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Single and pair production of doubly charged Higgs bosons at hadron colliders

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Current searches for doubly charged Higgs bosons ($H^{\pm\pm}$) at the Fermilab Tevatron are sensitive to single production of $H^{\pm\pm}$, although the pair production mechanism $q\overline{q} \to H^{++}H^{--}$ is assumed to be dominant. In the context of a Higgs Triplet Model we study the mechanism $q'\overline{q} \to H^{\pm\pm}H^{\mp}$ at the Tevatron and CERN Large Hadron Collider, and show that its inclusion can significantly improve the search potential for $H^{\pm\pm}$. Moreover, assuming that the neutrino mass is generated solely by the triplet field Yukawa coupling to leptons, we compare the branching ratios of $H^{\pm\pm} \to l^{\pm}l^{\pm}$ and $H^{\pm\pm} \to H^{\pm}W^{*}$ for the cases of a normal hierarchical, inverted hierarchical, and degenerate neutrino mass spectrum.

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I. INTRODUCTION

The quest for Higgs bosons is of utmost importance at high energy colliders [1–3]. In the standard model (SM), one isospin $I = 1/2$, hypercharge $Y = 1$ complex scalar doublet breaks the electroweak symmetry and provides mass for the fermions, $W^{\pm}$ and Z. One neutral scalar ϕ^0 remains as a physical degree of freedom—"the SM Higgs boson." Such a framework predicts $\rho(= M_W^2/M_Z^2\cos^2\theta_W) = 1$ at tree level, a result which is in impressive agreement with the experimental measurement of $\rho \approx 1$ [4]. More generally, any Higgs sector composed solely of $I = 1/2$, $Y = 1$ doublets assures $\rho = 1$ at tree level, with calculable one-loop corrections [5].

Predicting $\rho = 1$ at tree level is certainly an attractive feature of $I = 1/2$, $Y = 1$ doublet representations, although models with isospin triplets ($I = 1$) also can be considered [1]. Such models have various virtues and deficiencies. If the neutral member of the triplet acquires a vacuum expectation value (VEV) then $\rho = 1$ at tree level is no longer guaranteed, and the triplet VEV must be very small in order to comply with the measured value $\rho \approx 1$. However, unlike doublets, $Y = 2$ triplets can give rise to neutrino masses and mixings whose magnitude is proportional to the triplet vacuum expectation value multiplied by an arbitrary Yukawa coupling (h_{ij}) *without* invoking a right-handed neutrino [6,7].

A clear phenomenological signature of $Y = 2$ triplets would be the observation of a doubly charged Higgs boson $H^{\pm\pm}$. Such $H^{\pm\pm}$ have been searched for at the e^+e^- collider LEP, resulting in mass limits of the order $m_{H^{\pm\pm}} > 100$ GeV [8–11]. Their existence also can affect a wide variety of processes, such as Bhabha scattering, the anomalous magnetic moment of the muon $(g-2)_{\mu}$, and lepton flavor violating $\mu^{\pm}$ and $\tau^{\pm}$ decays [12–17]. The Fermilab Tevatron recently performed the first search for $H^{\pm\pm}$ at hadron colliders. The production process $p\overline{p} \to \gamma, Z \to H^{++}H^{--}$ was assumed, with subsequent decay $H^{\pm\pm} \to l^{\pm}l^{\pm}$. D0 Collaboration [18] searched for $H^{\pm\pm} \to \mu^{\pm}\mu^{\pm}$ while CDF Collaboration [19] searched for three final states $H^{\pm\pm} \to \mu^{\pm}\mu^{\pm}$, $\mu^{\pm}e^{\pm}$, and $e^{\pm}e^{\pm}$. Mass limits of the order $m_{H^{\pm\pm}} > 130$ GeV were obtained with an integrated luminosity of 240 pb^{-1}, assuming the branching ratio (BR) of $(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm})$ is 100% [19] in a given channel. These are the strongest direct mass limits on any type of Higgs boson, which shows the strong search capability of hadron colliders in the channel $H^{\pm\pm} \to l^{\pm}l^{\pm}$.

Given this strong search potential, in this paper we consider the phenomenological effect of relaxing these simplifying assumptions for the dominant production mechanism and decay modes of $H^{\pm\pm}$. Although work along these lines has appeared previously [14,20–24] we develop and expand the preceding analyses. For example, if h_{ij} are solely responsible for the currently favored form of the neutrino mass matrix then $\mathrm{BR}(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm}) <$ 100% in a given channel [14]. In this paper we study in detail the alternative production mechanism $q'\overline{q} \to W^* \to H^{\pm\pm}H^{\mp}$ [24], which can be as large as $q\overline{q} \to \gamma, Z \to H^{++}H^{--}$. Since the current search strategy at the Tevatron is in fact sensitive to *single* production of $H^{\pm\pm}$, we introduce the inclusive single production cross section ($\sigma_{H^{\pm\pm}}$) as the sum of the single and pair production cross sections. We point out that the contribution of $q'\overline{q} \to W^* \to H^{\pm\pm}H^{\mp}$ to $\sigma_{H^{\pm\pm}}$ strengthens the Tevatron mass limit on $H^{\pm\pm}$, which in general has a dependence on $m_{H^{\pm}}$. Moreover, we quantify the impact of the potentially important decay mode $H^{\pm\pm} \to H^{\pm}W^{*}$ [22] in the light of recent neutrino data. Although such a decay can weaken the $H^{\pm\pm}$ search capability in the leptonic channel, observation of $H^{\pm\pm} \to H^{\pm}W^{*}$ together with one or more leptonic channels might permit an order of magnitude estimate of h_{ij} [21,23].

Our work is organized as follows: In Sec. II we introduce the Higgs triplet model. In Sec. III we study the production mechanism $q'\overline{q} \to H^{\pm\pm}H^{\mp}$ and its phenomenological effect on the $H^{\pm\pm}$ search at the Tevatron and LHC. In Sec. IV we quantify the impact of the decay $H^{\pm\pm} \to H^{\pm}W^{*}$, while Sec. V considers the search potential of the Tevatron in the

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generalized scenario. Finally, in Sec. VI we present our conclusions.

II. THE HIGGS TRIPLET MODEL

Higgs $I = 1$ triplet representations arise in several well-motivated models of physics beyond the SM [1,20]. For example, left-right (L-R) symmetric models built on the gauge group $SU(2)_R \times SU(2)_L \times U(1)$ contain both left- and right-handed $I = 1$, $Y = 2$ triplet representations. Such models also require extra gauge bosons and can provide naturally light neutrino masses via the seesaw mechanism. Little Higgs models [25] also require $I = 1$, $Y = 2$ triplet representations, as well as new gauge bosons and fermions. However, Higgs triplets can be considered as a minimal addition to the SM [26]—for a review see [27]. We will focus on a particularly simple model [6,7], which merely adds a $I = 1$, $Y = 2$ complex (left-handed) Higgs triplet to the SM Lagrangian, hereafter referred to as the "Higgs triplet model" or "HTM." Such a model can provide a Majorana mass for the observed neutrinos without the need for a right-handed neutrino via the gauge invariant Yukawa interaction:[1]

$$\mathcal{L} = h_{ij}\psi_{iL}^T C i\tau_2 \Delta \psi_{jL} + \text{H.c.} \tag{1}$$

