Coherent Higgs-boson production in relativistic heavy-ion collisions

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Coherent heavy-ion collisions at ultrarelativistic energies can be used as a powerful source for the two-photon production of new particles within or beyond the standard model. Although the coherency condition limits the mass range, for beam energies of few TeV/nucleon, particles with masses up to few hundred GeV will be produced with large cross sections $\sim Z^4$ and low hadronic background. This can be achieved by accelerating, e.g., heavy ions at the CERN Large Hadron Collider and the Superconducting Super Collider (SSC) up to energies of 3.5–8 TeV/nucleon. Then the cross section for the production of a 200-GeV heavy Higgs boson in an 8-TeV/nucleon UU collision is 300 pb, and is larger than that expected from *pp* collisions at the SSC. Another example is the production of a H^+H^- pair of charged Higgs bosons, predicted by the different extensions of the standard model. In a 1–10-TeV/nucleon gold-on-gold collision the cross section for $M_{H^+}=20$ GeV is 0.3 nb–0.3 mb, compared to the cross section expected at the CERN e^+e^- collider LEP of 1 pb–0.3 nb.

I. INTRODUCTION

Spontaneous symmetry breaking in the standard model requires the existence of at least one scalar particle, the Higgs boson,[1] but there might as well exist more than one Higgs multiplet, as, for example, in supersymmetry,[2] or the Higgs boson might be a composite object, according to technicolor models.[3] Common to all phenomenologically interesting extensions of the standard model is the prediction of additional charged and neutral Higgs particles. Since there exist no definite theoretical predictions about the masses of the Higgs sector, experiments have to cover the entire mass range. Even the Linde-Weinberg mass bound of $M_{H^0}\geq 7$ GeV (Ref. 4) does not apply if the top quark is heavy, $M_t > 80$ GeV. A great deal of theoretical studies exist concerning the production and detection of the Higgs scalars in e^+e^-, *ep*, and *pp* colliders.[5–7]

We suggest to look for the two-photon production of neutral and charged Higgs particles in coherent ultrarelativistic heavy-ion collisions. By "coherent" we mean such collisions, where the nuclei do not break. The advantages are large cross sections due to the Z^4 factor, and the absence of the *WW*, *gg*, and *qq* contributions to the hadronic background.

In Sec. II we compare relativistic heavy-ion colliders to existing and planned e^+e^- and *pp* machines as far as two-photon and other parton-parton processes are concerned. We give an overall idea of what kind of physics might be interesting to look at in coherent heavy-ion collisions. In Sec. III we calculate the coherent two-photon production of neutral Higgs bosons and estimate the contribution of the neutral currents. The production of W^+W^- pairs as the main background to the production of heavy exotic particles has been calculated also. The cross section for charged Higgs bosons is given in Sec. IV.

II. TWO-PHOTON LUMINOSITIES

In the equivalent-photon approximation (EPA) (Ref. 8) cross sections for two-photon particle production (Fig. 1) can be written in a simple and universal way:

$$\frac{d\sigma}{dW^2}(Z_1Z_2\to Z_1Z_2X)=\frac{dL_{\gamma\gamma}}{dW^2}\hat{\sigma}(\gamma\gamma\to X)\ . \qquad (1)$$

The main physics is contained in the real two-photon subprocess $\hat{\sigma}(\gamma\gamma\to X)$. The two-photon luminosity function $dL_{\gamma\gamma}$ depends only on the type of the colliding particles and is very useful for comparing the relative efficiency of the different colliders as far as $\gamma\gamma$ or other parton-parton processes are concerned. It is given by integration over the equivalent-photon spectra of the two colliding ions[9]

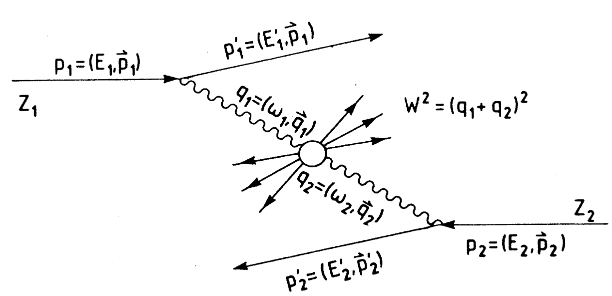

FIG. 1. Particle production in a two-photon heavy-ion collision.

$$dL_{\gamma\gamma} \simeq \int dn_1(\omega_1) \int dn_2(\omega_2) \delta(W^2 - 4\omega_1\omega_2) \tag{2}$$

with

$$dn_i = \frac{Z_i^2 \alpha}{2\pi} \ln \frac{q_{i\max}^2}{q_{i\min}^2} \frac{d\omega_i}{\omega_i} \ . \tag{3}$$

$W^2 = (q_1 + q_2)^2$ is the invariant $\gamma\gamma$ mass squared, Z_i the nuclear charges, and ω_i the photon energies in the center of mass of the colliding nuclei. In the high-energy approximation, where the EPA applies, the photon mass is related to the photon energy through

$$q^2 \geq \frac{\omega^2}{\gamma^2} \tag{4}$$

so that we set

$$q_{i\min}^2 \simeq \frac{\omega_i^2}{\gamma^2} \ . \tag{5}$$

$\gamma = E^{\text{c.m.}}/M$ is the relativistic factor. To take care of the nuclear elastic form factor, which drops rapidly for large momentum transfer, we set

