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Higgs-particle production through vacuum excitations

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We consider here a nonperturbative mechanism for the production of Higgs particles through vacuum excitations. It appears that it could already have been seen in "chiron" events, "halo" events, and some other signals in high-energy cosmic-ray collisions.

I. INTRODUCTION

The search for Higgs particles has been one of the prime interests of present and future accelerators, since the standard model of weak and electromagnetic interactions is incomplete without the Higgs bosons being discovered. Conceptually this is particularly important since the Salam-Weinberg symmetry breaks spontaneously. We shall give here a quantum-field-theoretic description of Salam-Weinberg-symmetry breaking as vacuum realignment, extend it to include temperature dependence and identify the possible production of Higgs particles through local vacuum destabilization, including some experimental signatures in cosmic rays.[1]

We organize the paper as follows. To fix our ideas, in Sec. II we give a quantum description of phase transition with U(1) symmetry through construction of a nonperturbative vacuum as a coherent state and next generalize it to Salam-Weinberg phase transition at zero temperature. In Sec. III we shall bring in the idea of temperature and construct the "thermal vacuum" using the techniques of thermofield dynamics of Umezawa, Matsumoto, and Tachiki.[2] The methodology used here is variational and nonperturbative with a self-consistent determination of a temperature-dependent Higgs-boson mass and the effective potential. In Sec. IV we look into the possibility of *local* heating of the vacuum. The discussions here will be qualitative. In Sec. V we discuss the experimental signatures for such a process. We shall also consider cosmic-ray events, such as "chirons," "halos," and muon anomalies from the directions of Cygnus X-3 and Hercules X-1 in terms of local vacuum destabilization. In Sec. VI we shall summarize our results.

II. SPONTANEOUS SYMMETRY BREAKING: A QUANTUM DESCRIPTION

Let us first consider here the quantum description of the conventional spontaneous symmetry breaking (SSB) for U(1) symmetry: The Lagrangian here is

$$\mathcal{L}=(D_\mu\phi)^*(D^\mu\phi)-V(\phi)-\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} \tag{1}$$

with

$$V(\phi)=-m^2\phi^*\phi+\lambda(\phi^*\phi)^2 \; . \tag{2}$$

$D_\mu=\partial_\mu-igA_\mu$ is the covariant derivative. Classically, when $m^2>0$ symmetry breaking occurs with $\langle\phi^*\phi\rangle=m^2/2\lambda$. To give a quantum description to SSB of the Higgs mechanism, let us write $\phi=(1/\sqrt{2})(\phi_1+i\phi_2)$ with ϕ_1 and ϕ_2 as real fields. Thus the equal-time algebra

$$[\phi(\mathbf{x},t),\dot\phi(\mathbf{y},t)^*]=i\delta(\mathbf{x}-\mathbf{y}) \tag{3}$$

is equivalent to

$$[\phi_i(\mathbf{x},t),\dot\phi_j(\mathbf{y},t)]=i\delta_{ij}\delta(\mathbf{x}-\mathbf{y}) \; . \tag{4}$$

This equal-time algebra is for interacting fields, as it holds good for interacting operators[1-3] and is consistent with the expansions

$$\phi_1(\mathbf{x},0)=\frac{1}{\sqrt{2\lambda_x}}[a(\mathbf{x})+a(\mathbf{x})^\dagger] \tag{5a}$$

and

$$\dot\phi_1(\mathbf{x},0)=i\left[\frac{\lambda_x}{2}\right]^{1/2}[-a(\mathbf{x})+a(\mathbf{x})^\dagger] \tag{5b}$$

with a and $a^\dagger$ satisfying the commutation relation

$$[a(\mathbf{x}),a(\mathbf{y})^\dagger]=\delta(\mathbf{x}-\mathbf{y}) \; , \tag{6}$$

and, for free fields, $\lambda_x=(-\nabla_x^2+m^2)^{1/2}$, but here could be arbitrary. The Lagrangian (1) has been normal ordered with respect to $|\text{vac}\rangle$ defined by $a(\mathbf{x})|\text{vac}\rangle=0$. Let us now define a coherent state $|\text{vac}'\rangle$ as[4,5]

$$|\text{vac}'\rangle=\exp\left[\xi\int\left[\frac{\lambda z}{2}\right]^{1/2}[a(\mathbf{z})^\dagger-a(\mathbf{z})]d\mathbf{z}\right]|\text{vac}\rangle \equiv U|\text{vac}\rangle \; , \tag{7}$$

where U is a unitary operator. We wish to define an annihilation operator corresponding to $|\text{vac}'\rangle$, such that

$$\phi_1'^{\text{an}}(\mathbf{z})|\text{vac}'\rangle=0 \; . \tag{8}$$

Clearly we then have

$$\phi_1'^{\mathrm{an}}(\mathbf{z}) \equiv \frac{1}{\sqrt{2\lambda_z}} a'(\mathbf{z}) = U\phi_1^{\mathrm{an}}(\mathbf{z})U^{-1}$$
$$= \phi_1^{\mathrm{an}}(\mathbf{z}) - \frac{\xi}{2} . \tag{9}$$

Hence for the complex field ϕ, the vacuum expectation value (VEV) is given as

$$\langle \mathrm{vac}'|\phi(\mathbf{z})|\mathrm{vac}'\rangle = \frac{\xi}{\sqrt{2}} . \tag{10}$$

The expectation value of the Hamiltonian density, using Eqs. (1), (2), and (7), is given as

$$\epsilon_0 = \langle \mathrm{vac}'|\mathcal{T}^{00}|\mathrm{vac}'\rangle$$
$$= -\tfrac{1}{2}m^2\xi^2 + \frac{\lambda}{4}\xi^4 . \tag{11}$$