Here $h_{ij}(i, j = 1, 2, 3)$ is an arbitrary coupling, C is the Dirac charge conjugation operator, $\psi_{iL} = (\nu_i, l_i)_L^T$ is a left-handed lepton doublet, and Δ is a 2×2 representation of the $Y = 2$ complex (left-handed) triplet fields:

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}. \tag{2}$$

The Higgs potential [14] is as follows, with $\Phi = (\phi^+, \phi^0)^T$:

$$\begin{aligned} V = {}& m^2(\Phi^\dagger\Phi) + \lambda_1(\Phi^\dagger\Phi)^2 + M^2\,\mathrm{Tr}(\Delta^\dagger\Delta) \\ &+ \lambda_2[\mathrm{Tr}(\Delta^\dagger\Delta)]^2 + \lambda_3 \det(\Delta^\dagger\Delta) \\ &+ \lambda_4(\Phi^\dagger\Phi)\,\mathrm{Tr}(\Delta^\dagger\Delta) + \lambda_5(\Phi^\dagger\tau_i\Phi)\,\mathrm{Tr}(\Delta^\dagger\tau_i\Delta) \\ &+ \left(\frac{1}{\sqrt{2}}\mu(\Phi^T i\tau_2\Delta^\dagger\Phi) + \text{H.c.}\right). \end{aligned} \tag{3}$$

The term $\mu\Phi\Delta\Phi$, where μ is a dimensionful trilinear coupling, gives rise to a VEV v_Δ for the neutral member of the triplet Δ^0:

$$v_\Delta \simeq \mu v^2/2M^2. \tag{4}$$

Here M is the common triplet mass ($M^2\Delta^\dagger\Delta$). Since we are interested in the case of light triplets we take $M \approx v$, and so $v_\Delta \approx \mu$. A nonzero v_Δ gives rise to the following mass matrix for neutrinos:

$$m_{ij} = 2h_{ij}\langle\Delta^0\rangle = \sqrt{2}h_{ij}v_\Delta. \tag{5}$$

Note that the HTM is free from a massless Goldstone boson (Majoron) arising from the violation of the lepton number (L) global symmetry, because the Higgs potential contains the term $\mu\Phi\Delta\Phi$ term which explicitly violates lepton number when Δ is assigned $L = -2$. Cosmological data provides a constraint on the neutrino masses m_i, $\Sigma m_i \lesssim 0.75$ eV [28]. Lepton flavor violating (LFV) processes involving μ and τ provide the strongest upper limits on h_{ij} and hence v_Δ cannot be arbitrarily small if the HTM is to accommodate the currently favored form of the neutrino mass matrix. A rough lower bound $v_\Delta \gtrsim 10$ eV can be derived. An upper limit on v_Δ can be obtained from considering its effect on ρ. In the HTM ρ is given by (where $x = v_\Delta/v$):

$$\rho \equiv 1 + \delta\rho = \frac{1 + 2x^2}{1 + 4x^2}. \tag{6}$$

From the measurement of $\rho \approx 1$ a purely tree level analysis gives the bound $v_\Delta/v \lesssim 0.03$. We will comment on the one-loop expression for $\delta\rho$ below [29–31]. In this paper we will assume

$$10\ \text{eV} \lesssim v_\Delta \lesssim 10\,000\ \text{eV}. \tag{7}$$

Hence the tree level value of ρ is essentially equal to 1, thus easily satisfying the experimental constraint on $\delta\rho$. Such small values of v_Δ can be explained by a two-loop mechanism [14] or in the context of extra dimensions [32,33]. Moreover, such values of v_Δ would permit some h_{ij} to be sufficiently large to enhance various LFV μ and τ decays to the sensitivity of current and forthcoming experiments [14,15,17] and also are consistent with the requirement that any primordially generated baryon asymmetry is not erased by the lepton number violating triplet interactions [34].

The HTM has seven Higgs bosons (H^{++}, H^{--}, H^+, H^-, H^0, A^0, h^0). While $H^{\pm\pm}$ is purely triplet ($= \Delta^{\pm\pm}$), the remaining eigenstates would be in general mixtures of the doublet and triplet fields. Such mixing is proportional to the triplet VEV, and hence small *even if* v_Δ assumes its largest value of a few GeV. Therefore the first six eigenstates are essentially composed of triplet fields, while the $I = 1/2$ doublet gives rise to a SM like h^0 and the Goldstone bosons $G^\pm$, G^0. The most striking signature of the HTM would be the observation of $H^{\pm\pm}$.[2] In the HTM there exists the following relationships among the masses of the physical Higgs bosons:

$$\begin{gathered} m_{H^{\pm\pm}}^2 \simeq M^2 + 2\frac{(\lambda_4 - \lambda_5)}{g^2}M_W^2, \\ m_{H^\pm}^2 \simeq m_{H^{\pm\pm}}^2 + 2\frac{\lambda_5}{g^2}M_W^2, \qquad m_{H^0,A^0}^2 \simeq m_{H^\pm}^2 + 2\frac{\lambda_5}{g^2}M_W^2. \end{gathered} \tag{8}$$

[1]Note that the analogous term for a $Y = 0$ triplet is forbidden by gauge invariance.

[2]The dominantly triplet eigenstates $H^\pm$, H^0, and A^0 can have a different phenomenology to the analogous Higgs bosons in doublet ($I = 1/2$, $Y = 1$) representations.

Here M is the triplet mass term, while λ_4, λ_5 are dimensionless quartic couplings. For $\lambda_5 > 0$ ($\lambda_5 < 0$) one has the following hierarchy $m_{H^{\pm\pm}} < m_{H^{\pm}} < m_{H^0,A^0}$ ($m_{H^{\pm\pm}} > m_{H^{\pm}} > m_{H^0,A^0}$). Clearly M sets the scale for the mass of the triplet fields, while the mass splitting among the eigenstates is determined by the quartic couplings and can be $\mathcal{O}(M_W)$. We will focus on Higgs boson masses of interest for the Tevatron and LHC, and hence we assume $M \gtrsim 1$ TeV.

At the one-loop level the Higgs sector contribution to $\delta\rho$ is a function of v_Δ and the Higgs boson masses. Although a quantitative analysis in the context of the HTM is still lacking, explicit formulas for the contributions of $Y = 2$ triplets to the self-energies of the W and Z in the context of L-R symmetric models and little Higgs models can be found in [29,30]. In particular, such contributions are sensitive to the mass splittings of the Higgs bosons. In the HTM the triplet Higgs boson mass splitting is determined by the quartic coupling λ_5, with $\lambda_5 = 0$ giving rise to degenerate triplet scalars of mass M. We will present results for both the degenerate case and for mild splittings of up to 20 GeV in our discussion of $H^{\pm\pm}$ phenomenology at the Tevatron.