$$q_{i\max}^2 \simeq \frac{1}{R_i^2} \ . \tag{6}$$

This corresponds to a nuclear charge density with a sharp edge at the radius R, since $q_{\max}^2 \simeq 1/b_{\min}^2$, where b is the impact of the nucleus. This approximation is good enough for heavy nuclei and guarantees the coherency of the process. Since the nuclear radius R varies between 2 and 8 fm, $q_{\max}^2$ is for heavy nuclei very small ~ 600 MeV^2. In other words, coherent particle production occurs in distant, i.e., peripheral nuclear collisions. After integration over the photon energy

$$\frac{R_2 W^2}{4\gamma} \leq \omega_1 \leq \frac{\gamma}{R_1} \ , \tag{7}$$

the differential two-photon luminosity in the center of mass of the colliding nuclei is simply given by

$$\frac{dL_{\gamma\gamma}}{dW^2} = \frac{16}{3} \frac{Z_1^2 Z_2^2 \alpha^2}{\pi^2} \frac{1}{W^2} \ln^3 \frac{2\gamma}{\sqrt{R_1 R_2} W} \ . \tag{8}$$

In the case of an electron, the photon spectrum involves additional terms of order $O(\omega/E)$ to the leading-log term of Eq. (3), since here the maximum photon energy can be approximately equal to the electron energy. So for an e^+e^- collision the two-photon luminosity in Eq. (8) is replaced by Low's approximate expression[10]

$$\frac{dL_{\gamma\gamma}^{e^+e^-}}{dW^2} = \frac{8\alpha^2}{\pi^2} \frac{1}{W^2} \ln^2 \frac{\sqrt{s}}{m_e} \left[\ln \frac{\sqrt{s}}{W} - \frac{3}{4} \right] , \tag{9}$$

where s is the center-of-mass energy squared and m_e the electron mass.

If we define by Ω the ratio of the nuclear to the e^+e^- two-photon luminosity, we see that heavy-ion collisions can be by a factor $\sim Z^4$ more effective photon sources compared to e^+e^- collisions, if we have the same relativistic factor in both cases. On the other hand, if we compare heavy-ion and e^+e^- collisions at the same incident energy per particle, the absence of a form factor for the electron and its very small mass tend to compensate partly the Z^4 factor in the heavy-ion $\gamma\gamma$ luminosity. This effect is shown in Fig. 2, where the ratio of the gold-on-gold to the e^+e^- two-photon luminosity Ω is plotted as a function of the two-photon invariant mass at different values of the center-of-mass energy per particle denoted by $x = E^{\text{c.m.}}/A$. For $x = 100$ GeV, e.g., we have a comparison of the CERN e^+e^- collider LEP 2 to the relativistic heavy-ion collider (RHIC) at Brookhaven. As expected, the higher the incident energy is, the larger the mass range where relativistic heavy-ion colliders are superior to e^+e^- colliders. To search for new particles beyond the reach of present colliders, i.e., $W \geq 50$ GeV, we need heavy-ion beams with an energy of few TeV/nucleon. Exceptions exist in the case where light particles such as neutral Higgs bosons, couple very feebly to photons, quarks, and gluons, and have not been seen at present colliders. The RHIC with its large two-photon luminosity in the low-mass region will present a unique possibility to look for them.

TeV energies for nuclei are already available in the upper spectrum of cosmic rays, where iron nuclei hit the top of the atmosphere with $E^{\text{lab}} \leq 10^{18}$ eV (Ref. 11). Here, of course, one cannot expect any reasonable rates since cosmic-ray fluxes are of the order of 10^{-7}–10^{-6} m^{-2}sr^{-1}sec^{-1} for iron nuclei, but one can look for single exotic events. A very interesting option, which is still under discussion, is to accelerate in the LEP tunnel ions up to energies of 3.5–9 TeV/nucleon. This would be an extension of the CERN Large Hadron Collider (LHC) project,[12] originally planned as a $\sqrt{s} \sim 18$ TeV pp collider at very high luminosities $\mathcal{L} \sim 10^{33}$ cm^{-2}sec^{-1}. Since for very heavy nuclei the charge to mass ratio goes down to $\sim 1/2.5$, only 40% of the proton energy can be achieved for each nucleon in gold or uranium beams. Assuming that protons can be accelerated up to an energy of $\sqrt{s} \sim 40$ TeV at the Superconducting Super Collider

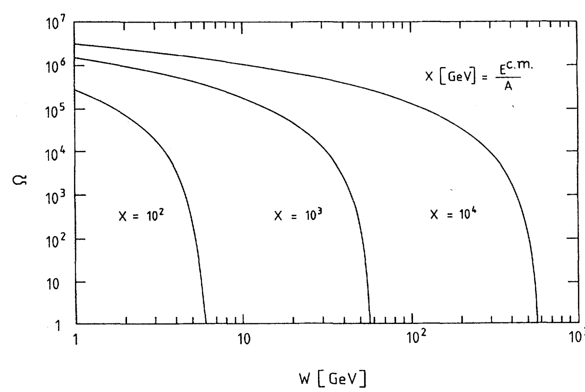

FIG. 2. The ratio Ω of the nuclear to the e^+e^- two-photon luminosity is plotted as a function of the invariant mass W. The nuclear luminosity has been calculated for $^{197}_{79}$Au-$^{197}_{79}$Au collisions. The different curves correspond to different center-of-mass energies per particle $x = E^{\text{c.m.}}/A$.

(SSC) the maximum beam energy for uranium there would then be 8 TeV/nucleon.