A minimization of the energy density ϵ_0 with respect to ξ now gives the usual result

$$\xi \equiv \xi_0 = \left[\frac{m^2}{\lambda}\right]^{1/2} , \tag{12}$$

which makes ϵ_0 negative for $m^2 > 0$, thus demonstrating that $|\mathrm{vac}'\rangle$ is the ground state. This gives a description of phase transition as a vacuum realignment at the quantum level, which we shall use later to discuss Higgs-boson production with local vacuum excitations. Let us now reorder the Lagrangian (1) with respect to ϕ' operators. Then we have, e.g.,

$$:(D_\mu\phi)^*(D^\mu\phi): = N_{\phi'}((D_\mu\phi')^*(D^\mu\phi')) + \frac{g^2\xi^2}{2}A_\mu A^\mu$$
$$+ \frac{ig\xi}{\sqrt{2}}[A_\mu(\partial^\mu\phi') - A_\mu(\partial^\mu\phi'^*)] , \tag{13}$$

where $N_{\phi'}$ denotes normal ordering with respect to ϕ' operators. Thus gauge bosons now have a mass $g\xi_0$. Similarly, the mass of the physical scalar field becomes as usual $\sqrt{2\lambda}\xi$. Thus the conventional Higgs mechanism is reproduced through reordering with respect to the new quantum state $|\mathrm{vac}'\rangle$.

We now consider the Salam-Weinberg theory. Here

$$\mathcal{L} = -\tfrac{1}{4}F^a_{\mu\nu}F^{a\mu\nu} - \tfrac{1}{4}B_{\mu\nu}B^{\mu\nu} + |D_\mu\phi|^2 + \bar{\psi}_L\gamma^\mu\mathcal{D}_\mu\psi_L$$
$$+ \bar{e}_R\gamma^\mu(i\partial_\mu - g'B_\mu)e_R - \lambda_1(\bar{\psi}_L\phi e_R + \bar{e}_R\phi^\dagger\psi_L) - V(\phi) . \tag{14}$$

In the above,

$$V(\phi) = -m^2\phi^\dagger\phi + \lambda(\phi^\dagger\phi)^2 , \tag{15}$$

with ϕ being a complex scalar doublet given as

$$\phi = \begin{bmatrix} \phi^{(+)} \\ \phi^{(0)} \end{bmatrix} . \tag{16}$$

For the present description of the Higgs mechanism, we may now identify $\phi^{(0)}$ of Eq. (16) with ϕ of Eq. (1) and run through the algebra as quoted. We then obtain that $|\mathrm{vac}'\rangle$ defined in Eq. (7) is the description of electroweak theory with SSB and

$$\langle \mathrm{vac}'|\phi|\mathrm{vac}'\rangle = \begin{bmatrix} 0 \\ \frac{\xi}{\sqrt{2}} \end{bmatrix} .$$

The tree-level potential is then given by, parallel to (11),

$$V(\xi) = -\frac{1}{2}m^2\xi^2 + \frac{\lambda}{4}\xi^4 . \tag{17}$$

We can then see that the present gauge-covariant kinetic term $|D_\mu\phi|^2$ with respect to $|\mathrm{vac}'\rangle$ as in Eq. (13), now yields the usual masses of the gauge bosons and Higgs particles as

$$m_{W^\pm} = \frac{g\xi_0}{2}, \quad m_{Z^0} = \frac{(g^2+g'^2)^{1/2}}{2}\xi_0, \quad m_H = \sqrt{2\lambda}\xi_0 , \tag{18}$$

where $\xi^0 = (m^2/\lambda)^{1/2}$ is the field expectation value corresponding to the minimum of the potential as in Eq. (17). We note that the above results are identical with the results for the Higgs mechanism starting with an equation parallel to Eq. (9).

III. TEMPERATURE DEPENDENCE OF VACUUM

We shall now consider the temperature dependence of $|\mathrm{vac}'\rangle$, where the above nonperturbative solution of field theory corresponds to the zero temperature. This is clear when we note that for any operator $\hat{O}$,

$$\lim_{\beta\to\infty} \frac{\mathrm{Tr}[\exp(-\beta H)\hat{O}]}{\mathrm{Tr}[\exp(-\beta H)]} = \frac{\langle \mathrm{vac}'|\hat{O}|\mathrm{vac}'\rangle}{\langle \mathrm{vac}'|\mathrm{vac}'\rangle} , \tag{19}$$

where $|\mathrm{vac}'\rangle$ now is the *state of lowest energy* corresponding to zero temperature. We shall now generalize this to finite temperatures. For this purpose we shall use the methodology of thermo field dynamics.[2] In the Appendix, we summarize the salient features of the same. The idea is to calculate an effective potential at finite temperature as an expectation value over a "thermal vacuum" which at zero temperature reduces to Eq. (11). As shown in the Appendix, the thermal vacuum is given by

$$|\mathrm{vac}',\beta\rangle = U(\beta)|\mathrm{vac}'\rangle , \tag{20}$$

with

$$U(\beta) = \exp(B^\dagger - B) , \tag{21}$$

where thermal modes are created with[2]

$$B^\dagger = \int \theta(k,\beta)a(\mathbf{k})'^\dagger\tilde{a}(-\mathbf{k})^\dagger d\mathbf{k} . \tag{22}$$