We now briefly discuss present mass bounds on the Higgs bosons of the HTM, which differ in some cases from the commonly quoted mass bounds in the two Higgs doublet model (2HDM). If H^0 and A^0 were the lightest, they could have been produced at LEP via the mechanism $e^+e^- \to A^0H^0$ (note that $e^+e^- \to ZH^0$ is proportional to v_Δ and hence negligible). However, since A^0 and H^0 would both decay invisibly to $\nu\bar{\nu}$, Ref. [35] suggested using LEP data on $\gamma\nu\bar{\nu}$ (where γ arises from bremsstrahlung from e^+ or e^-) and derived the mass limit $m_{H^0,A^0} \gtrsim 55$ GeV. Concerning $H^{\pm}$, LEP searched for $H^{\pm} \to cs$ or $\tau\nu_\tau$, which are expected to be the dominant decays in doublet models, and obtained mass limits around $m_{H^{\pm}} \gtrsim 80$ GeV. For the triplet $H^{\pm}$ the decays $H^{\pm} \to e^{\pm}\nu, \mu^{\pm}\nu$ may have large branching ratios. However, in this scenario one presumably could use data from slepton searches $e^+e^- \to \tilde{l}^+\tilde{l}^- \to l^+l^-\chi^0\overline{\chi}^0$ to derive similar mass limits ($\gtrsim 80$ GeV) [36]. A recent quantitative analysis of the above decays in the context of a little Higgs model can be found in [37].

Concerning $H^{\pm\pm}$, LEP searched for both left-handed $H_L^{\pm\pm}$ and right-handed $H_R^{\pm\pm}$ (which we will not consider in this paper) via several mechanisms:

(i) Pair production via $e^+e^- \to \gamma^*$, $Z^* \to H^{++}H^{--}$ followed by decay to $l^+l^+l^-l^-$ ($l^{\pm} = e^{\pm}, \mu^{\pm}, \tau^{\pm}$); the cross section is determined by gauge couplings and leads to mass limits of $m_{H^{\pm\pm}} > 100$ GeV [8–10].

(ii) Single production of $H^{\pm\pm}$ via $e^+e^- \to H^{\pm\pm}e^{\mp}e^{\mp}$; the rate is determined by the coupling h_{11} and leads to excluded regions in the plane (h_{11}, $m_{H^{\pm\pm}}$), with sensitivity up to $m_{H^{\pm\pm}} \lesssim 180$ GeV. Limits of $10^{-2} \to 10^{-1}$ were set on h_{11} [11].

(iii) The effect of $H^{\pm\pm}$ on Bhabha scattering $e^+e^- \to e^+e^-$; as in (ii) above this leads to excluded regions in the plane (h_{11}, $m_{H^{\pm\pm}}$) [10,11], with sensitivity up to $m_{H^{\pm\pm}} \lesssim 2$ TeV. Limits of $10^{-2} \to 10^{-1}$ were set on h_{11}.

The direct searches for $H^{\pm\pm}$ will continue at the hadron colliders, Tevatron and LHC.

III. PRODUCTION OF $H^{\pm\pm}$ AT THE TEVATRON

A distinct signature of $H^{\pm\pm}$ would be a pair of same sign charged leptons ($e^{\pm}$ or $\mu^{\pm}$) with high invariant mass. At hadron colliders such a signal has a relatively high detection efficiency and enjoys essentially negligible background from standard model processes. Earlier theoretical studies of the search potential for $H^{\pm\pm}$ at such colliders can be found in [20,21], with a recent analysis at the LHC in [38]. The decays of $H^{\pm\pm}$ to states involving $\tau^{\pm}$ are more problematic at hadron colliders, although simulations in these channels [21,38] promise sensitivity to values of $m_{H^{\pm\pm}}$ beyond the LEP limits. The decays $H^{\pm\pm} \to W^{\pm}W^{\pm}$ are proportional to v_Δ and can be neglected in the case of very small v_Δ of interest to us.

In 2003 the Tevatron performed the first search for $H^{\pm\pm}$ at a hadron collider. D0 Collaboration [18] has searched for $H^{\pm\pm} \to \mu^+\mu^-$ while CDF Collaboration [19] searched for three final states: $H^{\pm\pm} \to e^{\pm}e^{\pm}$, $e^{\pm}\mu^{\pm}$, $\mu^{\pm}\mu^{\pm}$. The assumed production mechanism for $H^{\pm\pm}$ is $q\overline{q} \to \gamma^*$, $Z^* \to H^{++}H^{--}$.[3] This cross section depends on only one unknown parameter, $m_{H^{\pm\pm}}$, and importantly is not suppressed by any small factor such as a Yukawa coupling h_{ij} or a triplet VEV. The search assumes that $H^{\pm\pm}$ is sufficiently long-lived to decay in the detector, which corresponds to $h_{ll} > 10^{-5}$. A search for a long-lived $H^{\pm\pm}$ decaying outside the detector has been performed in [39]. The cross section also depends on the hypercharge of the Higgs representation, which is $Y = 2$ in the HTM. This value of Y is assumed also in the experimental searches. The explicit partonic cross section at leading order (LO) is as follows (where $q = u, d$):

$$\sigma_{\rm LO}(q\overline{q} \to H^{++}H^{--}) = \frac{\pi\alpha^2}{9Q^2}\beta_1^3\left[e_q^2e_H^2 + \frac{e_qe_H{\rm v}_q{\rm v}_H(1 - M_Z^2/Q^2) + ({\rm v}_q^2 + {\rm a}_q^2){\rm v}_H^2}{(1 - M_Z^2/Q^2)^2 + M_Z^2\Gamma_Z^2/Q^4}\right]. \tag{9}$$

Here ${\rm v}_q = (I_{3q} - 2e_qs_W^2)/(s_Wc_W)$, ${\rm a}_q = I_{3q}/(s_Wc_W)$, and ${\rm v}_H = (I_{3H} - e_Hs_W^2)/(s_Wc_W)$. The third isospin component is denoted by I_{3q} (I_{3H}) and e_q (e_H) is the electric charge of the quark q ($H^{\pm\pm}$). s_W and c_W are $\sin\theta_W$ and $\cos\theta_W$, respectively. Q^2 is the partonic center-of-mass energy. α is the QED coupling evaluated at the scale Q, M_Z is the Z

[3]The model-dependent contribution from any Z' (which can enhance the cross section [24,38]) is currently not considered.

boson mass, Γ_Z is the Z boson width, and $\beta_1 = \sqrt{1 - 4m_{H^{\pm\pm}}^2/Q^2}$. Order α_s QCD corrections modify the LO cross section by a factor $K \approx 1.3$ at the Tevatron for $m_{H^{\pm\pm}} < 200$ GeV and $K \approx 1.25$ at the LHC for $m_{H^{\pm\pm}} < 1000$ GeV [40]. We neglect the gluon-gluon fusion (α_s^2) contribution to $H^{++}H^{--}$ production, which has no compensatory enhancement factor analogous to the $\tan^4\beta$ term for doublet $H^{\pm}$ production via $gg \to H^+H^-$ [41].

Assuming that $H^{\pm\pm}$ production proceeds via this pair production process, the absence of signal enables a limit to be set on the product:

$$\sigma(p\overline{p} \to H^{++}H^{--})\mathrm{BR}(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm}). \qquad (10)$$

Clearly the strongest constraints on $m_{H^{\pm\pm}}$ are obtained assuming $\mathrm{BR}(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm}) = 100\%$. Currently these mass limits stand at 133, 115, 136 GeV for the $e^{\pm}e^{\pm}$, $e^{\pm}\mu^{\pm}$, $\mu^{\pm}\mu^{\pm}$ channels, respectively [19]. In the HTM one expects $\mathrm{BR}(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm}) \neq 100\%$ if Eq. (5) is required to explain the currently favored form of the neutrino mass matrix [14].

