

We are grateful to Professor W. Heitler for various illuminating discussions on the interpretation of high-energy events. Professor E. Fermi has very kindly sent us some comment on a draft of this work; his remarks have enabled us to clarify, we hope, the issue involved.

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APPENDIX

We briefly indicate one of the difficulties the conception of independent zones meets with if applied to nucleons.

The interaction volume has the depth $s=R/\xi$, with $\xi=E/Mc^2$. The volume of one zone must be of order of magnitude s^3. The number of zones must be of order of magnitude ξ^2. Thus, with increasing primary energy the number of zones increases more quickly ($\propto \xi^2$) than the energy itself ($\propto \xi$), and the energy deposited in one zone must *decrease*. Therefore, the higher the energy, the less it seems justified to use statistical arguments if they are to refer to a single zone. For suitably high energies ($\xi\geq 10^3$, $E'\geq 10^{15} ev$), one is led to the picture that in most of the zones nothing at all happens, while in some of them one single outgoing particle appears. Such a picture, while not necessarily wrong, is yet quite different from what one would expect for "strong coupling" interactions.

[11] In Fermi's notation.