In the above the operator $\tilde{a}(\tilde{a}^\dagger)$ are the annihilation (creation) operators in the extra Hilbert space and satisfy the same quantum algebra. We note that $|\mathrm{vac}\rangle$ of the Appendix is replaced by $|\mathrm{vac}'\rangle$ here, and we consider the thermal excitations over $|\mathrm{vac}'\rangle$, the stable vacuum after Salam-Weinberg phase transition at $T=0$. For bosons, the function $\theta(k,\beta)$ is given by[2]

$$\sinh^2\theta(k,\beta) = \frac{1}{\exp[\beta\omega(k,\beta)]-1} , \tag{23}$$

corresponding to the Bose distribution of the number operator $a'(\mathbf{k})^\dagger a'(\mathbf{k})$ for the physical particles and commutator algebra. The Bogoliubov transformation corresponding to Eq. (A6) becomes

$$\begin{bmatrix} a(\mathbf{k},\beta)' \\ \tilde{a}(-\mathbf{k},\beta)^\dagger \end{bmatrix} = \begin{bmatrix} \cosh\theta(k,\beta) & -\sinh\theta(k,\beta) \\ -\sinh\theta(k,\beta) & \cosh\theta(k,\beta) \end{bmatrix} \begin{bmatrix} a(\mathbf{k})' \\ \tilde{a}(-\mathbf{k})^\dagger \end{bmatrix} . \tag{24}$$

A few remarks regarding construction of thermal vacuum as in Eq. (20) may be relevant. Clearly, at zero temperature, with $\beta=\infty$, the Hilbert space consisting of only physical particles decouples in the extended Hilbert space. Further, $|\text{vac}',\beta\rangle$ contains Higgs-particle excitations as is clearly seen from Eqs. (21) and (22). In principle we should include the gauge bosons and fermion excitation as well in $B^\dagger$. As a first approximation we ignore these channels to illustrate the qualitative features and possible experimental signatures with only the Higgs sector.

As is clear from the Lagrangian in Eq. (14), and the fact that $|\text{vac}',\beta\rangle$ contains only ϕ_1' quanta, the contribution to energy density will come from ϕ_1'-dependent terms only when expectation value of the Hamiltonian with respect to thermal vacuum is taken. We thus have the expression for the *effective Hamiltonian* density $\mathcal{T}^{00}_{\text{eff}}$ as

$$\mathcal{T}^{00}_{\text{eff}} = \tfrac{1}{2}\left[\dot{\phi}_1'^2 + \phi_1'(-\nabla^2)\phi_1' + \frac{\lambda}{2}\phi_1'^4 + (3\lambda\xi^2 - m^2)\phi_1'^2 + \frac{\lambda}{2}\xi^4 - m^2\xi^2 \right] . \tag{25}$$

The energy density at temperature β now becomes

$$V(\xi,\beta) \equiv \epsilon(\beta) = \langle \text{vac}',\beta | \mathcal{T}^{00}_{\text{eff}} | \text{vac}',\beta \rangle = \tfrac{1}{2}(2\pi)^{-3}\int \frac{\omega(k,\beta)^2 + \mathbf{k}^2 + 3\lambda\xi^2 - m^2}{\omega(k,\beta)\{\exp[\beta\omega(k,\beta)]-1\}} d\mathbf{k} + \frac{3\lambda}{4}\left[(2\pi)^{-3}\int \frac{1}{\omega(k,\beta)\{\exp[\beta\omega(k,\beta)]-1\}} d\mathbf{k} \right]^2 + \frac{\lambda}{4}\xi^4 - \frac{m^2}{2}\xi^2 , \tag{26}$$

where

$$\omega(k,\beta) = (\mathbf{k}^2 + m_H(\beta)^2)^{1/2} , \tag{27}$$

with $m_H(\beta)$ being the Higgs-boson mass at temperature $1/\beta$. Equation (26) is an implicit definition for the effective potential with the parameters to be determined self-consistently as explained below.

In the thermo field method, temperature-dependent field theory needs the physical mass of the Higgs particle as input while using the distribution function in the calculation of temperature-dependent effective potential as in the right-hand side of Eq. (26) through $\omega(k,\beta)$. At finite temperatures Lorentz invariance is broken and thus mass is not well defined. We shall, however, still follow the method of defining mass through the second derivative of the effective potential at its minimum along with a self-consistency requirement. The mass of the physical Higgs particles shall be taken as $d^2V(\xi,\beta)/d\xi^2|_{\xi=\xi_{\min}}$, where $\xi_{\min}$ corresponds to the expectation value of the field for the minimum of effective potential as on the left-hand side of Eq. (26). However, for the calculation of $V(\xi,\beta)$ in Eq. (26) we need a value for the masslike parameter $m_H(\beta)$ on the right-hand side. We shall call the mass on the right-hand side of Eq. (26) for such calculations as the input mass and the square root of the expression $d^2V(\xi,\beta)/d\xi^2|_{\xi=\xi_{\min}}$ after $V(\xi,\beta)$ is evaluated as the output mass. Self-consistency demands that both be the same. Such a determination of mass through self-consistency is very much like solving the integral equation for self-energy in perturbative calculations with self-energy both as an input as well as an output. The situation here is, however, technically different in the sense that the effective mass enters on the right-hand side of Eq. (26) for using thermal distributions, and on the left-hand side through a second-order derivative of the potential at its minimum. We determine $m_H(\beta)$ in Eq. (26) through an iterative procedure until the input mass equals the output mass. The methodology here is not a loop calculation, but self-consistency simulates the dynamical effects of something like loop calculations. We note that this definition, however, is not the same as through the pole of the propagator and an alternative definition of the mass might give somewhat different results.