The RHIC at Brookhaven National Laboratory (BNL), which is planned for the mid 1990s, will provide ions with masses spanning the full periodic table and energies up to 100 GeV per nucleon.[13] The particles that can be produced coherently in a two-photon collision cannot be heavier than 5–18 GeV. This is due to the fact that in order that the nuclei do not break up, the maximum momentum transfer to the nucleus should be $\sim 1/R$. This factor enters the photon-luminosity in Eq. (8) and leads to the constraint $W \leq 2\gamma/R$. The dependence of the two-photon luminosity on the particles' invariant mass is shown for different ion pairs in Fig. 3. Here we multiplied the dimensionless quantity $W^2\, dL_{\gamma\gamma}/dW^2$ by the machine luminosities $\mathcal{L}$ ($\mathrm{cm}^{-2}\mathrm{sec}^{-1}$), which are planned for the acceleration of different nuclei species at the RHIC (Table I). The resulting effective luminosities dL/dW^2 are quite low for heavy ions. This affects the production rates and partly cancels the advantage of the high charge of the heavy ions. The problem with heavy ions lies on the considerably larger energy losses present in the beam. Coulomb dissociation, bremsstrahlung pair production, beam-beam and beam-gas nuclear reactions, and intrabeam scattering become for heavy beams ($A \geq 100$) severe restrictions on the luminosity of the machine. This is a disadvantage of relativistic heavy-ion colliders compared to e^+e^- and pp machines. It is, however, possible to increase the luminosity values of Table I by 1–2 orders of magnitude.

We may give now a short idea of the two-photon processes that will be interesting to examine at the RHIC. The production of light Higgs bosons will be discussed in the next section. Charm production ($m_{c\bar{c}} \sim 3.5$ GeV) is favored by the copper system, $b\bar{b}$ states ($m_{b\bar{b}} \sim 10.3$ GeV) will be produced mainly in carbon collisions. To compare we show in Fig. 4 the two-photon effective luminosities for e^+e^- colliders, which are in operation or still under construction, together with the gold-on-gold $\gamma\gamma$ effective luminosity at the RHIC. Using the values of Table I the charmonium production at the RHIC will be comparable to LEP 1. Bottom production is out of the question, since the rate lies even below that of the DESY e^+e^- storage ring PETRA. We emphasize that this

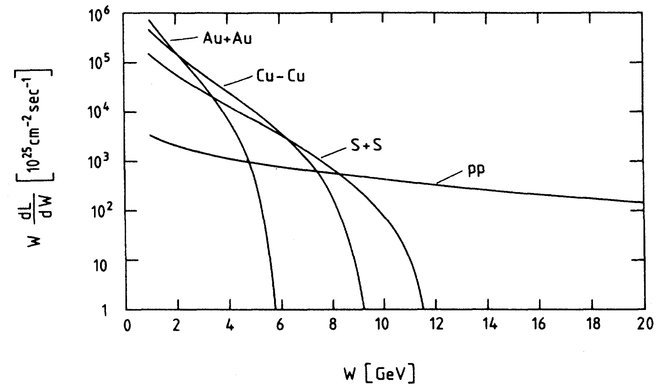

FIG. 3. The dependence of the effective two-photon luminosity on the invariant mass W is shown for different ion pairs at RHIC.

changes if the machine luminosity is increased. On the other hand, a great deal of low-mass phenomena can be studied at the RHIC at very high luminosities. We give two examples. Pushing down the limit for the two-photon coupling of the $\eta(1440)$, $f_2(1720)$, and $\zeta(2220)$ will help establish them eventually as glueball states. The widths of the scalar mesons $f_0(975)$ and $a_0(980)$ seem to be small compared with those of the tensor mesons, which indicates that they might be compound states of two $q\bar{q}$ pseudoscalar pairs.[14] Controversies in the theory can be clarified by more accurate experiments.

We show now that heavy-ion colliders are interesting when compared with planned pp colliders such as the LHC and the SSC. In a hadron collider the beam energy is shared among the consituents of the proton, so that the center-of-mass energy in the parton-parton collisions $\hat{s} = W^2$ is smaller than the one in the pp system (approximately $\frac{1}{6}$). Proton cross sections can be written in terms of the parton cross sections $\hat{\sigma}(ij \to x)$ and the parton-parton luminosities dL_{ij}/dz (where $z = W^2/s$) as in the two-photon case. Since $\hat{\sigma} \sim C/W^2$ with $C = \alpha_s^2/\pi^2$ for a strong process and $C = \alpha^2/\pi^2$ for an electroweak one, the quantity $\sigma_0 = (C/W^2)\, dL_{ij}/dz$ with the dimension of a

TABLE I. Beam energies and effective luminosities at present and future high-energy colliders.

Collider	Colliding particles	E_B (GeV)	$\mathcal{L}$ ($\mathrm{cm}^{-2}\mathrm{sec}^{-1}$)
PETRA (DESY)	e^+e^-	17	4×10^{30}
SPS (CERN)	$p\bar{p}$	270	2×10^{29}
SLC (SLAC)	e^+e^-	50	2×10^{30}
LEP 1 (CERN)	e^+e^-	50	2×10^{31}
LEP 2 (CERN)	e^+e^-	100	5×10^{31}
HERA (DESY)	ep	260 e, 820 p	2×10^{31}
RHIC (BNL)	pp	$250A$	1.2×10^{31}
	${}^{12}_{6}\mathrm{C}\,{}^{12}_{6}\mathrm{C}$	$125A$	6×10^{29}
	${}^{22}_{16}\mathrm{S}\,{}^{32}_{16}\mathrm{S}$	$125A$	6×10^{28}
	${}^{64}_{28}\mathrm{Cu}\,{}^{64}_{28}\mathrm{Cu}$	$125A$	2×10^{28}
	${}^{197}_{79}\mathrm{Au}\,{}^{197}_{79}\mathrm{Cu}$	$100A$	1.2×10^{27}