The current search strategy is in fact sensitive to any *singly produced* $H^{\pm\pm}$, i.e. signal candidates are events with *one pair* of same sign leptons reconstructing to $m_{H^{\pm\pm}}$. This requirement is sufficient to reduce the SM background to negligible proportions. Hence the search potential of the Tevatron merely depends on the signal efficiencies for the signal (currently $\approx 34\%$, 34%, 18% for $\mu\mu$, ee, $e\mu$) and the integrated luminosity. With these relatively high efficiencies and an expected $\mathcal{L} = 4$–8 fb^{-1} by the year 2009, discovery with >5 events will be possible for $\sigma_{H^{++}H^{--}}$ of a few fb, which corresponds to a mass reach $m_{H^{\pm\pm}} < 200$ GeV.

Although single $H^{\pm\pm}$ production processes such as $p\overline{p} \to W^{\pm} \to W^{\mp}H^{\pm\pm}$ can be neglected[4] due to the strong triplet VEV suppression, the mechanism $p\overline{p} \to W^* \to H^{\pm\pm}H^{\mp}$ is potentially sizeable. This latter process proceeds via a gauge coupling constant and is not suppressed by any small factor. The LO partonic cross section is as follows:

$$\sigma_{\mathrm{LO}}(q'\overline{q} \to H^{++}H^{-}) = \frac{\pi\alpha^2}{144 s_W^4 Q^2} C_T^2 p_W^2 \beta_2^3. \qquad (11)$$

Here C_T arises from the $H^{\pm\pm}H^{\mp}W^{\mp}$ vertex and $C_T = 2$ for $I = 1$, $Y = 2$ triplet fields (the doublet component of $H^{\pm}$ is negligible); $\beta_2 = \sqrt{(1-(m_{H^{\pm}}+m_{H^{\pm\pm}})^2/Q^2)(1-(m_{H^{\pm}}-m_{H^{\pm\pm}})^2/Q^2)}$ and $p_W = Q^2/(Q^2 - M_W^2)$. For simplicity, we take the same $K = 1.3$ as for $\sigma(q\overline{q} \to H^{++}H^{--})$ at the Tevatron and $K = 1.25$ at the LHC. Explicit calculations for the K factor for the process $\sigma(q'\overline{q} \to H^{\pm}A^0)$ in the minimal supersymmetric standard model (MSSM) [44] (which shares the same K factor as $q'\overline{q} \to H^{++}H^{-}$) give $K \approx 1.2$. In this paper we will study in detail the magnitude and relative importance of $\sigma(q'\overline{q} \to H^{\pm\pm}H^{\mp})$. Although we work in the HTM, our numerical analysis is relevant for other models which possess a $I = 1$, $Y = 2$ Higgs triplet (e.g. L-R symmetric models and little Higgs models).

[4] Single production of a right-handed triplet via $q'\overline{q} \to W_R^{\pm} \to W_R^{\mp}H^{\pm\pm}$ [42] and $W_R^{\pm}W_R^{\pm}$ fusion [43] can be sizeable at the LHC.

A previous quantitative study of this mechanism can be found in [24]. Cross sections were given at both LHC and Tevatron energies for $m_{H^{\pm\pm}} > 200$ GeV with the simplifying assumption $m_{H^{\pm\pm}} = m_{H^{\pm}}$. It was shown that $\sigma(q'\overline{q} \to H^{\pm\pm}H^{\mp})$ can be of comparable size to $\sigma(q\overline{q} \to H^{++}H^{--})$.

In this paper we first generalize the work of [24] as follows:

(i) In our discussion at the Tevatron we consider masses in the range 100 GeV $< m_{H^{\pm\pm}} <$ 200 GeV which will be probed during run II and allow mild mass splittings $|m_{H^{\pm\pm}} - m_{H^{\pm}}| \leq 20$ GeV.

(ii) In our discussion at the LHC we consider larger mass splittings $|m_{H^{\pm\pm}} - m_{H^{\pm}}| \leq 80$ GeV.

(iii) For both the Tevatron and LHC we study in detail the relative magnitude of $\sigma(q'\overline{q} \to H^{\pm\pm}H^{\mp})$ and $\sigma(q\overline{q} \to H^{++}H^{--})$.

Moreover, motivated by the fact that the currently employed Tevatron search strategy is sensitive to *single production* of $H^{\pm\pm}$, we advocate the use of the inclusive single production cross section ($\sigma_{H^{\pm\pm}}$) when comparing the experimentally excluded region with the theoretical cross section. This leads to a strengthening of the mass bound for $m_{H^{\pm\pm}}$ which now carries a dependence on $m_{H^{\pm}}$. We introduce the single production cross section as follows:

$$\sigma_{H^{\pm\pm}} = \sigma(p\overline{p}, pp \to H^{++}H^{--}) + \sigma(p\overline{p}, pp \to H^{++}H^{-}) + \sigma(p\overline{p}, pp \to H^{--}H^{+}). \qquad (12)$$

At the Tevatron $\sigma(p\overline{p} \to H^{++}H^{-}) = \sigma(p\overline{p} \to H^{--}H^{+})$ while at the LHC $\sigma(pp \to H^{++}H^{-}) > \sigma(pp \to H^{--}H^{+})$. If a signal for $H^{\pm\pm}$ were found in the two-lepton channel, subsequent searches could select signal events with three or four leptons, in order to disentangle $q\overline{q} \to H^{++}H^{--}$ and $q'\overline{q} \to H^{\pm\pm}H^{\mp}$. In our numerical analysis we utilize the CTEQ6L1 parton distribution functions [45]. We take the factorization scale (Q) as the partonic center-of-mass energy ($\sqrt{s}$). Our results for $\sigma(q\overline{q} \to H^{++}H^{--})$ agree with those in [21,40]. Our results for $\sigma(q'\overline{q} \to H^{\pm\pm}H^{\mp})$ agree with those in [24] (and taking $C_T = 1$ agree with $\sigma(q'\overline{q} \to H^{\pm}A^0)$ in the 2HDM/MSSM [44]). The above cross sections evaluated with MRST02 parton distribution functions [46] agree with those evaluated with CTEQ6L1 to within 10%–15%.