For numerical evaluations it is useful to rewrite Eq. (26) in terms of the dimensionless quantities with the substitutions

$$z = \frac{\xi}{\xi_0}, \quad \mu = \frac{m_H(\beta)}{\xi_0}$$

and

$$y = \beta\xi_0 , \tag{28}$$

where $\xi_0 = (m^2/\lambda)^{1/2}$ is the value of $\xi_{\min}$ for *zero temperature*. The expression for the effective potential now becomes

$$V(z,y) = \xi_0^4\left[\frac{\lambda}{4}z^4 - \frac{\lambda}{2}z^2 + \frac{1}{2}I_1(z,y) + \frac{3\lambda}{4}[I_2(z,y)]^2 \right] \equiv \xi_0^4 V_1(z,y) , \tag{29}$$

where

$$I_1(z,y) = \frac{1}{2\pi^2}\int_0^\infty \frac{x^2[\omega(x)^2 + x^2 + \lambda(3z^2-1)]}{\omega(x)\{\exp[y\omega(x)]-1\}} dx \tag{30a}$$

and

$$I_2(z,y) = \frac{1}{2\pi^2}\int_0^\infty \frac{x^2}{\omega(x)\{\exp[y\omega(x)]-1\}} dx \tag{30b}$$

with

$$\omega(x)=(x^2+\mu^2)^{1/2} . \tag{30c}$$

In the above μ is the Higgs-particle mass in ξ_0 units and as stated is to be determined self-consistently. $V(z,y)$ as a function of ξ and β is accepted as the effective potential at finite temperature only after an iterative determination of μ or $m_H(\beta)$ through self-consistency. The results of the calculations are stated below.

In Fig. 1 we have plotted $V(\xi,\beta)-V(0,\beta)$ for different temperatures. For $T=T_c\approx 2.1\xi_0$, the shape of the potential changes, which thus determines the critical temperature T_c, which is the same as obtained by Dolan and Jackiw.[6] In Fig. 2 we have then plotted in curve I $m_H(\beta)^2$ as a function of temperature T. For $T=T_c$, $m_H(\beta)^2$ goes to zero as expected. However, it again rises for temperatures $T>T_c$. We have also plotted in curve II of the same figure $m^2(\beta)\equiv d^3V(\xi,\beta)/d\xi^2|_{\xi=0}$ corresponding to the "mass square" of the effective Lagrangian at $\xi=0$ for the double-well potential, which as expected is negative below T_c and becomes the same as $m_H(\beta)^2$ above T_c as $\xi_{\min}=0$. We have next plotted the order parameter $\xi_{\min}$ as a function of temperature in curve III of the same figure. The order parameter goes to zero as temperature approaches the critical temperature T_c. T_c could be determined through any of these three curves and turns out to be $2.1\xi_0\simeq 525$ GeV. We did not use here loop expansion or a high-temperature approximation[6] but used the thermofield method with a self-consistent variational ansatz giving results similar to those of Ref. 6. It is reassuring that the present nonperturbative numerical technique yields the same result in spite of the expressions being different. The gap in energy density of the thermal vacuum with respect to the vacuum at zero temperature is given by

$$\Delta\epsilon(\beta)=V(\xi_{\min},\beta)-V(\xi_{\min},\beta=\infty)$$
$$=\xi_0^4\left[V_1(z_{\min},y)+\frac{\lambda}{4}\right] . \tag{31}$$

The number density of Higgs particles at temperature β is here given as

$$N(\beta)=\langle \text{vac}',\beta|a(z)^\dagger a(z)|\text{vac}',\beta\rangle$$
$$=(2\pi)^{-3}\int\frac{1}{\exp[\beta\omega(k,\beta)]-1}d\mathbf{k} . \tag{32}$$

IV. LOCAL DESTABILIZATION OF VACUUM

We shall now consider possible local "heating" of the vacuum with particle collision, which, from the very nature of the problem, will be intuitive and heuristic and can only lead to qualitative conclusions. As in condensed-matter physics, such a destabilization of a vacuum can occur if we pump in enough energy into a small macroscopic volume which thermalizes *locally* and becomes hot. The dynamics of such a "bubble" formation is known to be complex,[7] and shall not be tackled here. We shall rather consider the possibility of excitation of a

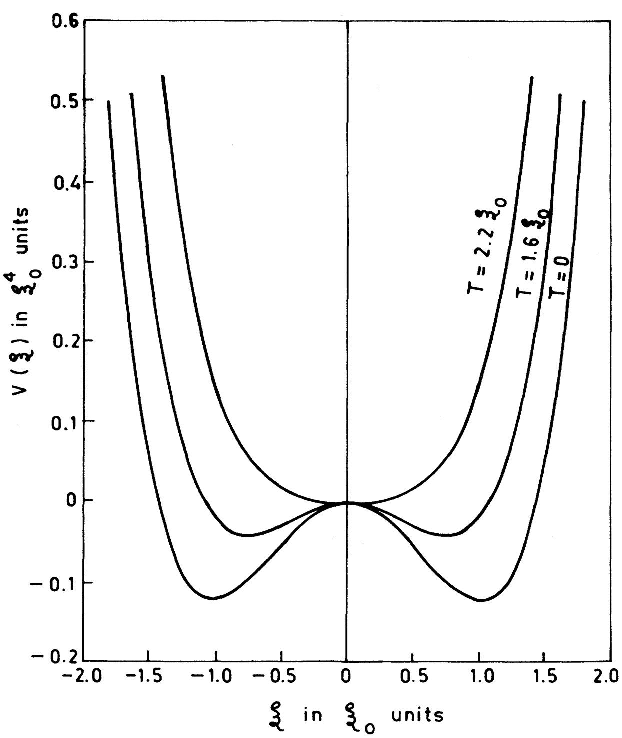

FIG. 1. Effective potential $V(\xi,\beta)-V(0,\beta)$ has been plotted at different temperatures $T=1/\beta$. Critical temperature $T_c\approx 2.1\xi_0\approx 525$ GeV.