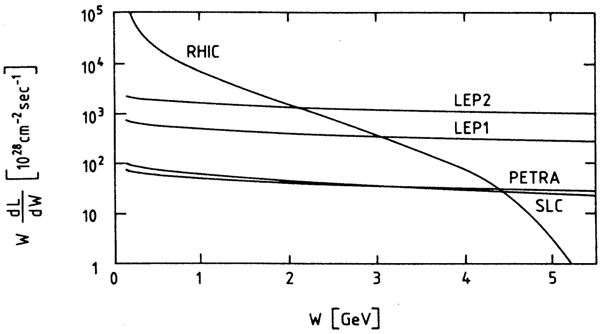

FIG. 4. Effective two-photon luminosities as a function of the invariant mass W for various e^+e^- colliders and for gold-on-gold collisions at RHIC.

cross section provides a nice measure of the reach of a hadron collision at a given energy and luminosity. The comparison of the $\gamma\gamma$ and gg luminosities in Ref. 15 showed that two-photon luminosities at LEP I can be up to 4 orders of magnitude smaller than the two-gluon luminosities at the CERN SPS. This can be easily compensated by the Z^4 factor at the RHIC. Scaling this behavior into the TeV energy region we expect that two-photon collisions with heavy ions will give comparable or even larger cross sections than pp collisions for the production of particles with masses up to few hundred GeV. The main advantages in the first case will be rather clear signals since the production is quasielastic and there exist no uncertainties in the two-photon luminosity distribution as in the case of gluons or quarks.

III. PRODUCTION OF NEUTRAL HIGGS BOSONS

Since the coupling of the Higgs boson to light particles, such as the partons of a proton beam or the virtual photons in an e^+e^- collision, is very small, the present experimental mass bound is quite low $M_{H^0} \geq 5.4$ GeV, leaving a mass window of 18–210 MeV still open.[16]

LEP 200 will be capable of finding the Higgs boson in the process $e^+e^- \rightarrow ZH$, if $M_{H^0} \leq 100$ GeV. To explore the few-hundred-GeV mass region one has to go to hadron colliders such as the LHC and the SSC. In pp collisions the dominant production mechanisms are through gg, WW, and ZZ fusion, leading to cross sections of a few tenths of picobarns.[17]

A. Coherent two-photon production

The Higgs boson couples to two photons through a triangle loop (Fig. 5). To this contribute quarks, the vector bosons, and any other charged particles that may exist, e.g., Higgs scalars, squarks, etc. We concentrate first on Higgs-boson production in the standard model with no additional flavor generation. The two-photon width of the Higgs boson is given by[18]

$$\Gamma(H^0 \rightarrow \gamma\gamma) = \frac{G_F M_{H^0}^3}{8\sqrt{2}\pi} \left[\frac{\alpha}{\pi}\right]^2 |I|^2 , \tag{10}$$

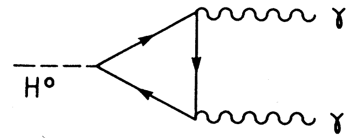

FIG. 5. Two photons coupled to a neutral Higgs boson.

where $G_F = 1.2 \times 10^{-5}$ GeV^{-2} is the Fermi constant and I the sum over the different contributions:

$$I = \sum_q q^2 I_q + \sum_l I_l + I_W$$

with the quark charge q and

$$\begin{aligned} &I_q = 3[2\lambda_q + \lambda_q(4\lambda_q - 1) f(\lambda_q)] , \\ &I_l = 2\lambda_l + \lambda_l(4\lambda_l - 1) f(\lambda_l) , \\ &I_W = 3\lambda_W(1 - 2\lambda_W) f(\lambda_W) - 3\lambda_W - \tfrac{1}{2} , \end{aligned} \tag{11}$$

where

$$f(\lambda) = -2 \left[\arcsin \frac{1}{2\sqrt{\lambda}}\right]^2 \quad \text{if } \lambda > \tfrac{1}{4} ,$$

$$f(\lambda) = \frac{1}{2} \left[\ln \frac{\eta^+}{\eta^-}\right]^2 - \frac{\pi^2}{2} + i\pi \ln \frac{\eta^+}{\eta^-} \quad \text{if } \lambda < \tfrac{1}{4} ,$$

and $\eta^{\pm} = \frac{1}{2} \pm (\frac{1}{4} - \lambda)^{1/2}$. If all $\lambda_i = M_i^2 / M_{H^0}^2 \ll 1$ (M_i are the masses of the quarks and vector bosons), only the vector-boson loop contributes and $I \sim -\frac{1}{2}$. If all $\lambda_i \gg 1$, i.e., for extremely light Higgs bosons: $I \sim \sum_q q^2 + \frac{1}{3} - \frac{7}{4} \sim \frac{1}{4}$.

The two-photon width is in general very small:

$$\Gamma(H \rightarrow \gamma\gamma) \simeq 1.8 \times 10^{-3} M_{H_0}^3 |I|^2 , \tag{12}$$

where Γ is in eV and M_{H^0} is in GeV. It lies between 10^{-3} eV for $M_{H^0} \simeq 1$ GeV and 1 MeV for $M_{H^0} \sim 600$ GeV. The dependence of the width on the t-quark mass is important, since according to present experimental limits this mass can be comparable or even larger than the vector-boson masses, so it will couple strongly to the Higgs boson. Neglecting the contribution of the light quarks and the leptons, we see that for $M_t \sim M_H \sim M_W$ the t quark and the W contribute with opposite sign and $|I|^2 \simeq 1.8$. For light Higgs bosons with a mass between 1 and 40 GeV, $|I|^2$ is almost independent of the top mass. The values for $M_H = 1$, 10, and 20 GeV are, respectively, equal to $|I|^2 = 0.2$, 1.6, and 1.9. In Table II we give further values of $|I|^2$ for different top-quark and Higgs-boson masses. We see a rather strong M_t dependence for Higgs-boson masses of a few hundred GeV. Notice that in all cases $|I|^2$ grows with M_H up to a maximum value, which is attained for $M_H \sim 160$ GeV and then falls to a constant.