In Fig. 1(a) we plot $\sigma_{H^{\pm\pm}}$ as a function of $m_{H^{\pm\pm}}$ at the Tevatron for three different values of $m_{H^{\pm}}$. We take $K = 1.3$. The current excluded regions from the $e^{\pm}e^{\pm}$, $e^{\pm}\mu^{\pm}$,

$\mu^\pm\mu^\pm$ searches correspond to the area above horizontal lines at roughly 40, 70, 35 fb, respectively. The present mass limits for $m_{H^{\pm\pm}}$ are where the curve for $H^{++}H^{--}$ intersects with the above horizontal lines and read as 133, 115, 136 GeV, respectively, for $\mathrm{BR}(H^{\pm\pm} \to l_i^\pm l_j^\pm) = 100\%$. With the inclusion of the $H^{\pm\pm}H^{\mp}$ channel, these mass limits increase to 150, 130, 150 for $m_{H^\pm} = m_{H^{\pm\pm}} + 20$ GeV, strengthening to 160, 140, 160 for $m_{H^\pm} = m_{H^{\pm\pm}} - 20$ GeV. Clearly the search potential of the Tevatron (i.e. the mass limit on $m_{H^{\pm\pm}}$) increases significantly when one includes the contribution to $\sigma_{H^{\pm\pm}}$ from $p\overline{p} \to H^{\pm\pm}H^{\mp}$. Note that the above mass limits strictly apply to the case when $H^{\pm\pm}$ decays leptonically and with BR = 100% in a given channel. However, if h_{ij} are to provide the currently favored form of the neutrino mass matrix then $\mathrm{BR}(H^{\pm\pm} \to l_i^\pm l_j^\pm) < 100\%$ in a given channel. Moreover, if $m_{H^{\pm\pm}} > m_{H^\pm}$ then the decay channel $H^{\pm\pm} \to H^\pm W^*$ would be open. As shown in [22], this decay can be sizeable and thus reduces $\mathrm{BR}(H^{\pm\pm} \to l_i^\pm l_j^\pm)$. We will return to these issues in Sec. V.

In Fig. 1(b) we plot the ratio of cross sections R at the Tevatron as a function of $m_{H^{\pm\pm}}$, where R is defined as follows:

$$R \equiv \frac{\sigma(p\overline{p}, pp \to H^{++}H^{-}) + \sigma(p\overline{p}, pp \to H^{--}H^{+})}{\sigma(p\overline{p}, pp \to H^{++}H^{--})}. \tag{13}$$

The $m_{H^{\pm\pm}}$ dependence arises from the phase space functions β_1 and β_2 in Eqs. (9) and (11). As can be seen, $0.8 < R < 2.2$ and thus $q'\overline{q} \to H^{\pm\pm}H^{\mp}$ contributes significantly to $\sigma_{H^{\pm\pm}}$.

In Fig. 2 we plot the analogies of Fig. 1 for the LHC. In Fig. 2(a) we plot $\sigma_{H^{\pm\pm}}$ for three values of $m_{H^{\pm\pm}}$ and for larger mass splittings ($|m_{H^{\pm\pm}} - m_{H^\pm}| \le 80$ GeV) than in Fig. 2. We take $K = 1.25$. As before, the inclusion of $q'\overline{q} \to H^{\pm\pm}H^{\mp}$ significantly increases the search potential e.g. if sensitivity to $\sigma_{H^{\pm\pm}} = 1$ fb is attained, the mass reach extends from $m_{H^{\pm\pm}} < 600$ GeV ($H^{++}H^{--}$ only) to 750 GeV for ($m_{H^\pm} = m_{H^{\pm\pm}} - 80$ GeV). Recently [38] performed a simulation of the detection prospects at the LHC for $q\overline{q} \to H^{++}H^{--}$ for the cases where three and four leptons are detected. With 100 fb^{-1}, sensitivity to $m_{H^{\pm\pm}} \lesssim 800$ GeV (three leptons) and $m_{H^{\pm\pm}} \lesssim 700$ GeV (four leptons) is expected. We are not aware of a simulation for the case where only two leptons are detected. Presumably even larger values of $m_{H^{\pm\pm}}$ ($\gtrsim 800$ GeV) could be probed. In Fig. 2(b) we plot R as a function of $m_{H^{\pm\pm}}$. One can see that $R > 1$ for the upper two curves for all $m_{H^{\pm\pm}}$, while for the lower curve $R > 1$ for $m_{H^{\pm\pm}} > 260$ GeV. Note that the dependence of R on $m_{H^{\pm\pm}}$ differs from that observed in Fig. 1(b), which can be attributed to the different parton luminosity functions at the Tevatron and LHC.

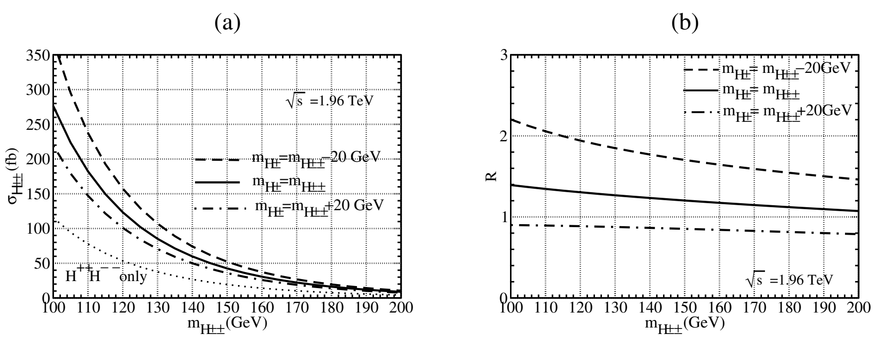

FIG. 1. (a) Single production cross section of $H^{\pm\pm}$ ($\sigma_{H^{\pm\pm}}$) at the Tevatron as a function of $m_{H^{\pm\pm}}$ for different values of $m_{H^\pm}$. (b) Ratio R as a function of $m_{H^{\pm\pm}}$. We use CTEQ6L1 parton distribution functions.

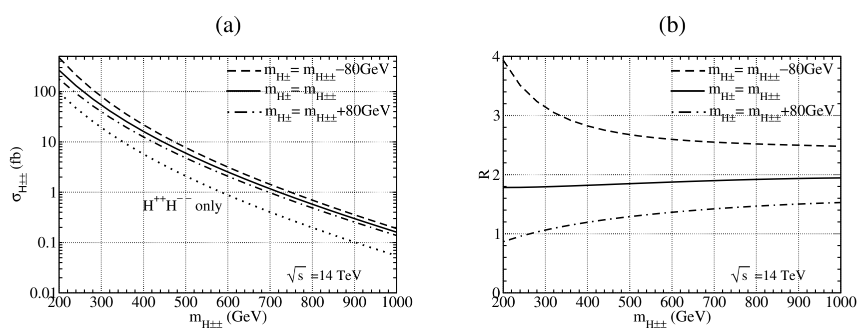

FIG. 2. (a) Single production cross section of $H^{\pm\pm}$ ($\sigma_{H^{\pm\pm}}$) at the LHC as a function of $m_{H^{\pm\pm}}$ for different values of $m_{H^\pm}$. (b) Ratio R as a function of $m_{H^{\pm\pm}}$. We use CTEQ6L1 parton distribution functions.