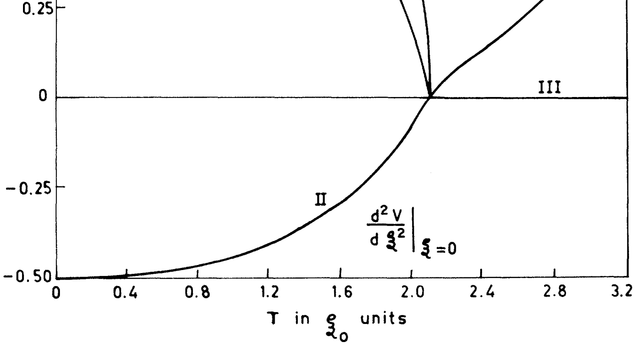

FIG. 2. In curves I, II, and III we have plotted $m_H(\beta)^2$, $(d^2V/d\xi^2)|_{\xi=0}$ and $\xi_{\min}(\beta)$, respectively, as functions of temperature in units of $\xi_0\approx 250$ GeV. Critical temperature T_c can be determined from any of the above curves as $2.1\xi_0$.

vacuum with a bubble formation which has a nonzero temperature and local thermal equilibrium. With vacuum as the medium in which collisions take place, we conjecture here that during particle collisions a part of the collision energy might excite the vacuum locally.

We shall now apply the above ideas to such a situation with temperature defined locally inside the bubble. Thus the energy density gap is no longer constant throughout the volume but is maximum at center of collision and decreases to zero away from it. The total energy of such a locally excited region or the bubble shall be given as

$$E_B=\int \Delta\epsilon(\beta(\mathbf{r}))d\mathbf{r} , \tag{33}$$

where $\Delta\epsilon(\beta(\mathbf{r}))$ is the same function as in Eq. (32) except that β is spatially dependent. Similarly, the number of Higgs particles inside the bubble becomes

$$N_B=\int N(\beta(\mathbf{r}))d\mathbf{r} . \tag{34}$$

As stated, the determination of $\beta(\mathbf{r})$ as a function of $\mathbf{r}$ from first principles is impossible. We shall therefore look into the qualitative structure only through a Gaussian distribution for the same. For the temperature distribution inside the bubble we shall therefore take

$$\beta(r)^{-1}=T(r)=T_0\exp(-ar^2) . \tag{35}$$

In the above, T_0 is the temperature at the center of bubble and the parameter a essentially decides the region over which vacuum is excited with the corresponding "volume" being approximately $a^{-(3/2)}$.

We have noted in Eq. (32) that such a bubble will contain Higgs particles. Such particles in Salam-Weinberg theory are, however, coupled to fermions and get converted to quark or lepton pairs[8] as the bubble cools. The Higgs particles in the bubble will primarily go to heavy-fermion–antifermion pairs. The decay width of such a process for free Higgs particles at finite temperature is given as[8]

$$\Gamma(\beta)=\tan\left[\frac{\beta m_H(\beta)}{4}\right]\Gamma_0 , \tag{36}$$

where Γ_0 is the decay width for $H\to f\bar{f}$ at zero temperature. Then the average decay width of the Higgs particles in the bubble may be calculated to be

$$\Gamma_{\rm avg}=\Gamma_0\int \tanh\left[\frac{\beta m_H(\beta(r))}{4}\right]N_B(\beta(r))d\mathbf{r} . \tag{37}$$

Inside the bubble, the masses of the fermions decrease in the same manner as the vacuum expectation value (VEV) or mass of the Higgs particles, since the masses of the fermions are generated through the vacuum expectation values. As they come out of the bubble, these will become more massive as the temperature outside decreases, resulting in a cooling of the bubble. The average mass of the *decaying* Higgs particles will depend on both their lifetimes as well as the number of Higgs particles with a specific mass. Thus we may calculate this average mass as

$$M_{\rm avg}=\int N_B(\beta(r))\frac{\Gamma(\beta(r))}{\Gamma_{\rm avg}}m_H(\beta(r))d\mathbf{r} = \frac{\int N_B(r)\tanh[\beta m_H(\beta(r))/4]m_H(\beta(r))d\mathbf{r}}{\int N_B(\beta(r))\tanh[\beta m_H(\beta(r))/4]d\mathbf{r}} , \tag{38}$$

which will correspond to the average transverse momentum of the decaying particles at any fixed time. This, however, will be a dynamic phenomenon as the bubble cools here mainly through the decay of Higgs particles with a corresponding loss of energy. To get an idea of the mass distribution inside the bubble we may calculate the variance of the same as

$$\sigma^2=[(M^2)_{\rm avg}-(M_{\rm avg})^2] , \tag{39}$$

where $(M^2)_{\rm avg}$ is given by

$$(M^2)_{\rm avg}=\int N_B(\beta(r))\frac{\Gamma(\beta(r))}{\Gamma_{\rm avg}}m_H(\beta(r))^2d\mathbf{r} . \tag{40}$$

We note that the conjecture quoted in Eq. (35) for the temperature distribution in a bubble has not been so far utilized. We shall now explicitly take it to get a qualitative idea of possible experimental signals. For this purpose, the total energy as in Eq. (33) should be available for vacuum excitation from some collision process. Further, the number of Higgs particles in the bubble as in Eq. (34) should be large enough for thermal equilibrium to be possible. With this in mind, we now look for qualitative experimental signals.