What makes the two-photon width of the Higgs boson an interesting quantity by itself is that it is quite sensitive to additional heavy masses in the loop. For each heavy

TABLE II. Values of $|I|^2$ for different top-quark and Higgs-boson masses.

M_{H^0} (GeV)	$|I|^2$ $M_t=40$ GeV	$M_t=100$ Gev	$M_t=200$ GeV
40	1.9	1.9	1.9
80	1.5	2.0	2.0
100	1.6	2.1	2.2
200	6.2	4.9	5.6
300	2.7	1.3	2.5
500	0.8	0.3	0.3

quark $I_q \simeq 1$ and for each heavy lepton $I_l \simeq \frac{1}{3}$. If charged scalars are present, for example, charged Higgs bosons, they contribute a factor $I_s = -\lambda_s[1+2\lambda_s f(\lambda_s)]$ to I. Because of the partial cancellations between the fermion and boson loops the dependence of the two-photon width on the various heavy masses in the theory is in general quite complicated.

There are also theories such as supersymmetry (SUSY) that predict additional neutral Higgs bosons. Now the Yukawa couplings of these Higgs bosons to the fermions contain additional parameters, which for certain values can lead to large enhancements.[15] Then

$$\frac{\Gamma(S^0 \to \gamma\gamma)}{\Gamma(H^0 \to \gamma\gamma)} \simeq \frac{M_{S_0}^3}{M_{H^0}^3} Y_f^2 , \tag{13}$$

where Y_f are the enhancement factors of the Yukawa couplings of the extra neutral Higgs boson S^0 to quarks and leptons. In minimal supersymmetry $Y_f^2 \sim 10^2 - 10^3$ (Ref. 19).

The two-photon production of H^0 in a nucleus-nucleus collision is given by integrating Eq. (1) over W^2 with the resonance cross section taken in the narrow-width approximation:

$$\sigma_{\gamma\gamma \to H^0} \simeq 8\pi^2 \frac{\Gamma_{\gamma\gamma}}{M_{H^0}} \delta(W^2 - M_{H^0}^2) . \tag{14}$$

Then the total cross section is given by

$$\sigma(Z_1Z_2 \to Z_1Z_2H^0) \simeq \frac{128}{3}(Z_1Z_2\alpha)^2 \frac{\Gamma_{\gamma\gamma}}{M_{H^0}^3} \times \ln^3 \frac{2\gamma}{\sqrt{R_1R_2}M_{H^0}} . \tag{15}$$

Notice that the dependence of the total cross section on the Higgs-boson mass is only logarithmic, since $\Gamma_{\gamma\gamma} \sim M_{H^0}^3$.

For low masses $M_{H^0} \ll 2\gamma/R$ the cross section is approximately equal to

$$\sigma \sim g(I) Z^4 \ln^3 \frac{2\gamma}{R} \tag{16}$$

with $g(I) = 3.7$ ab (1 ab $= 10^{-42}$ cm^2). In a gold-on-gold collision with a relativistic factor of $\gamma \simeq 10^2 - 10^4$ this gives a cross section of $\sigma \sim 0.7-37$ nb. Without the Z^4 factor this cross section would have been beyond any practical reach.

For large masses $M_{H^0} \simeq 2\gamma/R$ the cross section drops quite fast, due to the argument of the logarithm. The dependence of the cross section on the incident energy and on the mass of the Higgs boson is shown in Fig. 6. To compare with heavy-Higgs-boson production in pp collisions[17] we have plotted in Fig. 7 the cross section for UU and pp collisions at the maximum energy attainable at the LHC and the SSC. Assuming a proton energy of 9 GeV at the LHC, for uranium the maximum energy will be 3.5 TeV/nucleon and the UU cross section is larger than the pp cross section for $M_H \le 100$ GeV. At the SSC the beam energies will be 20 TeV for protons and 8 TeV/nucleon for heavy ions. At this energy coherent uranium collisions will be more effective Higgs-boson sources than proton collisions if the mass of the Higgs boson is less than 300 GeV. Notice the dependence of the cross section on the mass of the top quark.

The dominant decays of a heavy Higgs boson are into W^+W^- and Z^0Z^0 pairs. The $\Gamma(H^0 \to W^+W^-)$ decay width is $\sim 4\pi^2/\alpha^2 \simeq 10^6$ times larger than the two-photon width and the $\Gamma(H^0 \to Z^0Z^0)$ by a factor half as large. Since W^+W^- pairs will be produced also directly in $\gamma\gamma$ fusion with even larger cross sections (for a discussion see the end of this section), there is no way to identify the Higgs boson in this channel, but it can be identified in the Z^0Z^0 channel by looking, for example, at the four-lepton spectrum $(l^+l^-)(l^+l^-)$. The branching ratio for this decay channel is $B(H^0 \to Z^0Z^0 \to (l^+l^-)(l^+l^-)) \sim 10^{-3}$. However, one should be careful with the spectrum of very light lepton pairs such as electrons, since at very high energies $\gamma \ge 5 \times 10^2$ multiple pair production dominates. The process $\gamma\gamma \to (e^+e^-)(e^+e^-)$ is asymptotically constant and equal to $\sigma_0 \simeq 7 \times 10^{-30}$ cm^2 (Ref. 20), so the total cross section for double pair production in a heavy-ion collision grows like