IV. NEUTRINO MASS HIERARCHY AND THE DECAY $H^{\pm\pm} \to H^{\pm}W^*$

The current experimental searches assume that the sole decay mode of $H^{\pm\pm}$ is $H^{\pm\pm} \to l_i^{\pm} l_j^{\pm}$ mediated by the arbitrary Yukawa couplings h_{ij}. The decay rate for $H^{\pm\pm} \to l_i^{\pm} l_j^{\pm}$ is given by

$$\Gamma(H^{\pm\pm} \to l_i^{\pm} l_j^{\pm}) = S\frac{m_{H^{\pm\pm}}}{8\pi}|h_{ij}|^2, \tag{14}$$

where $S = 1(2)$ for $i = j$ ($i \neq j$). Clearly $\Gamma(H^{\pm\pm} \to l_i^{\pm} l_j^{\pm})$ depends crucially on the *absolute* value of the h_{ij}, although the leptonic BRs are determined by the *relative* values. In this section we consider the impact of the decay mode $H^{\pm\pm} \to H^{\pm}W^*$ on the BRs of the leptonic channels. It has been known for some time that $\mathrm{BR}(H^{\pm\pm} \to H^{\pm}W^*)$ is potentially sizeable and a quantitative analysis can be found in [22]. The decay rate for $H^{\pm\pm} \to H^{\pm}W^*$ (summing over all fermion states for $W^* \to ff$ excluding the t quark) is given by

$$\Gamma(H^{\pm\pm} \to H^{\pm}W^*) = 9G_F^2 M_W^4 m_{H^{\pm\pm}} C_T^2 P/(16\pi^3), \tag{15}$$

where P is the phase space term (which we calculate by numerical integration) and $C_T(=2)$ is from the coupling $H^{\pm\pm}H^{\pm}W$. P depends on the mass difference Δm defined by $\Delta m = m_{H^{\pm\pm}} - m_{H^{\pm}}$, and $P = 0$ for $\Delta m = 0$. If $m_{H^{\pm}} < m_{H^{\pm\pm}}$ this decay can compete with $H^{\pm\pm} \to l_i^{\pm} l_j^{\pm}$ since the phase space suppression of the virtual W^* is compensated by the gauge strength coupling [21]. Reference [22] showed that $H^{\pm\pm} \to H^{\pm}W^*$ can dominate over $H^{\pm\pm} \to l_i^{\pm} l_j^{\pm}$ if Δm is sizeable (> 40 GeV) and h_{ij} are of order 10^{-3} or less. A large $\mathrm{BR}(H^{\pm\pm} \to H^{\pm}W^*)$ would debilitate the $H^{\pm\pm}$ search potential in the leptonic channel. However, as emphasized in [23], observation of $H^{\pm\pm} \to H^{\pm}W^*$ together with one or more of the leptonic channels could provide information on the absolute values of h_{ij}. If only $\mathrm{BR}(H^{\pm\pm} \to l_i^{\pm} l_j^{\pm})$ are measured then only the *relative* values of the h_{ij} can be evaluated. The decay rate for $H^{\pm\pm} \to H^{\pm}W^*$ is theoretically calculable once $m_{H^{\pm}}$ and $m_{H^{\pm\pm}}$ are known experimentally, and thus it can be used as a benchmark decay with which to estimate the total width of $H^{\pm\pm}$. It is known that the BRs of the leptonic channels depend on which solution to the neutrino mass matrix is realized [14]. However, a quantitative analysis of the impact of $H^{\pm\pm} \to H^{\pm}W^*$ in the various allowed scenarios is still lacking and will be presented below. We are not aware of any experimental simulation of $H^{\pm\pm} \to H^{\pm}W^*$. The signature would depend crucially on the decay products of $H^{\pm}$, which are either $H^{\pm} \to l^{\pm}\nu_l$ (driven by h_{ij}), or possibly $H^{\pm} \to H^0W^*$, A^0W^*.

We now briefly review relevant results and formulas from neutrino physics. The neutrino mass matrix is diagonalized by the MNS (Maki-Nakagawa-Sakata) matrix V_{MNS} [47]. Using Eq. (5) one can write the couplings h_{ij} as follows:

$$h_{ij} = \frac{1}{\sqrt{2}v_\Delta} V_{\mathrm{MNS}} \mathrm{diag}(m_1, m_2, m_3) V_{\mathrm{MNS}}^T. \tag{16}$$

Here we take the basis in which the unitary matrix responsible for diagonalizing the charged-lepton mass matrix is a unit matrix. The MNS matrix in the standard parametrization is as follows:

$$V_{\mathrm{MNS}} = \begin{pmatrix} c_1c_3 & s_1c_3 & s_3e^{-i\delta} \\ -s_1c_2 - c_1s_2s_3e^{i\delta} & c_1c_2 - s_1s_2s_3e^{i\delta} & s_2c_3 \\ s_1s_2 - c_1c_2s_3e^{i\delta} & -c_1s_2 - s_1c_2s_3e^{i\delta} & c_2c_3 \end{pmatrix} \times \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i\varphi_1/2} & 0 \\ 0 & 0 & e^{i\varphi_2/2} \end{pmatrix}, \tag{17}$$

where $s_i \equiv \sin\theta_i$ and $c_i \equiv \cos\theta_i$, δ is the Dirac phase, and φ_1 and φ_2 are the Majorana phases.

Neutrino oscillation experiments involving solar [48], atmospheric [49], and reactor neutrinos [50] are sensitive to the mass-squared differences and the mixing angles and give the following preferred values:

$$\Delta m_{12}^2 \equiv m_2^2 - m_1^2 \simeq 8.0 \times 10^{-5}\ \mathrm{eV}^2, \quad |\Delta m_{13}^2| \equiv |m_3^2 - m_1^2| \simeq 2.1 \times 10^{-3}\ \mathrm{eV}^2, \tag{18}$$

$$\sin^2 2\theta_1 \simeq 0.8, \qquad \sin^2 2\theta_2 \simeq 1, \qquad \sin^2 2\theta_3 \lesssim 0.16. \tag{19}$$

Since the sign of Δm_{13}^2 and the mass of the lightest neutrino are both undetermined at present, distinct neutrino mass hierarchy patterns are classified as follows: *Normal hierarchy* (NH) ($m_1 < m_2 \ll m_3$), *Inverted hierarchy* (IH) ($m_2 > m_1 \gg m_3$), *Quasidegenerate* (DG) ($m_1 \sim m_2 \sim m_3 \gg \sqrt{|\Delta m_{13}^2|}$). From Eq. (5) and Eq. (16) it can be shown that

$$\sum_{i,j} h_{ij}^2 v_\Delta^2 \propto \sum_i m_i^2. \tag{20}$$

Hence the total leptonic decay width depends on the absolute mass of the neutrinos, and the value of $\sum_i m_i^2$ depends on which solution to the neutrino mass matrix (NH, IH, DG) is realized. The minimum value of $\sum_i m_i^2$ is $|\Delta m_{13}^2|$ while the maximum is given by the cosmological constraint.