V. EXPERIMENTAL SIGNATURES

The experimental signatures of bubble formation as compared to *heating* of any material in condensed-matter physics is extremely sharp, since the bubble here has energy dissipation through the conversion of Higgs particles to fermion and antifermion pairs, which has no earlier parallel. The signature here is particularly strong since there will be a *preferential* production of heavy-fermion and antifermion pairs, as the coupling of fermions to Higgs particles is proportional to their masses. Such a preferential production of heavy fermions is absent for the conventional mechanisms for particle production. Hence, irrespective of the details of the modeling, excessive heavy-fermion production in a collision can be a signature for bubble formation. The explicit modeling will decide whether we could anticipate to observe the same anywhere, or interpret existing data with the association of Higgs particles. Before doing this, however, we shall note some preliminary calculations for different E_B and a in Eqs. (33) and (35) as inputs. We shall consider the possibility of a destabilization of vacuum over a volume $\approx a^{-3/2}$ due to a collision. A natural ansatz here appears to be the length scale associated with the cross section of the collision process. Hence, for vacuum excitation, we shall tentatively assume that $\sigma\simeq\pi/a$, and compute the possibility of signals. It is quite possible that such a hypothesis is not correct. If so, the scales quoted below will change, and the process might be visible in a different way or not visible at all. The above assumption is made for conceptual convenience and seems to correspond to some signals in cosmic rays.

For definiteness let us consider a cross section of the order of microbarns, which corresponds to a value $0.02\xi_0^2$ for a in Eq. (35). The corresponding bubble energy versus number of Higgs particles is plotted for $\lambda \simeq 0.02$ corresponding to Higgs-boson mass of 50 GeV in Fig. 3 as curve I. The results are not very sensitive to λ except for kinematics. We observe that the multiplicity increases almost linearly with bubble energy for E_B larger than 1 TeV. We note that for E_B greater than 2–3 TeV, there is a reasonable number of Higgs particles. The possibility of vacuum excitation will depend upon the fraction of the total energy of collision that goes to E_B for local excitation and bubble formation. What this fraction is and how it depends on energy is not known. In curve II of Fig. 3 we plot the central temperature against bubble energy E_B for N_B around and larger than 30 since thermal effects for small N_B will not be sensible. We note that for temperatures of the order of 150 GeV, which is much less than critical temperature of 525 GeV, the number of Higgs particles appears to be sizable, as may be seen in Fig. 3. We note here that with the inclusion of gauge particles and fermions T_c might become smaller than 525 GeV. In such a case, the present mechanism would show up at a lower temperature. This is because N_B will become appreciable at a lower temperature since the Higgs-boson mass at that temperature will be less compared to the present case.

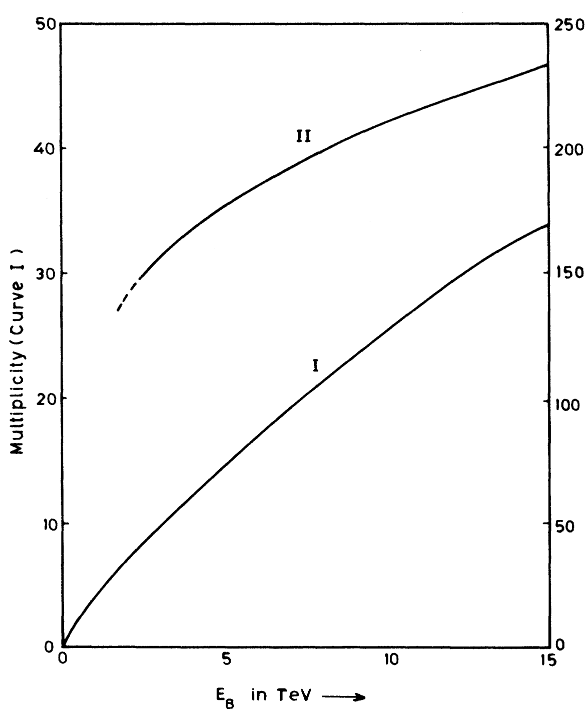

FIG. 3. In curve I we plot Higgs-particle multiplicity N_B against bubble energy E_B in units of TeV corresponding to microbarn size. In curve II we plot the central temperature T_0 up to 250 GeV against bubble energy E_B.

The bubble energy will depend on the fraction of collision energy that goes towards vacuum excitation. This is not known. However, we note from Fig. 3 that for E_B around 2 TeV or more, the above process could become relevant. Curve I here gives the multiplicity for on-mass-shell Higgs particles, which is temperature dependent. Such a bubble will primarily lose energy and cool rapidly through Higgs particles getting converted to fermion pairs dynamically, thus using up bubble energy. This production of fermions shall include the fact that the fermion masses are proportional to the Higgs-boson masses above. As the fermions come out, the bubble cools and the fermions become more massive. The characteristic feature of such a process leading to Higgs particle decay yields, e.g., more high-P_T muon pairs and an overall excess of muons due to heavy flavor production. This will be larger than the expected probability for them because of preferential heavy-flavor production for Higgs-particle "decay." Such events shall be rare, but could be just visible in accelerator experiments beyond the TeV range. As stated, these will have unusually high multiplicity and an excess of muons, which shall signify the onset of the above mechanism. We note that the above comments are qualitative, but shall leave visible signatures in spite of the dynamics not being understood.