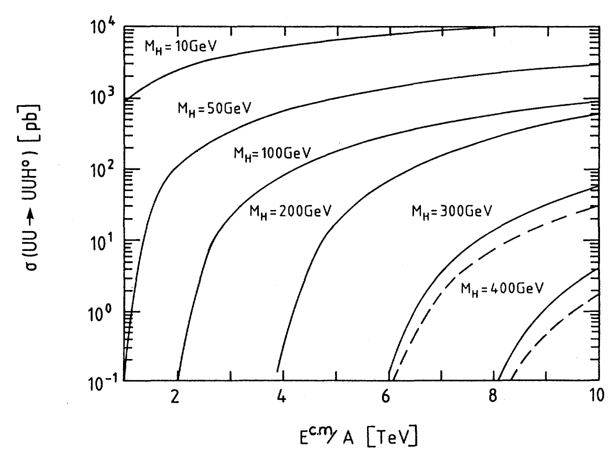

FIG. 6. Total cross section for the production of Higgs bosons in $^{238}_{92}$U-$^{238}_{92}$U collisions as a function of the incident energy per nucleon. Solid lines correspond to $M_t = 50$ GeV and dashed lines to $M_t = 100$ GeV.

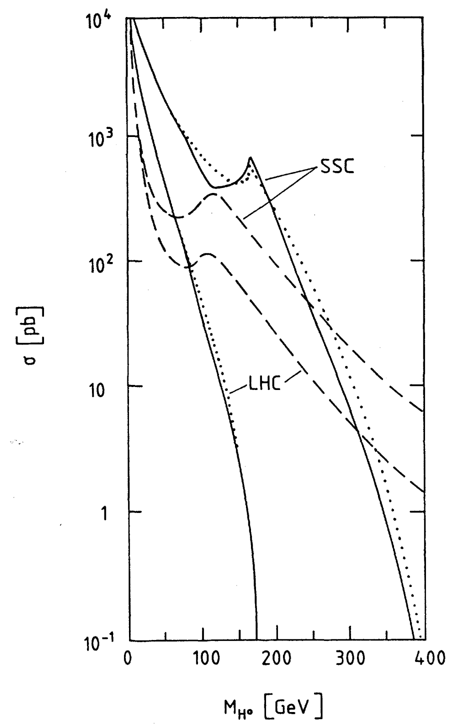

FIG. 7. Total cross section for the production of Higgs bosons in ${}^{238}_{92}U$-${}^{238}_{92}U$ collisions (solid and dotted lines) and in pp collisions (dashed lines). We have assumed that at the LHC the beam energy will be 9 TeV for protons and 3.5 TeV for uranium. At the SSC the corresponding beam energies are 20 and 8 TeV. Solid and dashed lines have been computed with $M_T = 50$ GeV, dotted lines with $M_t = 100$ GeV.

$$\sigma \simeq \frac{8}{3}\frac{Z^4\alpha^2}{\pi^2}\sigma_0 \ln^4(2\gamma) \ . \tag{17}$$

Notice that in the case of e^+e^- pair production one has to use $q^2_{\max} \simeq 2m_e^2$ instead of the inverse nuclear cutoff, which we used previously [Eq. (6)]. This is due to the approximations entering the EPA. In a 1-TeV/nucleon gold-on-gold collision, for example, this cross section is extremely high ~ 1.7 b. It is possible that at these energies e^+e^- pairs behave like massless bosons and show a quasi-infrared divergency, similar to the soft-photon radiation problem in QED, so one has to sum a large number of e^+e^- pairs to get the right answer.

To identify the Higgs boson in the other decay channels of the Z^0Z^0 to the $(q\bar{q})(q\bar{q})$ or the $(q\bar{q})(l^+l^-/\nu\bar{\nu})$ final states, one has to know the background coming from the four-jet and two-jet production in two-photon collisions.

A Higgs boson with a mass $2M_t \le M_{H^0} \le 2M_W$ will decay into a $t\bar{t}$ pair. On the other hand, $t\bar{t}$ pairs will be also produced directly in a two-photon collision. The ratio of the Higgs-boson contribution to the $t\bar{t}$ background is roughly determined by

$$R_{t\bar{t}} \sim \frac{8\pi M_t^2}{\alpha^2}\frac{\Gamma_{\gamma\gamma}}{M_H^3} \tag{18}$$

which for $M_t = 50$ GeV is $\sim 2\times 10^{-3}|I|^2$. Also, angular cuts do not really affect the value of $R_{t\bar{t}}$, so in order to see a Higgs boson in this mass range one has to look for other decay modes. One possibility is the decay $H^0 \to Z^0\sum q\bar{q}$, the sum being over all quarks except the top quark. The branching ratio ranges from 10^{-3} for a Higgs boson with a mass of 160 GeV to 10^{-4} for a Higgs boson with a mass of 120 GeV. For masses around the $t\bar{t}$ threshold the branching ratio is too small (10^{-6}) for practical use.

To conclude we discuss now the effect of the detector acceptance on the two-photon particle production.