In Fig. 3 we show contours of $\mathrm{BR}(H^{\pm\pm} \to H^{\pm}W^*)$ in the plane $(m_{H^{\pm\pm}}, v_\Delta)$ for three different solutions to the neutrino mass matrix. We assume that $m_{1(3)} = 0$ for NH (IH) and $m_1 = 0.2$ eV for DG. We take $m_{H^{\pm}} = m_{H^{\pm\pm}} - 20$ GeV. From Eq. (16), all h_{ij} are determined once v_Δ is specified. In order to comply with current experimental upper limits on LFV decays of $\mu^{\pm}$ and $\tau^{\pm}$, one can derive the bound $v_\Delta > 10$ eV for NH and IH, and $v_\Delta > 100$ eV

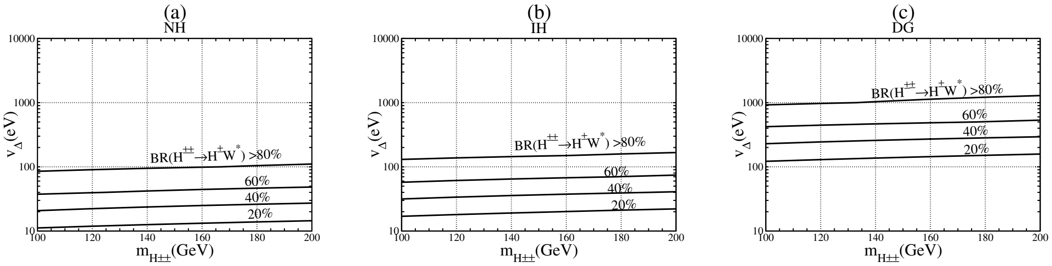

FIG. 3. Contours of BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) in the plane ($m_{H^{\pm\pm}}$, v_Δ), for NH (a), IH (b), and DG (c). We take $m_{H^\pm} = m_{H^{\pm\pm}} - 20$ GeV.

for DG. The stronger constraint on v_Δ in DG arises because $\sum_i m_i$ in DG is larger than those in NH and IH.

From Fig. 3 it is clear that BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) can be sizeable and approaches 100% for larger v_Δ. For a fixed value of v_Δ, one can see that BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) is relatively more important in NH and IH than in DG. This can be understood from Eq. (20), since DG requires heavier neutrinos (and thus larger h_{ij}) which in turn reduces BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$). One can consider three distinct scenarios with very different magnitudes for BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) and BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$):

(i) BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) $\gg$ BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$).—In this case the current search strategy (which requires $H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$ decay) is ineffective. Simulations have not been carried out for the decay $H^{\pm\pm} \rightarrow H^{\pm} W^*$ although one might naïvely expect sensitivity comparable to that for the decay $H^{\pm\pm} \rightarrow \tau^{\pm}\tau^{\pm}$, as suggested in [21,23].

(ii) BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) $\approx$ BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$).—The search for $H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$ would be effective and $H^{\pm\pm}$ could be discovered in one or more leptonic channels. If $H^{\pm\pm} \rightarrow H^{\pm} W^*$ is also observed then information on the absolute value of h_{ij} might be possible: Using Eqs. (14) and (15), the ratio of leptonic events ($N_{l_i l_j}$) to $H^{\pm} W^*$ events ($N_{H^\pm W^*}$) is given as follows:

$$\frac{N_{l_i l_j}}{N_{H^\pm W^*}} \sim \frac{h_{ij}^2}{P}. \tag{21}$$

Observation of the leptonic channel provides $m_{H^{\pm\pm}}$. If $m_{H^\pm}$ can be roughly measured then P (and hence the partial width for $H^{\pm\pm} \rightarrow H^{\pm} W^*$) can be calculated. From the above equation one can obtain an order of magnitude estimate of h_{ij}.

(iii) BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) $\ll$ BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$).—In this case the current search strategy ($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$) is effective. If BR($H^{\pm\pm} \rightarrow l_i^{\pm} l_j^{\pm}$) are measured then the ratios of h_{ij} can be evaluated. This can be compared with Eq. (5) in order to see which neutrino solution is realized [14]. The absolute values of h_{ij} cannot be measured unless a LFV decay of μ and/or τ is observed.

V. TEVATRON SEARCH POTENTIAL IN HTM

We now study the search potential of the Tevatron for the generalized case in the HTM where $p\overline{p} \rightarrow H^{\pm\pm} H^{\mp}$ is included, BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$) $\neq 0\%$ and h_{ij} are required to reproduce a phenomenologically acceptable neutrino mass matrix. We relax the assumptions for the Majorana phases and take $\varphi_1, \varphi_2 = 0$ or π, which leads the seven distinct solutions:

NH: $m_1 < m_2 \ll m_3$	
IH1: $m_2 > m_1 \gg m_3$	IH2: $-m_2 > m_1 \gg m_3$
DG1: $m_1 \simeq m_2 \simeq m_3$	DG2: $m_1 \simeq m_2 \simeq -m_3$
DG3: $m_1 \simeq -m_2 \simeq m_3$	DG4: $m_1 \simeq -m_2 \simeq -m_3$

In the HTM, BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$) are predicted and different in each of the seven distinct solutions (NH, IH1, IH2, DG1–DG4), and their ratios were evaluated in [14]. Note that such predictions of BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$) are a feature of the HTM in which the couplings h_{ij} are the sole origin of neutrino mass. This direct correlation between BR($H^{\pm\pm} \rightarrow l^{\pm} l^{\pm}$) and the neutrino mass matrix may not extend to $H^{\pm\pm}$ of other models in which neutrinos can acquire mass by other means e.g. the seesaw mechanism in L-R models or by a combination of mechanisms which may or may not include the h_{ij} couplings [37,51]. In contrast, the production process $\sigma(p\overline{p} \rightarrow H^{\pm\pm} H^{\mp})$ is certainly relevant in any model with $Y = 2$ triplets.

In Figs. 4–6 we plot σ_{ll} as a function of $m_{H^{\pm\pm}}$, where σ_{ll} is the total leptonic ($l = e, \mu, \tau$) cross section defined as

$$\sigma_{ll} = \sigma(p\overline{p} \rightarrow H^{++}H^{--}) B_{ll}(2 - B_{ll}) + 2\sigma(p\overline{p} \rightarrow H^{++}H^{-}) B_{ll}. \tag{22}$$

The contribution to σ_{ll} from $\sigma(p\overline{p} \rightarrow H^{++}H^{--})$ falls more slowly with decreasing B_{ll} since signal candidates are events with at least 2 leptons. Equation (22) simplifies to Eq. (10) in the limit where $\sigma(p\overline{p} \rightarrow H^{\pm\pm} H^{\mp}) = 0$ and $B_{ll} = 1$. Figure 4(a) shows σ_{ll} for the NH with $m_{H^\pm} = m_{H^{\pm\pm}}$, which leads to $B_{ll} = 1$. In this case $\sum \sigma_{ll} = \sigma_{H^{\pm\pm}}$. For the other figures we take $m_{H^\pm} = m_{H^{\pm\pm}} - 20$ GeV, which induces a sizeable (but not dominant) BR($H^{\pm\pm} \rightarrow H^{\pm} W^*$), and hence $\sum \sigma_{ll} < \sigma_{H^{\pm\pm}}$. We set $v_\Delta = 10$ eV in Figs. 4 and 5 and $v_\Delta = 100$ eV in Figs. 6. We only plot σ_{ll} for ee, $e\mu$, $\mu\mu$ since the Tevatron already has performed

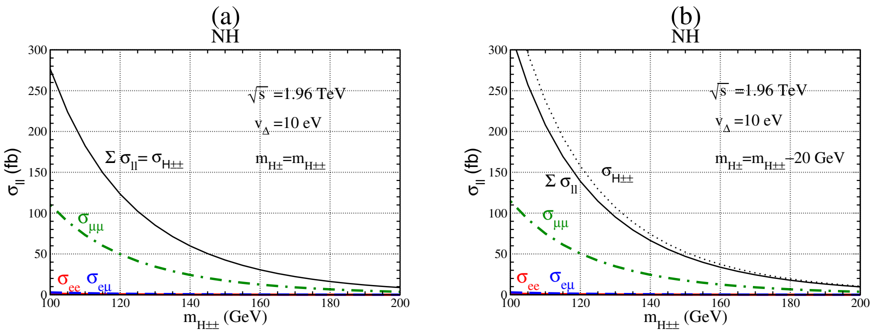

FIG. 4 (color online). σ_{ll} as a function of $m_{H^{\pm\pm}}$ for NH with (a) $m_{H^{\pm}} = m_{H^{\pm\pm}}$ and (b) $m_{H^{\pm}} = m_{H^{\pm\pm}} - 20$ GeV.