Since as per the present calculations, the accelerator energies are inadequate, we state below some signatures of this type in cosmic rays. For this purpose let us examine the following events which appear not to be capable of being explained with simulation programs of known physics.

(i) Chiron events. These events[9] were observed by the Brasil-Japan Collaboration in Chacaltaya Emulsion Chamber experiment in chambers 19 and 21 and have the following characteristics. The total P_T is of the order of 5 GeV there. Further inside them we have "miniclusters" associated with as small a P_T as $\simeq 10$–20 MeV. These, we note, may be associated with Higgs-particle production as above. In the present scenario, as may be seen from Eq. (38), the high P_T of the order of 5 GeV will correspond to decay of the Higgs particle. These particles will produce heavy flavors and hence the "miniclusters" of 10–20-MeV transverse momentum might indicate D^* production. Production of miniclusters along with absence of π^0 could thus be interpreted as signature of Higgs-particle productions in these events.

One more observation in this context may be relevant. The total transverse momentum of the chirons will correspond to the mass of the Higgs particle when it decays. Higgs particles could decay during the process of cooling of the bubble even before zero temperature is reached. Thus the present identification of chiron events where large P_T comes from the decay of Higgs particles to heavy-fermion pairs does not determine its mass.

(ii) Halo events. Halo events in cosmic rays are some events with excessive multiplicity where the particle number could not be counted. As is clear from the E_B versus N_B plots, the multiplicity of Higgs particles rises almost linearly with the bubble energy. Hence, for energies beyond the threshold for the present process, the multiplicity will increase linearly with energy which is

much faster than what can be expected from ordinary physics. Hence the halo events[10] may indicate the onset of new physics through an unusual rise in the multiplicity resulting from excitation of vacuum. Further, the multiple cores[10] in halos just look like multiple bubble formation.

(iii) Cygnus X-3, Hercules X-1 signals. These are signals from extremely high-energy air showers originating from the directions of Cygnus X-3 and of Hercules X-1 which have a high muon content,[11] thus indicating *hadronic* interactions. It is then presumed that a neutral *stable* hadronic particle, e.g., cygnets, quark nuggets, etc. should be coming from Cygnus X-3 and Hercules X-1.[12] Instead, we propose that *local vacuum excitations* with Higgs-particle production as above could give rise to the excess muon signals through preferential heavy-flavor production.[14]

VI. CONCLUSIONS

The present mechanism of Higgs-particle production consists of three parts: (a) The bubble will get formed; (b) the temperature-dependent vacuum will have a quantum-mechanical description with production of off-mass-shell Higgs particles, and (c) the dissipation of the bubble takes place through particle production via conversion of Higgs particles to fermion pairs. In the present note we have concentrated on (b), i.e., the description of temperature-dependent local excitation of the vacuum. It is possible to obtain a signal for the process because of *preferential* heavy-flavor production during the dissipation of the bubble. Processes (a) and (c) are not only nonperturbative but also time dependent, and an understanding of the same is a nontrivial step.[7] However, heavy-flavor production for (c) leaves a trail which can be always seen if we can look for this.

A comment regarding thermalization may be relevant. Production of a "coherent bubble" as in Eq. (7) may not be instantaneous corresponding to a new value $\xi(\beta)$. However, time scales for local thermalization of a vacuum is not known and will be associated with bubble formation.[7] Once the existence of the process is confirmed with more observations, the time scales involved for this process for thermalization and the response of vacuum regarding thermal and transport properties for the same as a medium, need to be investigated and understood.

We note that with conventional consideration of field theory with high-temperature approximation, a careful and complete calculation has been done by Ferrer, de la Incera, and Shabad,[15] where three possible phases have been identified, including one with W condensates. Our calculations are complementary in the sense that the high-temperature limit is not taken and that experimental signatures at temperatures less than critical temperature could be observed. We may further observe that a recent analysis by Goldberg with scalars only shows the instability of the perturbative tree-level calculations regarding cross sections at ultrahigh energies.[16] In this context or otherwise, the nonperturbative techniques developed here may be more relevant and we may look for signatures of heavy-flavor production or unusual rise in multiplicity.

Thus, to sum up, the characteristic features of Higgs-particle production through vacuum excitation would be (i) relatively high P_T for Higgs-particle decay, (ii) preferential production of heavy flavors, and (iii) an unusual rise in multiplicity with energy as a result of this mechanism. We find that the observations of chirons and halo events as well as the hardonic signal associated with the direction of Cygnus X-3 and Hercules X-1 could indicate the onset of the present mechanism, and observation of the same in accelerators may be possible not too far in the future. It is significant to note that even though Salam-Weinberg phase transition takes place around a temperature of more than 500 GeV, the effects of local heating of a vacuum could be seen for temperatures as low as 150–200 GeV with the corresponding bubble energy being a few TeV.

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APPENDIX

We here summarize briefly the salient features of thermofield dynamics as used here.