In order to identify the Higgs boson at least some of its decay products should be seen in the detector. To estimate the reduction in luminosity resulting from the experimental cuts we rewrite the two-photon luminosity in terms of the two-photon rapidity:

$$y = \frac{1}{2}\ln\frac{\omega_1}{\omega_2} \tag{19}$$

as

$$\frac{d^2L_{\gamma\gamma}}{dy\, dW^2} = \frac{8(Z_1Z_2\alpha)^2}{\pi^2}\left[\ln^2\left[\frac{2\gamma}{W\sqrt{R_1R_2}}\right] - y^2\right]\frac{1}{W^2} \ . \tag{20}$$

In the absence of experimental cuts

$$y \le \ln\frac{2\gamma}{W\sqrt{R_1R_2}} \tag{21}$$

but in general it will be smaller. For an acceptance angle of, e.g., $\theta_c \simeq 10$ rad the maximum rapidity is only $y_c \simeq 1.6$ and the two-photon luminosity function of Eq. (8) is reduced by a factor

$$L_A = \frac{3y_c\left[\ln^2\frac{2\gamma}{W\sqrt{R_1R_2}} - y_c^2\right]}{\ln^3\frac{2\gamma}{W\sqrt{R_1R_2}}} \ . \tag{22}$$

The total cross section for an $M_H \sim 200$-GeV Higgs boson in a 8-TeV/nucleon UU collision is then smaller by a factor $\frac{1}{2}$ compared to the value given in Fig. 7. However, for $y_c \ge 1.8$ there would be no reduction.

We turn now to other coherent subprocesses in heavy-ion collisions, which may compete with two-photon fusion. Since normal nuclei are color neutral, gluon processes do not contribute to coherent particle production. Clearly the same applies for weak charged currents. There exist, however, coherent neutral-current contributions, which will be briefly discussed in the next section.

A scalar particle such as the Higgs boson also couple directly to the nucleus with a coupling constant, which becomes of order 1 for heavy nuclei, since it is proportional to G_F and to the nuclear mass squared. So a Higgs boson might be emitted with a large cross section directly by one of the nuclei in nuclear elastic scattering, if its mass is $\simeq 1/R \simeq 30-200$ MeV. The differential cross section scales with $1/(M_f^2 - M_i^2)^2$, as can be checked on purely kinematical grounds,[21] where M_i is the mass of the incident nucleus and M_f the mass of the final Higgs-boson–nucleus state, so it peaks at small Higgs-boson masses. Another factor entering this cross section is the diffractive nucleus-nucleus scattering, which can be quite large. A detailed calculation will be given elsewhere.

B. Coherent neutral currents in heavy-ion collisions

Isoscalar nuclei can couple to the weak neutral current as a single particle as well. On the scale of few-TeV/nucleon incident energies the Z^0 gauge boson can be viewed as a light particle emitted by the nucleus, just as the photon, and we can estimate the $Z^0 Z^0 \to H^0$ contribution, using the equivalent vector-boson approximation.[22] Since in this approximation the Z^0 is treated as quasireal, $\omega \geq M_{Z^0}$, and so only Higgs bosons with $M_{Z^0} \geq 2M_{Z^0}$ can be considered. If in the expression for the photon spectrum we replace the electromagnetic coupling $Z^2\alpha$ by

$$g_V^2 + g_A^2 = \frac{1}{4} \frac{g^2}{\cos^2\theta_W} [(T_3 - 2Z \sin^2\theta_W)^2 + T_3^2] \; , \quad (23)$$

we get after integration over the Z propagator an approximate expression for the transverse Z spectrum given by[23]

$$dn_{TZ} = \frac{g_V^2 + g_A^2}{4\pi^2} \ln \left[1 + \frac{1}{R^2 M_Z^2} \right] \frac{d\omega_i}{\omega_i} \; , \quad (24)$$

where we set again the maximum momentum transfer equal to $1/R$. T_3 is the third component of the isospin, θ_W the Weinberg angle, and $g^2 = \alpha/\sin^2\theta_W$.

For isoscalar nuclei $(Z = N = A/2)$ T_3 vanishes and the coupling constant in Eq. (24) reduces to $\alpha Z^2 tg^2\theta_W \sim 0.3Z^2\alpha$, which is comparable to the electromagnetic coupling in Eq. (3). However, because of the mass suppression in the Z^0 propagator and the restriction of the nuclear momentum transfer to small values the logarithm in Eq. (24) is extremely small, $\sim 10^{-5}-10^{-7}$. The ZZ transverse luminosity is roughly given by

$$\frac{dL_{TZZ}}{dW^2} \simeq \left[\frac{g_V^2 + g_A^2}{4\pi^2} \right]^2 \frac{1}{W^2} \ln^2 \left[1 + \frac{1}{R^2 M_Z^2} \right] \ln \frac{4\gamma^2}{RW^2} \quad (25)$$

and is by a factor of $\sim(10^{-11}-10^{-15})\ln^{-2}(2\gamma/RW)$ smaller than the two-photon one. So although the $ZZ \to H$ width of the Higgs boson is by 6 orders of magnitude larger than the $\gamma\gamma \to H$ width the corresponding cross section is by several orders of magnitude smaller.

Unlike photons, on-shell vector bosons possess also a longitudinal spectrum[23]

$$dn_{LZ} = \frac{g_V^2 + g_A^2}{4\pi^2} \left[1 + \frac{1}{R^2 M_Z^2} \right] \frac{d\omega_i}{\omega_i} \quad (26)$$

which in the case of nuclei is of the same order of magnitude as the transverse spectrum.

We conclude that neither the transverse part nor the longitudinal part of the ZZ luminosity can be competitive with the two-photon luminosity present in coherent heavy-ion collisions.