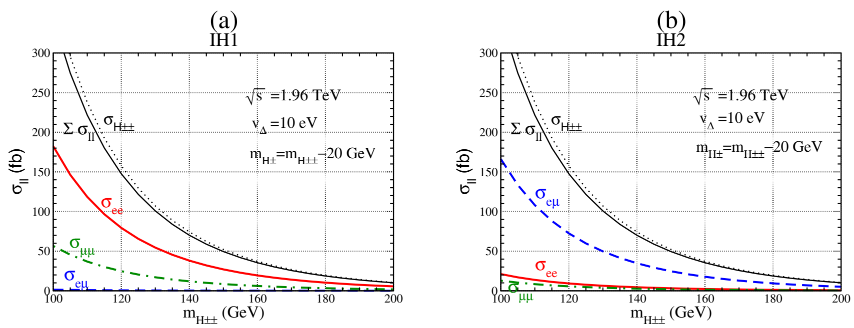

FIG. 5 (color online). σ_{ll} as a function of $m_{H^{\pm\pm}}$ for (a) IH1 and (b) IH2.

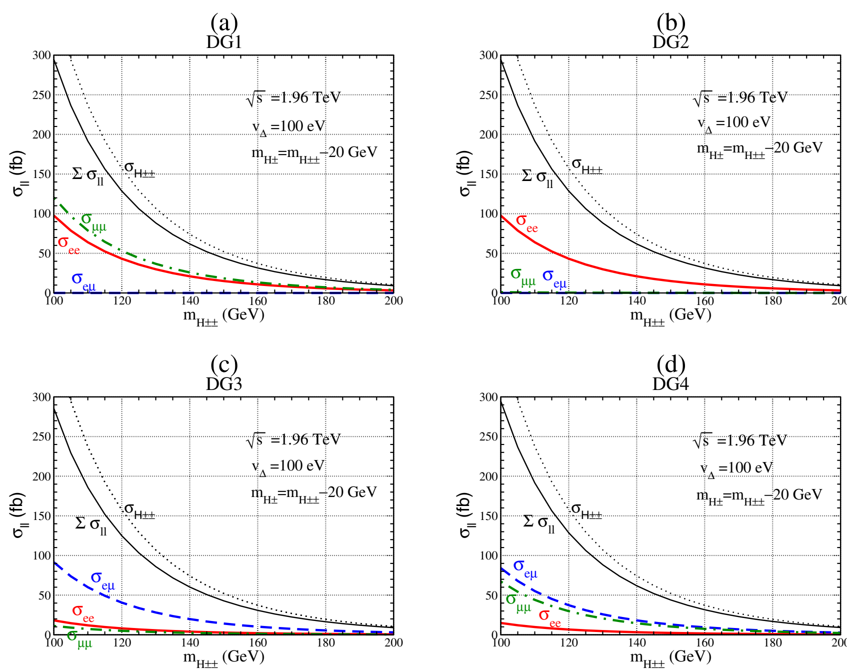

FIG. 6 (color online). σ_{ll} as a function of $m_{H^{\pm\pm}}$ for (a) DG1, (b) DG2, (c) DG3, and (d) DG4.

searches in these channels. Sensitivity to σ_{ll} of a few fb will be possible with the anticipated integrated luminosities of 4–8 fb^{-1}. There are plans to search for the three leptonic decays involving τ ($e\tau$, $\mu\tau$, $\tau\tau$) although the discovery reach in $m_{H^{\pm\pm}}$ is expected to be inferior to that for the ee, $e\mu$, $\mu\mu$ channels. In all figures we take $\theta_3 = 0°$. From the figures it is clear that $\sigma_{ee,e\mu,\mu\mu}$ differ considerably in each of the seven scenarios. Optimal coverage is for cases DG1 and DG4, which have $\sigma_{ee,\mu\mu} \geq 5$ fb and $\sigma_{e\mu,\mu\mu} \geq 5$ fb, respectively, for $m_{H^{\pm\pm}} \lesssim 180$ GeV. For NH, $\sigma_{\mu\mu} \geq 5$ fb for $m_{H^{\pm\pm}} \lesssim 190$ GeV but σ_{ee} and $\sigma_{e\mu}$ are both unobservable. Taking θ_3 at its largest experimentally allowed value results in minor changes to all figures, with the most noticeable effect being a significant reduction of $\sigma_{\mu\mu}$ in DG4. Clearly the Tevatron run II not only has strong search potential for $H^{\pm\pm}$, but is also capable of distinguishing between the various allowed scenarios for the neutrino mass matrix.

VI. CONCLUSIONS

We have studied the production of doubly charged Higgs bosons ($H^{\pm\pm}$) at hadron colliders in the Higgs triplet model (HTM), in which a complex $Y = 2$ scalar triplet is added to the standard model. The HTM can explain the observed neutrino mass matrix by invoking Yukawa couplings h_{ij} of the triplet fields to leptons. A definitive signal of the HTM would be the observation of the decay $H^{\pm\pm} \to l^{\pm}l^{\pm}$, which enjoys almost negligible background at hadron colliders, and whose branching ratios are correlated with the neutrino mass matrix. We studied the production mechanism $q'\overline{q} \to H^{\pm\pm}H^{\mp}$ which can be as large as the mechanism $q\overline{q} \to H^{++}H^{--}$ assumed in the current searches at the Tevatron. Since the present search strategy is sensitive to single production of $H^{\pm\pm}$, we advocated the use of the inclusive single production cross section ($\sigma_{H^{\pm\pm}}$) when comparing the experimentally excluded region with the theoretical cross section. This leads to a strengthening of the mass bound for $m_{H^{\pm\pm}}$, which now carries a dependence on $m_{H^{\pm}}$, and significantly improves the $H^{\pm\pm}$ search potential at the Tevatron and LHC. Although we performed our numerical analysis in the HTM, we emphasized that the introduction of $\sigma_{H^{\pm\pm}}$ also is relevant for any model which contains a $Y = 2$ Higgs triplet (e.g. L-R symmetric models and little Higgs models).

Moreover, we quantified the impact of the decay mode $H^{\pm\pm} \to H^{\pm}W^*$ for the case of a hierarchical, inverted hierarchical, and degenerate neutrino mass spectrum. On discovering a $H^{\pm\pm}$ it would be imperative to measure the absolute value of h_{ij} (and hence v_Δ) in order to reconstruct the low energy Higgs triplet Lagrangian. We stressed that an order of magnitude estimate of h_{ij} could be obtained if the channel $H^{\pm\pm} \to H^{\pm}W^*$ is observed *and* $m_{H^{\pm}}$ is roughly measured. We encourage a detailed experimental simulation of this decay mode at both the Tevatron and LHC.

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