In statistical mechanics, the thermal average of an operator $\hat{O}$ is given as, with $\beta = 1/kT$,

$$\langle \hat{O} \rangle_\beta = \frac{\mathrm{Tr}(e^{-\beta H}\hat{O})}{\mathrm{Tr}(e^{-\beta H})} , \tag{A1}$$

where the trace is taken over a complete basis of states. First we note that, in the zero-temperature limit, the above reduces to ground-state expectation value for the operator $\hat{O}$. This is easily seen as

$$\begin{aligned}\lim_{\beta\to\infty} \langle \hat{O} \rangle_\beta &= \lim_{\beta\to\infty} \frac{\langle 0|\hat{O}|0\rangle e^{-\beta\epsilon_0} + \langle 1|\hat{O}|1\rangle e^{-\beta\epsilon_1} + \cdots}{e^{-\beta\epsilon_0} + e^{-\beta\epsilon_1} + \cdots} \\ &= \lim_{\beta\to\infty} \frac{\langle 0|\hat{O}|0\rangle + \langle 1|\hat{O}|1\rangle e^{-\beta(\epsilon_1-\epsilon_0)} + \cdots}{1 + e^{-\beta(\epsilon_1-\epsilon_0)} + \cdots} \\ &= \langle 0|\hat{O}|0\rangle ,\end{aligned} \tag{A2}$$

where $|0\rangle$ corresponds to the state with the lowest energy. In thermo field method, one essentially generalizes (A2) to the case of finite temperature and defines a thermal vacuum such that the statistical average reduces to an expectation value with respect to the thermal vacuum. Thus we want that for some $|0(\beta)\rangle$ the relationship

$$\langle \hat{O} \rangle_\beta = \frac{\mathrm{Tr}(e^{-\beta H}\hat{O})}{\mathrm{Tr}(e^{-\beta H})} \equiv \langle 0(\beta)|\hat{O}|0(\beta)\rangle , \tag{A3}$$

where $|0(\beta)\rangle$ is defined as the thermal vacuum. This can be done if one doubles the degrees of freedom, i.e., corresponding to every physical operator A, a "tilde" operator

$\tilde{A}$ is introduced. For example, we generalize the free boson Hamiltonian as[2]

$$\tilde{H}_0 = \int d\mathbf{k}\, \omega(k)[a(k)^\dagger a(k) - \tilde{a}(k)^\dagger \tilde{a}(k)] \; . \qquad \text{(A4)}$$

In the above, $\tilde{a}(\mathbf{k})$ are the new operators named as "thermal modes." They are associated with negative energy, with conventional quantization, but do not have any physical significance in the sense of observation of these modes. In a zero-temperature vacuum, these modes are absent, so that conventional field theory holds. At finite temperature the ground state is replaced by $|0(\beta)\rangle$,

$$|0(\beta)\rangle \equiv U(\beta)|\text{vac}\rangle = \exp\left[\int \theta(\mathbf{k},\beta)[a(\mathbf{k})^\dagger \tilde{a}(-\mathbf{k})^\dagger - \text{H.c.}]d\mathbf{k}\right]|\text{vac}\rangle, \qquad \text{(A5)}$$

where $\tilde{a}(-\mathbf{k})^\dagger$ in the above corresponds to the extra Hilbert space. It is now convenient to define a thermal basis

$$\begin{bmatrix} a(\mathbf{k},\beta) \\ \tilde{a}(-\mathbf{k},\beta)^\dagger \end{bmatrix} = U(\beta) \begin{bmatrix} a(\mathbf{k}) \\ \tilde{a}(-\mathbf{k})^\dagger \end{bmatrix} U(\beta)^{-1} \; , \qquad \text{(A6)}$$

which amounts to the Bogoliubov transformation

$$\begin{bmatrix} a(\mathbf{k},\beta) \\ \tilde{a}(-\mathbf{k},\beta)^\dagger \end{bmatrix} = \begin{bmatrix} \cosh\theta(\mathbf{k},\beta) & -\sinh\theta(\mathbf{k},\beta) \\ -\sinh\theta(\mathbf{k},\beta) & \cosh\theta(\mathbf{k},\beta) \end{bmatrix} \begin{bmatrix} a(\mathbf{k}) \\ \tilde{a}(-\mathbf{k})^\dagger \end{bmatrix} \qquad \text{(A7)}$$

$a(\mathbf{k},\beta)$ and $\tilde{a}(\mathbf{k},\beta)$ are the annihilation and creation operators at temperature $\beta = 1/KT$ corresponding to the thermal vacuum such that $a(\mathbf{k},\beta)|0(\beta)\rangle = 0 = \tilde{a}(\mathbf{k},\beta)|0(\beta)\rangle$. We next take the function $\theta(k,\beta)$ in (A5) such that, for bosons,

$$\langle 0(\beta)|a(\mathbf{z})^\dagger a(\mathbf{z})|0(\beta)\rangle = (2\pi)^{-3}\int d\mathbf{k}\,\sinh^2\theta(k,\beta) \equiv (2\pi)^{-3}\int \frac{d\mathbf{k}}{e^{\beta\omega(k,\beta)}-1} \; ; \qquad \text{(A8)}$$

i.e., the expectation value of the number operator reproduces the Bose distribution. This determines $\theta(k,\beta)$ to be given by

$$\sinh^2\theta(k,\beta) = \frac{1}{e^{\beta\omega(k,\beta)}-1} , \qquad \text{(A9)}$$

so that when statistics are known, the corresponding Bogoliubov transformation relating the zero-temperature ground state with the thermal ground state of the extended Hilbert space is known. The ground state or the thermal vacuum obviously contains particles with appropriate distributions as in Eq. (A8). A self-consistent determination of mass in Eq. (A9) shall be sometimes needed if we wish to include the effect of interactions through a masslike parameter in taking $\omega(k,\beta) = (k^2 + m(\beta)^2)^{1/2}$. In the present paper we do this.

The methodology enables us to replace mixed states of statistical mechanics by pure states in an extended Hilbert space while generating correct distribution functions. The extra thermal modes enable us to do this.

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