IV. VECTOR-BOSON PAIR PRODUCTION

We saw previously that new particles expected to be seen in a future heavy-ion collider will decay predominantly into vector-boson pairs if their masses are above $2M_W$. On the other hand, W^+W^- pairs will be themselves produced directly in the two-photon channel. So their cross section has to be known if one is concerned with the identification of particles with masses above few hundred GeV. The $\gamma\gamma \to W^+W^-$ cross section is given by[24]

$$\sigma(\gamma\gamma \to W^+W^-) = \frac{8\pi\alpha^2}{W^2} \left[\frac{1}{t}(1 + \tfrac{3}{4}t + 3t^2)\Lambda - 3t(1-2t)\ln\frac{1+\Lambda}{1-\Lambda} \right] \quad (27)$$

with

$$t = \frac{M_W^2}{W^2} \text{ and } \Lambda = \sqrt{1-4t} \; .$$

For $t \ll 1$ the $\gamma\gamma \to W^+W^-$ cross section approaches a constant value

$$\sigma_W \sim \frac{8\pi\alpha^2}{M_W^2} \quad (28)$$

and so the total W^+W^- production in a heavy-ion collision rises with the fourth power of the logarithm of the

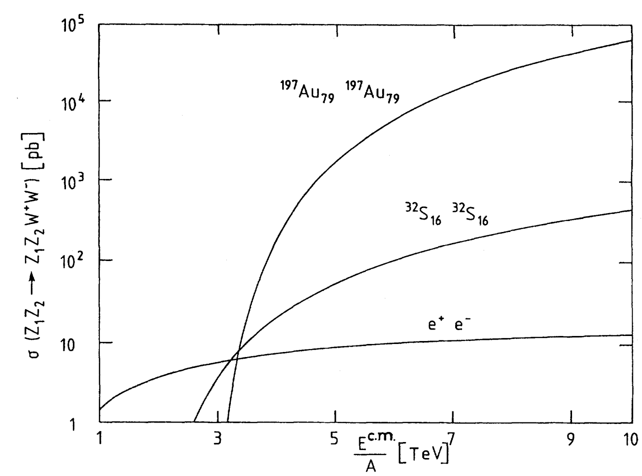

FIG. 8. Two-photon production of a W^+W^- pair in heavy-ion and e^+e^- collisions as a function of the incident energy per particle.

incident energy, while the Born cross section rises with the third logarithmic power:

$$\sigma(Z_1Z_2 \to Z_1Z_2W^+W^-) \sim \frac{8}{3}\frac{(Z_1Z_2\alpha)^2}{\pi^2}\sigma_W \times \ln^4 \frac{\gamma}{\sqrt{R_1R_2}M_W} . \quad (29)$$

The full expression is shown as a function of the incident energy per particle in Fig. 8 and is compared to the e^+e^- production. As expected we get large values of the cross section with increasing energy (pb–mb), so there is no chance to see heavy Higgs bosons in the *WW* channel, but maybe it is possible to see exotic charged particles.

V. PRODUCTION OF CHARGED HIGGS BOSONS

We turn now to the two-photon production of charged Higgs bosons in relativistic heavy-ion collisions. In certain models the masses are constrained, but since the calculation in this case is completely model independent we will treat the mass as a free parameter. The on-shell $\gamma\gamma$ cross section is given by the Born expression[6]

$$\sigma(\gamma\gamma \to H^+H^-) = \frac{2\pi\alpha^2}{W^2}\left\{(1+y)\sqrt{1-y} - 2y\left[1-\frac{y}{2}\right]\ln\left[\frac{1}{\sqrt{y}}+\left[\frac{1}{y}-1\right]^{1/2}\right]\right\} \quad (30)$$

with $y=4M_H^2/W^2$. In the large-W^2 approximation the total cross section for a heavy-ion collision is then given by integrating Eq. (1) for $4M_H^2 \le W^2 \le 4\gamma^2/R_1R_2$:

$$\sigma(Z_1Z_2 \to Z_1Z_2H^+H^-) \simeq \frac{8}{3}\frac{Z_1^2Z_2^2\alpha^4}{\pi M_H^2} \times (\ln^3 r - \tfrac{3}{2}\ln^2 r + \tfrac{3}{2}\ln r - \tfrac{3}{4}) , \quad (31)$$

where

$$r = \frac{\gamma}{\sqrt{R_1R_2}M_H} .$$

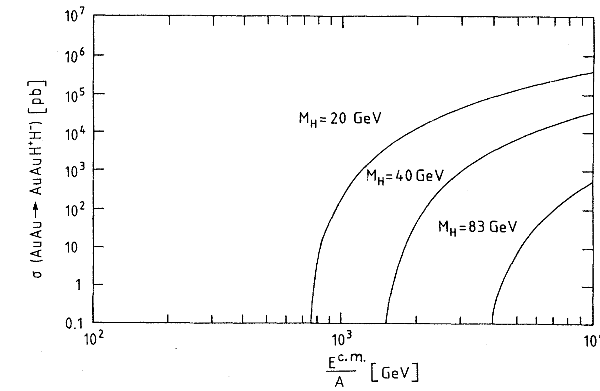

FIG. 9. Total cross section for the production of charged Higgs bosons in a nucleus-nucleus collision as a function of the incident energy per nucleon.

Notice here the leading-log behavior, which is common to all two-photon processes in the Born approximation. The results for the full expression are shown for different masses of the charged Higgs bosons in Fig. 9. To compare with LEP, we quote from Ref. 5 the maximum values of the $e^+e^- \to H^+H^-$ cross section expected to be measured there. For $M_H=20$, 40, and 80 GeV they are equal to 2×10^2, 60, and 0.2 pb.

VI. CONCLUSIONS

We have shown that the production of heavy Higgs particles with a mass $M_H \le 300$ GeV in the coherent electromagnetic field of two colliding nuclei is competitive to the Higgs-boson production at planned e^+e^- and *pp* colliders, if the incident energy is few TeV/nucleon. The measurement of the two-photon width of the neutral Higgs boson might help to establish the existence of new heavy charged particles.

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